Cover Photos

Top left: Shown in this photograph is the fourth-harmonic (263-nm) probe diagnostic system located on the upper deck of the OMEGA EP Target Area Structure. Senior Laboratory Engineer Jay Brown is shown inspecting the system.

Middle left: Installation of the single-line-of-sight time-resolved x-ray imager (SLOS-TRXI) diagnostic on the OMEGA target chamber. This was a joint project including General Atomics, Kentech Instruments, Lawrence Livermore National Laboratory, Sandia National Laboratories, and LLE. The system provides time-resolved x-ray images of the hot spot formed in cryogenic capsule implosions on OMEGA.

Bottom left: The 11 participants in the 2017 LLE Summer High-School Research Program are shown in the OMEGA EP Laser Bay. Two of the students were designated as “Scholars” in the prestigious nationwide Regeneron Science Talent Search for their work at LLE during the summer.

Top right: Graduate students and postdocs attending the Ninth Annual Omega Laser Facility Users Group Workshop. The workshop was attended by 110 researchers from five nations (U.S., U.K., France, Spain, and Hungary).

Middle right: The OMEGA and OMEGA EP Laser Systems require regular maintenance to ensure high reliability and effectiveness. Shown in the photograph is Laboratory Engineer Michael Scipione examining a brick of flash lamps on the OMEGA EP laser.

Bottom right: LLE designed and installed a novel Target Area System for the Dynamic Compression Sector (DCS) at the Advanced Photon Source located at Argonne National Laboratory near Chicago. The system includes a target chamber, target positioner, and optical train to deliver the DCS laser (also constructed by LLE). Dale Guy (left) and Robert Early (right) are shown working on the system.

This report was prepared as an account of work conducted by the Laboratory for Laser Energetics and sponsored by New York State Energy Research and Development Authority, the University of Rochester, the U.S. Department of Energy, and other agencies. Neither the above named sponsors, nor any of their employees, makes any warranty, expressed or implied, or assumes any legal liability or responsibility for the accuracy, completeness, or usefulness of any information, apparatus, product, or process disclosed, or represents that its use would not infringe privately owned rights. Reference herein to any specific commercial product, process, or service by trade name, mark, manufacturer, or otherwise, does not necessarily constitute or imply its endorsement, recommendation, or favoring by the United States Government or any agency thereof or any other sponsor. Results reported in the LLE Review should not be taken as necessarily final results as they represent active research. The views and opinions of authors expressed herein do not necessarily state or reflect those of any of the above sponsoring entities.

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LLE 2017 Annual Report

October 2016 – September 2017

Inertial Fusion Program and National Laser Users’ Facility Program
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Executive Summary

The federal fiscal year ending September 2017 (FY17) concluded the first 54 months of the fourth five-year renewal of Cooperative Agreement DE-NA0001944 with the U.S. Department of Energy (DOE). This annual report summarizes work carried out under the Cooperative Agreement at the Laboratory for Laser Energetics (LLE) during the past fiscal year including work on the Inertial Confinement Fusion (ICF) Campaign; laser, optical materials, and advanced technology development; operation of the Omega Laser Facility for the ICF and High-Energy-Density (HED) Campaigns, the National Laser Users’ Facility (NLUF), the Laboratory Basic Science (LBS) Program, and other external users; and programs focusing on the education of high school, undergraduate, and graduate students during the year.

Inertial Confinement Fusion Research

One of LLE’s principal missions is to conduct research in ICF with particular emphasis on supporting the goal of achieving ignition at the National Ignition Facility (NIF). This program uses the Omega Laser and NIF facilities and the full experimental, theoretical, and engineering resources of the Laboratory. During FY17, 2138 target shots were taken at the Omega Laser Facility (comprised of the 60-beam OMEGA UV laser and the four-beam, high-energy petawatt OMEGA EP laser). Of the facility’s 2138 target shots, 72% were designated for ICF and HED campaigns. LLE is the lead laboratory worldwide for the laser-direct-drive approach with research focused on cryogenic implosions on the 60-beam OMEGA laser. LLE is responsible for a number of critical elements within the Integrated Experimental Teams that support the demonstration of indirect-drive ignition on the NIF and is the lead laboratory for the validation of direct-drive ignition. LLE has also developed, tested, and constructed a number of diagnostics currently being used at both the Omega Laser Facility and on the NIF. During this past year, progress in the Inertial Fusion Research Program continued in three principal areas: ICF experiments and experiments in support of ICF; theoretical analysis and design efforts aimed at improving direct-drive–ignition capsule designs and advanced ignition concepts; and development of diagnostics for experiments on the NIF, OMEGA, and OMEGA EP facilities.

1. Inertial Confinement Fusion Experimental Highlights in FY17

Measurements of the deuterium–tritium (DT) to deuterium–deuterium (DD) fusion neutron yield ratio were used (p. 23) to evaluate species separation in OMEGA cryogenic target implosion experiments. No species separation was found in these experiments.

Beginning on p. 29 we report on OMEGA spherically converging shock experiments, comparing CH, Be, C, and SiO₂ ablators. CH gives 2 to 3× more hot electrons than other ablators and, as expected, a higher effective ablation pressure.

OMEGA cryogenic target experiments are presented (p. 36) in which x-ray self-emission diagnostics were deployed using the short-pulse OMEGA EP laser to drive a Si Heα x-ray backlighter. The evolution of nonuniform ablators, perturbations caused by mounting stalks, and carbon mix into the DT are observed for capsule adiabats lower than 4.

An overview of the development of a laser-driven MagLIF (magnetized liner inertial fusion) platform on the OMEGA laser is provided (p. 10). MagLIF was developed at the Z Pulsed Power Facility at Sandia National Laboratories and is a key target concept in the U.S. ICF Program. Laser-driven MagLIF on OMEGA is being developed to provide the data on energy scaling and to allow for more shots with better diagnostic access than Z, including proton radiography, to more directly measure embedded B fields, facilitating basic physics studies.

The mitigation of cross-beam energy transfer (CBET) in OMEGA direct-drive implosions by wavelength detuning the three separate legs of the system using a 3-D model is being investigated (p. 61). CBET redistributes power from the ongoing central portion of the outgoing edge of OMEGA beams, significantly increasing the root-mean-squared absorption nonuniformity and reducing the total absorbed power and ablation pressure. A wavelength shift of ±10 Å on two legs is found to be optimal for absorption and close to optimal for absorption uniformity.
The 3-D hydrodynamic code HYDRA and the neutron tracking code IRIS3D were used to interpret neutron emission measurements (p. 100). It is shown that background subtraction is important for inferring areal density from backscattered neutrons, but less important for forward-scattered neutrons, and is important for inferring ion temperature from D–D neutrons, but is insignificant when inferring ion temperature from D–T neutrons at the areal densities typical of OMEGA implosions. Asymmetries resulting in fluid flow in the core are shown to influence the absolute inferred ion temperatures from both reactions. Relative inferred temperatures reflect the underlying asymmetry of the implosion and residual kinetic energy at stagnation.

Planar laser–plasma interaction (LPI) experiments on the NIF have for the first time allowed access to the regimes of electron density scale length (500 to 700 nm), electron temperature (3 to 5 keV), and laser intensity (\(6 \times 10^{14}\) W/cm\(^2\) to \(16 \times 10^{14}\) W/cm\(^2\)) that are relevant to direct-drive ICF ignition (see p. 140). Scattered-light data on the NIF show that the near-quarter-critical LPI physics appears to be dominated by stimulated Raman scattering rather than by two-plasmon decay. These results have significant implications for the mitigation of LPI hot-electron preheat in direct-drive–ignition designs.

An effective method for determining the offsets of the cryogenic implosion cores generated in OMEGA’s ICF experiments was demonstrated (p. 146). This method utilizes images taken by the gated microscopic x-ray imaging diagnostic module. The cryogenic shot images are cross correlated onto images of their respective pulse-shape setup shot images. A true offset is then determined to be the average of the offsets calculated in the images, with the difference between those offsets being taken as the error. Initial offset results using this method indicate that the determined core offsets follow the core offsets at \(t_0\).

Tomographic x-ray images of noncryogenic targets imploded in the direct-drive configuration on the 60-beam OMEGA laser were used to measure 3-D drive asymmetry in target modes \(\lambda = 1, 2, \) and \(3\) at a convergence ratio of \(\sim 3\) (p. 152). Laser configurations were varied linearly with the corresponding modes. This made it possible to use the linear evolutions to determine the residual target mode amplitudes that remain when the laser beam energies are balanced and the laser mode-amplitude compensations are obtained. The analysis provides a means to determine the residual target modes (and the laser modes that compensate them) that agree with 3-D simulations, which predicts significant enhancements in fusion performance.

The first observation of CBET mitigation for direct-drive ICF implosions using wavelength detuning on the NIF has been reported. This article (p. 169) discusses the CBET results from two-beam energy exchange via seeded stimulated Brillouin scattering, which, as mentioned, detrimentally reduces ablation pressure and implosion velocity in direct-drive ICF. Direct-drive implosions on the NIF were conducted to reduce CBET by detuning the laser-source wavelengths (±2.3 Å UV) of the interacting beams over the equatorial region of the target. For the first time, wavelength detuning was shown experimentally to increase the equatorial region velocity by 16% and to alter the in-flight shell morphology. These experimental observations are consistent with design predictions of radiation–hydrodynamic simulations that indicate a 10% increase in the average ablation pressure.

2. Theoretical Design and Analysis

A significant portion of the cryogenic capsule implosion campaign in FY17 was dedicated to the 1-D implosion campaign used to develop a predictive model for direct-drive cryogenic implosions. The model includes a statistical mapping of the experimental results onto a simulation database of high-adiabat implosions (\(\alpha \sim 5\) having convergence ratios of \(\sim 11\) to 13). The nonlinear regression formula from the mapping is used to bridge the gap between experiments and simulations. Using this approach, neutron yields in this cryogenic capsule campaign have exceeded \(10^{14}\) neutrons and areal densities have exceeded 100 mg/cm\(^2\).

The effects of low-mode asymmetries on OMEGA direct-drive implosions using the 3-D Eulerian hydrodynamic code \(\text{ASTER}\) are analyzed beginning on p. 1. Beam-power balance, beam mispointing, beam mistiming, target offset, and variation in target-layer thickness are considered, using values determined from experimental measurements. \(\text{ASTER}\) indicates that implosion performance is mainly affected by target offset (\(\sim 10\) to 20 μm), beam-power balance (\(\sigma_{\text{rms}} \sim 10\%\)), and variation in target-layer thickness (\(\sim 5\%\)).

The impact of beam speckle and polarization smoothing on CBET is compared using the 3-D wave-based laser–plasma interaction code \(\text{LPSE}\) and ray-based models (p. 128). The results indicate that ray-based models underpredict CBET when the assumption of spatially averaged longitudinal incoherence across the CBET interaction region is violated. A model for CBET between linearly polarized speckled beams that uses ray tracing to solve for the real speckle pattern of the unperturbed laser beams within the eikonal approxima-
Two novel target designs for using direct laser ablation (direct drive) on the NIF to assemble and ignite cryogenic fuel using the existing indirect-drive beam configuration are presented (p. 176). These two designs are the first ignition-relevant “polar” direct-drive target designs to include the physical effects of CBET between laser beams and nonlocal electron heat transport. A wavelength-detuning strategy is used to reduce scattered-light losses caused by CBET, allowing for ignition-relevant implosion velocities. The designs include: (1) a moderate-adiabat alpha-burning design with a D–T neutron fusion yield of $1.2 \times 10^{17}$ and (2) a lower-adiabat ignition design with a gain of 27. Both designs have low in-flight aspect ratios, which imply improved hydrodynamic stability levels of perturbation growth during the implosion.

3. Diagnostics

During FY17, LLE installed and operated the single-line-of-sight time-resolved x-ray imager (TRXI) on OMEGA. TRXI is the result of a collaboration among General Atomics, Lawrence Livermore, LLE, and Sandia National Laboratories. The instrument was used on cryogenic capsule experiments and provided high-resolution (~10-μm spatial, ~40-ps temporal) imaging of highly compressed DT implosion cores. This enhanced imaging system will deliver significantly higher resolution than previously available, advancing stockpile stewardship diagnostics for ICF facilities throughout the Nuclear Security Enterprise.

Our LLE–LLNL collaborative effort demonstrated high-efficiency fifth-harmonic conversion of IR laser light. Initial tests were done using heated cesium lithium borate (CLBO) crystals and achieved 30% conversion efficiency. More recent tests use ammonium dihydrogen phosphate (ADP) coded to $-70^\circ$C in a specially designed two-chamber cryostat. The fifth-harmonic beams will eventually be used to probe dense plasma using Thomson scattering.

A new technique to extract electron density profiles from angular filter refractometry (AFR) (p. 48) that makes use of a simulated annealing algorithm was developed. A seven-parameter function was chosen for the electron density and used to generate an AFR image that is compared to the measurement using a $\chi^2$ test. The algorithm was applied to measurements of plasma expansion from a planar target and produced a fit with a statistical uncertainty of no more than 10% in the region of interest ($10^{20}$ to $10^{21}$ cm$^{-3}$).

A 16-image Kirpatrick–Baez-type x-ray microscope coupled to a high-speed framing camera is described beginning on p. 79. A temporal resolution of ~30 ps and a spatial resolution of ~6 μm were achieved with this diagnostic. The new diagnostic makes it possible to accurately determine the cryogenic implosion core emission size and shape at the peak of the stagnation. This system was used to determine cryogenic capsule core pressures in excess of 50 Gbar.

Beginning on p. 184, we report on the linearity of the photo-stimulated luminescence process to make repeated image-plate scanning a viable technique to extract a more-dynamic range. To obtain a response estimate for second and subsequent scans with a BAS-MS image plate and the Typhoon FLA 7000 scanner, a new model for the readout fading of the image plate is introduced; it relates the depth distribution of activated photostimulated luminescence centers within the image plate to the recorded signal. Model parameters are estimated from an image-plate scan series for the hard x-ray image-plate diagnostic over a collection of experiments, providing x-ray energy spectra whose approximate shape is a double exponential.

4. Target Technology

A model for tritium interaction with metals was developed (p. 87) that can predict the temporal evolution of tritium concentration profiles during exposure to tritium gas, during storage and during successive decontamination efforts, and could be used to develop surfaces that are less prone to absorbing tritium. Operation of a tritium facility requires an understanding of the migration of tritium in metals.

Thermal contraction anomalies seen in glow-discharge polymer (GDP) capsules with a layer of an equimolar mixture of DT on their interior, compared to GDP with only deuterium and polystyrene capsules permeated with only DT, are discussed beginning on p. 159. Thermal contraction of the GDP-mixture capsules from cooling do not exhibit expected contraction and retain their room-temperature diameter after cooling. It is speculated that the highly cross-linked GDP shell is under compressive stress after fabrication and experiences bond breakage when exposed to high-density DT during permeation and some of this compressive stress is relieved during bond cleavage, causing the capsule’s wall to swell, which counteracts contraction during cooling.
High-Energy-Density Science

During FY17, several high-energy-density science campaigns were conducted at the Omega Laser Facility including:

- measurements of the sound velocity and Grüneisen parameter in CH shocked to 800 GPa,
- equation-of-state measurements of CO\textsubscript{2} precompressed to 1.2 GPa and shock compressed to 980 GPa,
- the solid hP4 phase of Na was observed at \(~320\) GPa,
- Hugoniot measurements were conducted of Si shock compressed to 21 Mbar,
- a new high-pressure solid phase of dynamically compressed Al was observed, and
- in collaboration with Harvard University, the optical reflectance of dense hydrogen was measured as a function of energy in the 1.4- to 1.7-Mbar region and up to 2500 K. The data are consistent with the metallic hydrogen being a free-electron partially ionized plasma.

Picosecond time-resolved measurements of the shift of the 1\textit{s}2\textit{p}–1\textit{s}2 line in He-like Al as a function of electron density are reported (p. 73). Temperature and density are inferred from the Al He\textsubscript{a} complex using a nonlocal-thermodynamic equilibrium model. The measurements are broadly consistent with an analytic line-shift model based on calculations of a self-consistent field ion-sphere model.

In the article beginning on p. 13, we report on the evaluation of the equation of state of silicon using density-functional-theory dynamics simulations for densities from 0.001 to 500 gm/cm\textsuperscript{3} and temperatures from 2000 to 108 K. This first-principles equation of state (FPEOS) is compared to SESAME 3810. The Hugoniot from FPEOS is \(~20\)% softer than that from SESAME 3810 below \(10^4\) to \(10^5\) K, depending on density, and lower at higher temperatures. In LILAC simulations of a silicon-shell implosion, FPEOS gives \(~30\)%–higher areal density and \(~70\)%–higher neutron yield than SESAME 3810 because of the larger compressibility of silicon in FPEOS.

Lasers, Optical Materials, and Advanced Technology

LLE developed, constructed, installed, and activated a 100-J UV laser and Target Area System for the Dynamic Compression Sector at the Advanced Photon Source located at the Argonne National Laboratory near Chicago. This new research facility is operated by Washington State University (WSU) under sponsorship from the National Nuclear Security Administration (NNSA). LLE partnered with Logos Technologies to develop and build the high-energy laser, which is suitable for a broad range of applications.

A novel approach is introduced (p. 115) for controlling laser–plasma interactions that remove the need for long-focal-length systems or guiding structures to maintain high intensities over long distances and decouples the velocity of the focal spot from the group velocity of the light. This advanced focusing scheme, called a “flying focus,” enables a small-diameter laser focus to propagate nearly 100\times its Rayleigh length. By providing unprecedented spatiotemporal control over the laser focal volume, it allows the laser focus to co- or counter-propagate along its axis at any velocity.

A new laser-amplifier scheme has been proposed (p. 122) that will utilize stimulated Raman scattering in plasma in conjunction with a flying focus—a chromatic focusing system combined with a chirped pump beam that provides spatiotemporal control over the pump’s focal spot. Simulations show that this enables optimization of the plasma temperature and mitigates many of the issues that are known to have impacted previous Raman amplification experiments, in particular the growth of precursors.

A time-to-frequency converter was constructed using an electro-optic phase modulator as a time lens, allowing the pulse shape in time to be transferred to the frequency domain (p. 192). The device was used to record the temporal shape of infrared pulses at a wavelength of 1053 nm (width about 7 ps) and to compare these measurements to those made by using both a streak camera and an autocorrelator. Numerical simulations were used to establish that the time-lens–based system can accurately measure the shape of infrared pulses between 3 and 12 ps. The numerical model was also used to determine how such a system can be modified to measure pulses whose width lies in the range of 1 to 30 ps—a range of interest for the OMEGA EP laser.

Omega Laser Facility Users Group

The Ninth Omega Laser Facility Users Group (OLUG) Workshop was held at LLE on 26–28 April 2017. It was attended by 110 researchers, including scientists, postdoctoral
fellows, and students. The attendees represented institutions from five nations, including the United States, United Kingdom, France, Spain, and Hungary. The workshop included the presentation of invited talks, a talk on the NNSA perspective by Dr. Njema Frazier of ICF/NNSA, facility tutorials, poster papers, presentations on research at the national laboratories, and panel discussions. A summary of the OLUG Workshop is presented in an article starting on p. 198.

Education

As the only major university participant in the National ICF Program, education continues to be an important mission for LLE. The Laboratory’s education programs cover the range from high school (p. 204) to graduate education.

1. High School Program

During the summer of 2017, 11 students from Rochester-area high schools participated in the Laboratory for Laser Energetics’ Summer High School Research Program. The goal of this program is to excite a group of high school students about careers in the areas of science and technology by exposing them to research in a state-of-the-art environment. Too often, students are exposed to “research” only through classroom laboratories, which have prescribed procedures and predictable results. In LLE’s summer program, the students experience many of the trials, tribulations, and rewards of scientific research. By participating in research in a real environment, with an LLE advisor, the students often become more excited about careers in science and technology. In addition, LLE gains from the contributions of the many highly talented students who are attracted to the program. Three hundred and sixty-four high school students have now participated in the program since it began in 1989. Two of this year’s program participants, Nikhil Bose and Yujia Yang, were named Science Talent Search “Scholars” in the prestigious Regeneron Science Talent search for the research projects they carried out at LLE. Bose developed a simulation model to explore a novel method of improving the performance of LLE’s OMEGA EP laser and Yang carried out hydrodynamic simulations of a new fusion concept for the NIF. This year’s students were selected from approximately 60 applicants to the program.

2. Undergraduate Student Program

Forty undergraduate students participated in work or research projects at LLE this past year. Student projects included operational maintenance of the Omega Laser Facility; work in laser development, materials, and optical thin-film coating laboratories; computer programming; image processing; and diagnostics development. This is a unique opportunity for students, many of whom go on to pursue a higher degree in the area in which they gained experience at LLE.

3. Graduate Student Program

Graduate students are using the Omega Laser Facility as well as other LLE facilities for fusion and HED physics research and technology development activities (see Table I). These students are making significant contributions to LLE’s research program. Twenty-six faculty members from five University of Rochester academic departments collaborate with LLE scientists and engineers. In FY17, 72 graduate students were involved in research projects at LLE, and LLE directly sponsored 47 students pursuing Ph.D. degrees via the NNSA-supported Frank Horton Fellowship Program in Laser Energetics. Their research includes theoretical and experimental plasma physics, HED physics, x-ray and atomic physics, nuclear fusion, ultrafast optoelectronics, high-power laser development and applications, nonlinear optics, optical materials and optical fabrication technology, and target fabrication. Two of the 2017 LLE Ph.D. graduates are pursuing careers in NNSA programs. Dr. Michelle Gregor is now a post-doctoral fellow at Lawrence Livermore National Laboratory (LLNL) and Dr. Amanda Davis is an NNSA Graduate Fellow in the NNSA Defense Programs.

LLE also directly funds research programs within the MIT Plasma Science and Fusion Center, the State University of New York (SUNY) at Geneseo, and the University of Wisconsin. These programs involve a total of approximately 6 graduate students, 25 to 30 undergraduate students, and 10 faculty members. Over 340 graduate students have now conducted their graduate research work at LLE since its graduate research program began.

In addition, 170 graduate students and post-graduate fellows from other universities have conducted research at the Omega Laser Facility as part of the NLUF program. Thirty-three graduate students (Table 152.VII, p. 211) and approximately 30 undergraduate students were involved in NLUF research programs in FY17.

FY17 Omega Laser Facility Operations

During FY17, the Omega Laser Facility conducted 1353 target shots on OMEGA and 785 target shots on OMEGA EP for a total of 2138 target shots (see Tables 152.IV and 152.V, p. 206). OMEGA averaged 10.7 target shots per operating day with Availability and Experimental Effectiveness averages for FY17.
# Executive Summary

## FY17 Annual Report

Table I: University of Rochester graduate students conducting research at LLE in FY17.

<table>
<thead>
<tr>
<th>Student Name</th>
<th>Dept.</th>
<th>Faculty Advisor</th>
<th>LLE Advisor</th>
<th>Funding</th>
<th>Research Area</th>
<th>Status</th>
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</thead>
<tbody>
<tr>
<td>M. B. Adams</td>
<td>PAS</td>
<td>P.-A. Gourdain</td>
<td>Horton</td>
<td>Numerical studies of instabilities in plasma jets generated by the laser ablation of DT ice</td>
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<td>Y. E. Akbas</td>
<td>PAS</td>
<td>R. Sobolewski</td>
<td>Other</td>
<td>Intrinsic nanostructures: asymmetric nanochannel devices</td>
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<td>D. H. Barnak</td>
<td>PAS</td>
<td>R. Betti</td>
<td>Horton</td>
<td>Applications of magnetic fields in high-energy-density physics</td>
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<td>D. Bassler</td>
<td>CHEM</td>
<td>W. U. Schröder</td>
<td>Other</td>
<td>Radiochemistry of tritium transport</td>
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<td>K. Bauer</td>
<td>OPT</td>
<td>R. Brown and M. A. Alonso</td>
<td>J. D. Zuegel</td>
<td>351-mm characterization of scatter from distributed polarization rotator concepts for polar direct drive at the National Ignition Facility (NIF)</td>
<td></td>
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<tr>
<td>M. Berkgoetter</td>
<td>OPT</td>
<td>B. E. Kruschwitz</td>
<td>Other</td>
<td>Implosion dynamics of direct-drive targets in the presence of magnetic fields</td>
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<tr>
<td>A. Bose</td>
<td>PAS</td>
<td>R. Betti</td>
<td>Horton</td>
<td>Nonlinear effects in multimode fibers</td>
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<td>S. Buch</td>
<td>OPT</td>
<td>G. P. Agrawal and W. R. Donaldson</td>
<td>Other</td>
<td>Novel 1200-nm high-power laser beams for laser-plasma amplification</td>
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<td>S. Bucht</td>
<td>PAS</td>
<td>D. H. Froula</td>
<td>Horton</td>
<td>Terahertz detection and imaging with ultra-broadband high-energy lasers</td>
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<tr>
<td>L. E. Bukowski</td>
<td>OPT</td>
<td>W. H. Knox</td>
<td>J. D. Zuegel</td>
<td>Active vibration stabilization of NIF-scale direct-drive cryogenic targets</td>
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<tr>
<td>E. Burnham-Fay</td>
<td>ME</td>
<td>J. D. Ellis</td>
<td>D. W. Jacobs-Perkins</td>
<td>Defending Feb. 2018</td>
<td></td>
<td></td>
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<tr>
<td>J. Cady</td>
<td>ECE</td>
<td>R. Sobolewski</td>
<td>Other</td>
<td>Ultrafast semiconductor devices</td>
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<td>G. Chen</td>
<td>MSC</td>
<td>R. Sobolewski</td>
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<td>Time-domain THz spectroscopy</td>
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<td>Z. Chen</td>
<td>PAS</td>
<td>A. Frank</td>
<td>Horton</td>
<td>Evolution of binary stars</td>
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<td>D. A. Chin</td>
<td>PAS</td>
<td>G. W. Collins and J. R. Rygg</td>
<td>T. R. Boehly</td>
<td>Using extended x-ray absorption fine structure to characterize highly compressed matter</td>
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<td>B. P. Chock</td>
<td>CE</td>
<td>D. R. Harding</td>
<td>Horton</td>
<td>Enhancing electric-field–enabled droplet motion for target production</td>
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<td>A. R. Christopherson</td>
<td>ME</td>
<td>R. Betti</td>
<td>Horton</td>
<td>Measuring and understanding hydrodynamic performance of cryogenic implosions on the OMEGA laser</td>
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<td>L. Crandall</td>
<td>PAS</td>
<td>J. R. Rygg and G. W. Collins</td>
<td>T. R. Boehly</td>
<td>Equation of state of planetary fluids</td>
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<td>C. Danly</td>
<td>ME</td>
<td>S. P. Regan</td>
<td>Other</td>
<td>Implosions, stagnation, and neutron-imaging diagnostics</td>
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<td>A. Davies</td>
<td>PAS</td>
<td>D. H. Froula</td>
<td>D. Haberberger</td>
<td>Ultrafast measurements of electron temperature in a laser-plasma amplifier</td>
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<td>A. K. Davis</td>
<td>PAS</td>
<td>D. H. Froula</td>
<td>D. T. Michel</td>
<td>Three-dimensional measurements of direct-drive implosions with low-mode nonuniformities</td>
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<td>Y. Ding</td>
<td>ME</td>
<td>R. Betti</td>
<td>S. X. Hu</td>
<td>First-principles investigations on transport and optical properties of high-energy-density plasmas</td>
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*FY17 Annual Report*
Table I: University of Rochester graduate students conducting research at LLE in FY17 (continued).

<table>
<thead>
<tr>
<th>Student Name</th>
<th>Dept.</th>
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<th>LLE Advisor</th>
<th>Funding</th>
<th>Research Area</th>
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<tr>
<td>T. Eckert</td>
<td>PAS</td>
<td>P.-A. Gourdain</td>
<td>C. Forrest</td>
<td>Horton</td>
<td>Experimental nuclear physics</td>
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<td>M. Evans</td>
<td>PAS</td>
<td>P.-A. Gourdain</td>
<td>C. Forrest</td>
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<td>Experimental study of instabilities in plasma jets generated by laser ablation of DT ice</td>
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<td>C. Fagan</td>
<td>CHEM</td>
<td>W. U. Schröder</td>
<td>Other</td>
<td></td>
<td>Surface effects in tritium absorption</td>
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<td>P. Franke</td>
<td>PAS</td>
<td>D. H. Froula</td>
<td>Horton</td>
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<td>Control of ionization wave propagation with a flying focus</td>
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<td>J. M. Garcia- Figuero</td>
<td>CE</td>
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<td>Manufacture of low-atomic-number nonpolymeric materials</td>
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<td>M. Gates</td>
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<td>R. Sobolewski</td>
<td>Other</td>
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<td>M. K. Ginnane</td>
<td>ME</td>
<td>J. R. Rygg and G. W. Collins</td>
<td>T. R. Boehly</td>
<td>Other</td>
<td>Novel diagnostics to study highly compressed matter</td>
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<td>X. Gong</td>
<td>ME</td>
<td>G. W. Collins and J. R. Rygg</td>
<td>J. R. Rygg</td>
<td>Horton</td>
<td>Structural and optical changes in ramp-compressed alkalais</td>
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<td>V. Gopalswamy</td>
<td>ME</td>
<td>R. Betti</td>
<td>Other</td>
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<td>Hydrodynamics in 1-D physics campaigns</td>
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<td>M. C. Gregor</td>
<td>PAS</td>
<td>G. W. Collins</td>
<td>T. R. Boehly</td>
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<td>Properties of matter that are shocked and compressed to high-energy densities using high-power lasers</td>
<td>Graduated (LLNL Fellow)</td>
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<td>A. Hansen</td>
<td>PAS</td>
<td>D. H. Froula</td>
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<td>Nonlinear cross-beam energy transfer physics</td>
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<td>R. J. Henchen</td>
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<td>Hydrodynamic gradients in underdense plasmas</td>
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<td>B. J. Henderson</td>
<td>PAS</td>
<td>G. W. Collins and J. R. Rygg</td>
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<td>Horton</td>
<td>Optical properties of compressed matter deduced from reflectivity measurements</td>
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<td>J. Hinz</td>
<td>PAS</td>
<td>S. X. Hu</td>
<td>Other</td>
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<td>Using machines to learn to improve the free-energy functionals</td>
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<td>M. Huff</td>
<td>PAS</td>
<td>J. R. Rygg and G. W. Collins</td>
<td>T. R. Boehly</td>
<td>Other</td>
<td>Measurements of sound speed in shock-compressed materials</td>
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<td>G. Jenkins</td>
<td>OPT</td>
<td>J. Bromage</td>
<td>Horton</td>
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<td>Broadband seed generation and amplification at high-average powers</td>
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<td>B. Lam</td>
<td>OPT</td>
<td>J. Shojaie</td>
<td>Other</td>
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<td>Laser–plasma interactions (LPI) in magnetized plasmas</td>
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<td>L. Leal</td>
<td>PAS</td>
<td>R. Betti and A. V. Maximov</td>
<td>Other</td>
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<td>Plasma physics</td>
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<td>A. Lees</td>
<td>ME</td>
<td>H. Aluie</td>
<td>Horton</td>
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<td>High-performance code development</td>
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<td>S. Li</td>
<td>PAS</td>
<td>A. Frank</td>
<td>Other</td>
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<td>Heterogeneous flow in an interstellar medium</td>
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<td>W. Liu</td>
<td>ME</td>
<td>C. Ren</td>
<td>C. Ren</td>
<td>Horton</td>
<td>Magnetic fields and their effects on LPI</td>
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<td>N. Luciani</td>
<td>ME</td>
<td>R. Betti</td>
<td>Other</td>
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<td>Implosion physics</td>
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<td>O. M. Mannion</td>
<td>PAS</td>
<td>S. Y. BenZvi and C. J. Forrest and J. P. Knauer</td>
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<td>Moment analysis of neutron spectra</td>
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<td>A. L. Milder</td>
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<td>D. H. Froula</td>
<td>Horton</td>
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<td>Measurement of electron distribution function using collective Thomson scattering</td>
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<td>S. C. Miller</td>
<td>ME</td>
<td>V. N. Goncharov</td>
<td>P. B. Radha</td>
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<td>Hydrodynamics of ICF implosions</td>
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<td>Z. L. Mohamed</td>
<td>PAS</td>
<td>D. H. Froula</td>
<td>J. P. Knauer</td>
<td>Horton</td>
<td>Gamma emission from fusion reactions</td>
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Table I: University of Rochester graduate students conducting research at LLE in FY17 (continued).

<table>
<thead>
<tr>
<th>Student Name</th>
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<th>LLE Advisor</th>
<th>Funding</th>
<th>Research Area</th>
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<tr>
<td>D. Patel</td>
<td>ME</td>
<td>R. Betti</td>
<td>V. N. Goncharov</td>
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<td>Hybrid direct–indirect drive for ICF</td>
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<td>R. Paul</td>
<td>ME</td>
<td>S. X. Hu</td>
<td></td>
<td>Horton</td>
<td>High-pressure phase diagram of material</td>
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<td>D. N. Polsin</td>
<td>PAS</td>
<td>G. W. Collins</td>
<td>T. R. Boehly</td>
<td>Horton</td>
<td>Observation of a new high-pressure solid phase in dynamically compressed aluminum</td>
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<td>D. Saulnier</td>
<td>CE</td>
<td>K. L. Marshall and M. Anthamatten</td>
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<td>Horton</td>
<td>Liquid crystal chiroptical polarization rotators for the near-UV region</td>
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<td>E. M. Schiesser</td>
<td>OPT</td>
<td>J. Rolland</td>
<td>S.-W. Bahk</td>
<td>Horton</td>
<td>Applying freeform optics to scalable, compact beamlines</td>
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<td>A. Schwemmlein</td>
<td>PAS</td>
<td>W. U. Schröder</td>
<td>J. P. Knauer</td>
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<td>Thermonuclear fusion and breakup reactions between light nuclei</td>
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<td>J. R. Serafini</td>
<td>PAS</td>
<td>R. Sobolewski</td>
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<td>Horton</td>
<td>Ultrafast optical and electronic characterization of (Cd,Mg)Te single crystals</td>
<td>Graduated (postdoc RIT)</td>
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<td>A. Shramuk</td>
<td>ECE</td>
<td>R. Sobolewski</td>
<td></td>
<td>Other</td>
<td>Superconducting optoelectronic devices</td>
<td>M.S. graduated (in law school)</td>
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<td>R. Shrestha</td>
<td>ECE</td>
<td>R. Sobolewski</td>
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<td>Other</td>
<td>Terahertz time-domain spectroscopy of carbon nanotubes</td>
<td>M.S. graduated May 2017</td>
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<td>J. Slater</td>
<td>OPT</td>
<td>J. Bromage</td>
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<td>Characterization of DKDP crystals using spatially resolved Raman spectroscopy</td>
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<td>Z. Sprowal</td>
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<td>G. W. Collins and J. R. Rygg</td>
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<td>C. R. Stillman</td>
<td>PAS</td>
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<td>P. M. Nilson</td>
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<td>Hot dense matter</td>
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<td>G. Tabak</td>
<td>PAS</td>
<td>G. W. Collins and J. R. Rygg</td>
<td>M. Zaghoo</td>
<td>Horton</td>
<td>Study of precompressed materials using shock compression</td>
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<td>N. D. Viza</td>
<td>CE</td>
<td>D. R. Harding</td>
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<td>Horton</td>
<td>Integrated “lab-on-chip” microfluidic device for manufacturing foam targets</td>
<td>Defending Jan./Feb. 2018</td>
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<td>M. Wang</td>
<td>CE</td>
<td>D. R. Harding</td>
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<td>Horton</td>
<td>Use of two-photon polymerization to “write” millimeter-size structures with micron resolution</td>
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<td>I. West-Abdallah</td>
<td>PAS</td>
<td>G. W. Collins and J. R. Rygg</td>
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<td>K. M. Woo</td>
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<td>Horton</td>
<td>Three-dimensional ablative Rayleigh–Taylor instability</td>
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<td>E. Wright</td>
<td>PAS</td>
<td>D. H. Froula</td>
<td>W. Theobald</td>
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<td>Direct measurements of wave breaking using Thomson scattering</td>
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</table>
Table I: University of Rochester graduate students conducting research at LLE in FY17 (continued).

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<td>J.-C. Yang</td>
<td>CE</td>
<td>M. Anthamatten</td>
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<td>Horton</td>
<td>Crystallization in shape-memory polymer networks</td>
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<td>D. Zhao</td>
<td>ME</td>
<td>H. Aluie</td>
<td>Horton</td>
<td>Horton</td>
<td>Analyzing multiscale physics of multimode and turbulent Rayleigh–Taylor hydrodynamics</td>
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<tr>
<td>Y. Zhao</td>
<td>MS</td>
<td>W. R. Donaldson</td>
<td>Horton</td>
<td>Horton</td>
<td>Fabrication and testing of AlGaN photodiodes</td>
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CE: Department of Chemical Engineering
CHEM: Department of Chemistry
ECE: Department of Electrical and Computer Engineering
ME: Department of Mechanical Engineering
MSC: Department of Materials Science
OPT: The Institute of Optics
PAS: Department of Physics and Astronomy

of 95.7% and 94.4%, respectively. OMEGA EP was operated extensively in FY17 for a variety of internal and external users. A total of 773 target shots were taken in the OMEGA EP target chamber and 12 joint target shots were taken in the OMEGA target chamber. OMEGA EP averaged 8.7 target shots per operating day with Availability and Experimental Effectiveness averages for FY17 of 95.8% and 96.6%, respectively. Per the guidance provided by DOE/NNSA, the facility provided target shots for the ICF, HED, NLUF, and LBS programs. The facility also provided a small number of shots for Commissariat à l’énergie atomique et aux energies (CEA), Centre Lasers Intenses et Applications (CEMIA), and ARPA-E programs (see Figs. 1 and 2). Nearly 70% of the target shots in FY17 were taken for the ICF and HED programs.

Details of this work are contained in an article beginning on p. 206. Highlights of the Omega Laser Facility activities in FY17 included the following:

- The 100-Gbar Campaign worked to achieve higher implosion pressures through improved laser power balance and a cryogenic fill-tube target system.

- To mitigate CBET, a tunable wavelength UV beam will be produced by one of the OMEGA EP beams. This beam will be injected into OMEGA to study CBET mitigation via wavelength detuning (up to 30 Å).

- Various improvements to the laser systems were implemented in FY17 including a new ultrafast temporal diagnostic and a time-multiplexed pulse-shaping system to OMEGA EP, and an improved smoothing by spectral dispersion spectrometer on OMEGA.
• Target diagnostic improvements included a high-resolution spectrometer that was implemented on OMEGA EP; a single line-of-sight, time-resolved x-ray imager was developed in collaboration with General Atomics, LLNL, Sandia National Laboratories (SNL), and Kentech Instruments and deployed on OMEGA; an improved signal-to-noise ratio neutron time-of-flight diagnostic was deployed on port H10; and a powder x-ray diffraction diagnostic is being upgraded to acquire time-resolved images.

National Laser Users’ Facility and External Users Programs

The Fundamental Science Campaigns accounted for nearly 24% of the Omega Laser Facility target shots taken in FY17. Nearly 62% of these shots were taken for experiments for the NLUF Program, and the remaining shots were allotted to the LBS Program, comprising peer-reviewed fundamental science experiments conducted by the national laboratories and by LLE.

The Omega Laser Facility was also used for several campaigns by teams from CEA of France. These programs are conducted at the facility on the basis of special agreements put in place by DOE/NNSA and participating institutions.

The facility users during this year included 13 collaborative teams participating in the NLUF Program; 14 teams led by Los Alamos National Laboratory (LANL), LLNL, and LLE scientists participating in the LBS Program; many collaborative teams from the national laboratories [LANL, LLNL, SNL, and Naval Research Laboratory (NRL)] and LLE conducting ICF experiments; investigators from LLNL, LANL, and LLE conducting experiments for high-energy-density–physics programs; and scientists and engineers from CEA and CELIA. Nearly 60% of the facility target shots were provided to external users.

1. FY17 NLUF Program

During the first quarter of FY17, the Inertial Fusion Office of DOE/NNSA completed a solicitation, review, and selection process for NLUF experiments to be conducted at the Omega Laser Facility during calendar years 2017 and 2018. Twenty-eight proposals were submitted in response to the call for proposals and the shot requests totaled 60.5 shot days at the Omega Laser Facility. The proposals were peer reviewed by an independent review committee and ICF/NNSA selected 13 proposals for funding and shot allocation for the period calendar years 2017 and 2018. These NLUF projects (see Table 152.VI, p. 210) were allotted Omega Laser Facility shot time and conducted 319 target shots at the facility. The FY17 NLUF experiments are summarized beginning on p. 209.

2. FY17 Laboratory Basic Science Studies

Sixteen LBS projects previously approved for FY17 target shots were allotted Omega Facility shot time and conducted a total of 199 target shots at the Omega Facility in FY17 (see Table 152.VII, p. 230). The FY17 LBS experiments are summarized beginning on p. 230.

During FY17, LLE issued a solicitation for LBS proposals to be conducted in FY18. A total of 28 proposals were submitted. An independent review committee reviewed and ranked the proposals; on the basis of these scores, 16 proposals were allocated 21 shot days at the Omega Laser Facility in FY18. Table 152.IX, p. 231 lists the approved FY18 LBS proposals.

3. FY17 LLNL Experimental Programs

In FY17 LLNL’s HED Physics and Indirect-Drive Inertial Confinement Fusion (ICF-ID) Programs conducted numerous campaigns on the OMEGA and OMEGA EP Laser Systems. Overall, these LLNL programs led 413 target shots in FY17, with 282 shots using only OMEGA and 131 shots using only OMEGA EP. Approximately 27% of the total number of shots (78 OMEGA shots and 35 OMEGA EP shots) supported the ICF-ID Campaign. The remaining 73% (204 OMEGA shots and 96 OMEGA EP shots) were dedicated to experiments for HED physics. Highlights of the various HED and ICF-ID Campaigns are summarized beginning on p. 242.

In addition to these experiments, LLNL Principal Investigators (PI’s) led a variety of Laboratory Basic Science Campaigns using OMEGA and OMEGA EP, including 85 target shots using only OMEGA and 70 shots using only OMEGA EP.

Overall, LLNL PI’s led a total of 568 shots at LLE in FY17. In addition, LLNL PI’s also supported 30 NLUF shots on OMEGA and 46 NLUF shots on OMEGA EP, in collaboration with the academic community.

4. FY17 LANL Experimental Campaigns

In FY17, LANL carried out 17 shot days comprising 218 target shots on the OMEGA and OMEGA EP Laser Systems in the areas of HED science and ICF. In HED, the LANL focus was on areas of radiation flow, hydrodynamic turbulent mix and burn, warm-dense-matter equations of state, and coupled Kelvin–Helmholtz/Richtmyer–Meshkov instability growth. For ICF, the campaigns focused on the Priority
Research Directions (PRD’s) of implosion phase mix and stagnation and burn, specifically as they pertain to laser direct drive (LDD). Several focused shot days were also dedicated to transport properties in the kinetic regime. In addition, LANL continues to develop advanced diagnostics such as neutron imaging, gamma reaction history, and gas Cherenkov detectors at the Omega Laser Facility.

5. FY17 SNL Experimental Campaigns
   During FY17, SNL conducted a total of four shot days (one on OMEGA and three on OMEGA EP) aimed at characterizing the laser heating of underdense plasmas (D₂, Ar) at parameters that are relevant to the MagLIF ICF scheme. In total these accounted for 44 target shots at the facility. Reports on these campaigns are on p. 280.

6. FY17 NRL Experimental Campaigns
   During FY17, the NRL/LLE collaboration on laser imprint led to three successful shot days on OMEGA EP. The experiments showed that the application of a prepulse that pre-expands and lifts off the coating prior to the arrival of the main laser pulse gives an order of magnitude reduction of laser imprint, as expected on the basis of the original experiments on the Nike laser and an understanding of the mechanism of the imprint suppression. Further experiments demonstrated imprint reduction with prepulse times compatible with pulse durations available for implosions on OMEGA. LLE is evaluating the implementation of a suitable prepulse for imprint reduction with NRL.

7. FY17 CEA Experiments
   During FY17 CEA conducted 59 target shot experiments at the Omega Laser Facility (four shot days on OMEGA and one on OMEGA EP). The experiments included studies of rugby hohlraums in preparation for Laser Mégajoule campaigns, measurements of the Hugoniot of LiH, spectroscopy of direct-drive implosions, studies of beam propagation in gas-filled cavities, and laser-plasma interaction experiments. Some of these experiments are described (p. 282).

E. Michael Campbell
Director, Laboratory for Laser Energetics
Introduction
Direct-drive inertial confinement fusion (ICF) experiments conducted at the 30-kJ Omega Laser Facility are used to demonstrate the hydrodynamic equivalence of scaled-down cryogenic target implosions to ignition designs at MJ energies such as those available at the National Ignition Facility. OMEGA implosion experiments demonstrate good agreement between the measured and simulated efficiency of conversion of the laser energy into the kinetic energy of the imploding shell ($\sim 4\%$). The fuel-compression stage of cryogenic implosions significantly underperforms, however, typically showing that the implosion’s hot-spot pressure and deuterium–tritium (D–T) fusion neutron yield do not exceed $\sim 60\%$ of the values predicted in simulations using the one-dimensional (1-D) radiation–hydrodynamics code LILAC. This and other experimental evidence, including asymmetries of x-ray images of implosion shells and hot spots, nonspherical distribution of stagnated fuel shell $\rho R$, and $\sim 100$-km/s directional motions of hot-spot plasma, both inferred from neutron measurements, suggest that short- and long-scale nonuniformities in implosion shells can cause the observed performance degradation.

Short-scale nonuniformities (corresponding to Legendre modes $\ell \gtrsim 30$) can be seeded by laser imprint and small target-surface and structural defects. The effects of Rayleigh–Taylor (RT) growth of these nonuniformities likely dominate over other effects of performance degradation in low-adiabat ($\alpha \lesssim 3$) and high in-flight aspect ratio (IFAR $\gtrsim 25$) implosions. Here, the adiabat $\alpha$ is defined in 1-D simulations as the ratio of the pressure in the imploding DT fuel shell to the corresponding Fermi-degenerated pressure and the IFAR is defined as the ratio of the shell’s radius to its thickness (at a density level of 1 g/cm$^3$) at the moment when the ablation radius equals 2/3 of the initial radius of the inner shell. The short-scale RT-growth effects can be mitigated using mid- to high-adiabat ($\alpha \gtrsim 4$) and/or low-IFAR ($\lesssim 20$) implosions.

Large-scale nonuniformities (with modes $\ell \lesssim 10$) can develop because of laser illumination and structural asymmetries of implosion targets. The asymmetry of illumination is caused by the OMEGA laser’s 60-beam-port configuration in addition to target offset ($\sim 10$ to $20\ \mu m$) and inaccuracy of pointing, power balance, and timing of the beams (with typical $\sigma_{\text{rms}} < 10\ \mu m$, 10%, and 5 ps, respectively). The structural asymmetries include mounting stalks, variations of thickness and shape of plastic (CH or CD) ablator shells in warm and cryogenic targets (with $\sigma_{\text{rms}} < 1\ \mu m$), and variations in thickness of the DT ice layer in cryogenic targets (with $\sigma_{\text{rms}} \sim 1\ \mu m$). Large-scale modes are amplified by the secular and Bell–Plesset growths and by the RT growth during the deceleration and stagnation stages. Variations of $\alpha$ and IFAR have little effect on the growth of these modes.

Investigation of the effects of large-scale asymmetries and the development of strategies to mitigate them are important steps toward improving the performance of OMEGA implosions. To understand these effects, experimental observations of implosion asymmetries are simulated in detail employing the three-dimensional (3-D) radiation–hydrodynamics code ASTER. Results of 3-D simulations are post-processed to be directly compared with observables, which include x-ray images and deuterium–deuterium (DD) and/or DT fusion neutron spectra, among others.

This article describes recent progress in 3-D ASTER simulations of room-temperature and cryogenic OMEGA implosions focusing on large-scale ($\ell \lesssim 10$) target asymmetries as sources of the degradation in implosion performance. Simulations show that mode 1 is typically the most-destructive one in the case of both room-temperature and cryogenic implosions. The presence of this mode results in relatively large residual kinetic energy of implosion shells at maximum compression in comparison with that resulting from other modes ($\geq 2$) of similar amplitude. This large residual kinetic energy causes undercompression of the hot spot and a reduction of neutron yields down to values found in experiments. Mode 1 can be observed as an offset of the core emission in x-ray images with respect to the initial target center and as a directional variation of neutron spectra.
All above-mentioned sources of long-scale nonuniformities (except for that caused by the OMEGA discrete-beam illumination,11 which introduces a dominant mode $\ell = 10$) can contribute to mode-1 perturbations. Mount stalks and target offsets apparently result in such perturbations. Beam mistiming, mispointing, and imbalance, as well as initial target structural asymmetry, can be considered as quasi-random sources and result in perturbations having broad spectra, which peak at the lowest modes from 1 to $\sim 3$ and gradually decline toward higher modes. Recent 3-D simulations suggest that the latter sources can be important contributors to mode-1 asymmetries.

The goal of this work is to estimate the relative importance of different sources of large-scale nonuniformities in developing asymmetries in OMEGA implosions. This will help to specify improvements in both the OMEGA laser and target fabrication that can lead to improved implosion performance and a better understanding of the physics and robustness of the laser direct-drive approach. Understanding the sources of nonuniformities requires 3-D simulations assuming laser illumination and initial target structural asymmetries that are suggested by direct and indirect measurements and pre-shot target characterization. Results of these simulations are compared with asymmetries of implosion shells measured at different evolution stages, ranging from the beginning of shell acceleration until bang time.

The following sections (1) briefly describe the code ASTER and recent developments; (2) present results of 3-D ASTER simulations of room-temperature and cryogenic implosions and compare these results with experiments; and (3) present our discussion and conclusions.

The Numerical Method

Large-scale nonuniformities in OMEGA implosions were simulated using the 3-D radiation–hydrodynamics code ASTER. This code was tested against 1-D LILAC and two-dimensional (2-D) DRACO12 results, showing good agreement with both results.11

ASTER is an Eulerian code implemented on a spherical grid. Its hydrodynamic algorithm is based on the piecewise-parabolic Godunov method.13 This code uses a 3-D simplified laser-deposition model, which assumes inverse bremsstrahlung for light absorption and includes cross-beam energy transfer (CBET),14 and electron and ion Spitzer thermal transport without flux limitation. ASTER can use various on-the-fly and post-processing diagnostic routines that simulate, for example, neutron spectra and images, burn history, x-ray images, etc.

ASTER is characterized by low numerical noise that allows one to simulate nonuniform implosions without using any kind of diffusion or Fourier filtering to reduce the noise. Figure 149.1

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**Figure 149.1**

Three-dimensional ASTER test simulation assuming 1% perturbation of the mode $(\ell, m) = (10,5)$ in laser deposition. [(a),(b)] The power spectra $\sigma_\ell$ and $\sigma_m$ [see Eq. (1)] of the areal-density perturbation, respectively, at the end of the laser pulse, $t = 2.52$ ns; [(c),(d)] these spectra at $t = 2.805$ ns, which corresponds to $t_{bang} + 30$ ps. (e) An illustration of the shape of the hot spot at the latter time showing an isosurface of $T_i = 1$ keV.
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These spectra are defined as follows:

\[
\sigma_\ell = \sum_{m=-\ell}^{\ell} \sigma_{\ell,m}^2 \quad \text{and} \quad \sigma_m = \sum_{\ell=1}^{\ell_{\text{max}}} \sigma_{\ell,m}^2, \quad (1)
\]

where \( \sigma_{\ell,m}^2 = \left( C_{\ell,m} / C_{00} \right)^2 \) and \( C_{\ell,m} \) are the expansion coefficients on the real (tesseral) spherical harmonics. Figures 149.1(a) and 149.1(b) show these spectra at the end of the laser pulse, \( t = 2.52 \) ns, when the shell’s implosion velocity approaches its maximum. One can see in these figures that the fundamental modes \( \ell = 10 \) and \( m = 5 \) dominate by more than an order of magnitude over the level of background noise introduced by numerical effects. At this time, the fundamental mode experiences mainly secular growth and is insignificantly affected by RT growth because of its relatively large wavelength. Figures 149.1(c) and 149.1(d) show the same spectra at \( t = 2.805 \) ns, which is about 30 ps after bang time, or peak neutron rate. At this time the shell is at maximum compression and is just beginning to move outward. Here, the shell undergoes an efficient RT growth and the perturbations become nonlinear, so that harmonics with \( \ell = 20, 30, \) and 40 and \( m = 10, 15, ... \) are clearly visible and dominate over the background noise. These harmonics are still, however, below the amplitude of the fundamental mode \( (\ell, m) = (10,5) \). Figure 149.1(e) shows the 3-D structure of the hot spot at \( t = 2.805 \) ps, represented by a 1-keV ion temperature isosurface.

Recent developments of ASTER include the capability to simulate radiation transport using multigroup flux-limited diffusion. This development is important since it makes it possible to accurately simulate room-temperature plastic-shell implosions, in which radiative ablation of the inner edge of the dense shell at maximum compression is important. Radiation transport is implemented using the parallel geometric multigrid algorithm. The use of spherical grids with anisotropies near the poles and typically higher resolution in the radial direction (versus angular directions) requires modifications to the standard multigrid relaxation and coarsening procedures to retain optimal efficiency. To treat the polar anisotropies, the algorithm uses nonuniform coarsening strategies, in which the grid is coarsened only in regions and directions that have sufficient isotropic grid coverage. This is combined with line relaxation (using the marching algorithm) in the radial direction. The algorithm is adapted for parallel calculations using a domain decomposition approach similar to that used in the hydrodynamic part of ASTER. Intensive test simulations have been performed to check the accuracy of the radiation-transport routine in ASTER. Results of these simulations showed good agreement with corresponding results obtained using LILAC and DRACO.

Simulation Results

The goal of this study is to identify the effects of large-scale asymmetries in OMEGA implosions with the help of 3-D simulations including a variety of nonuniformities in laser illumination and target structure. The nonuniformities can be chosen only to investigate their effects based on measurements. In the latter case, simulation results are compared with experiments.

Laser-induced nonuniformities include those created by the OMEGA beam-port geometry, target offset, and beam power imbalance, mistiming, and mispointing. The initial target structure nonuniformities can be caused by a variation in the thickness and shape of plastic shells in room-temperature and cryogenic targets and DT-ice shells in cryogenic targets.

The effects of beam imbalance and mistiming in ASTER simulations are included by using the power history of individual laser beams measured on a particular shot. This history is measured before laser light enters the target chamber; therefore, it can be different from the actual on-target value, which is affected by beam-forming optics and protective blast windows. The effects of the latter two are included in simulations by applying time-independent “imbalance correction” factors, which increase or reduce the power of individual beams. These factors are inferred using cross-calibration analysis of time-integrated x-ray images of laser spots from all 60 beams illuminating 4-mm-diam gold sphere targets with a 1-ns square pulse. These targets are chosen to be larger than the nominal OMEGA targets (with radius \( R_t = 430 \) \( \mu \text{m} \)) to avoid the overlapping of laser spots (with radius \( R_0 \approx 430 \) \( \mu \text{m} \)). The imbalance correction factors are typically determined with the accuracy corresponding to about 1% to 2% of the beam power.

Beam mispointing is inferred using the same x-ray data from 4-mm-diam gold targets as in the case of the imbalance measurements. The mispointing data are determined with the accuracy of \( \sim 5 \) \( \mu \text{m} \) and assumed to be fixed in time. These data are provided as horizontal (\( \delta \alpha \)) and vertical (\( \delta \beta \)) displacements of laser spots with respect to their nominal positions on the target surface. ASTER models beam mispointing by
displacing the deposition regions for each beam by the angles of \( \delta \theta = \delta y / R_{\text{dep}} \) and \( \delta \phi = \delta x / R_{\text{dep}} \) in the spherical coordinates, where \( R_{\text{dep}} \) is the radius of the deposition region.

Target offset, or displacement of target center with respect to the laser pointing center, is measured using x-ray imaging\textsuperscript{20} with an accuracy of about \( \pm 3 \) to 5 \( \mu \text{m} \). Offsets are typically small for warm implosions (<5 \( \mu \text{m} \)) and can be significant for cryogenic implosions (~10 to 20 \( \mu \text{m} \)). \textit{ASTER} models target offsets by displacing the deposition region of each beam by angles \( \delta \theta \) and \( \delta \phi \), which are calculated depending on the offset and its directionality and the radius \( R_{\text{dep}} \).

Cryogenic and room-temperature targets are routinely used in OMEGA experiments to study implosion physics. While implosions of these targets share many common physical effects, there are important differences in experimental setups, initial target uniformity, and details of implosion physics that require separate considerations. First we will describe the \textit{ASTER} simulations of room-temperature implosions. These simulations reproduce well the amplitude of observed asymmetries in implosion targets but not the directionality of these asymmetries. Next we will consider the results of cryogenic implosion simulations, which yield similar conclusions: there is good reproduction of the asymmetry amplitudes, but not directionality. The lack of agreement with the directionality can be explained by an inaccuracy of the assumed nonuniformities, which are measured within the time and space resolution of the diagnostics, while some of them are inferred from indirect measurements.

### 1. Room-Temperature Implosions

Room-temperature implosions have several advantages with respect to their cryogenic counterparts that make them a preferable choice for an initial study of large-scale asymmetries: (1) the relatively low fabrication and operation costs that result in an increased shot rate, (2) the ability to add high-Z dopants to the shell that is not fully ablated and confines fuel at stagnation, (3) smaller target offsets, and (4) relatively small initial target nonuniformities. The latter two allow one to concentrate on studying laser-induced asymmetries, whereas the ability to add dopants can help to quantify implosion core asymmetry using self-emission x-ray radiography.

Figure 149.2 shows two warm implosion designs that correspond to OMEGA shots (a) 79638 and (b) 79972. These designs have an IFAR \( \approx 18 \) and 27, respectively, and are relatively stable with respect to high-mode \( (\ell \gtrsim 30) \) RT growth. Shot 79638 (a) uses a 10-atm \( \text{D}_2 \)-filled, 27-\( \mu \text{m} \)-thick plastic \((\text{CH})\) shell. Simulations of this shot are used to study implosion asymmetry during the laser drive and are compared with self-emission x-ray images (at \( h\nu > 1 \text{ keV} \)) of implosion shells.\textsuperscript{21} This x-ray emission comes mainly from a thin layer of plasma located immediately outside the ablation surface. Such images, therefore, can be used to measure the shape and outer radius of implosion shells.

Figure 149.2

Schematic target structure, laser pulse (in black), and simulated neutron rate (in red, left axis) of two warm implosion designs corresponding to OMEGA shots (a) 79638 and (b) 79972.
The design in Fig. 149.2(b) (shot 79972) uses a 15-atm D2–filled, 20-μm-thick plastic shell, which is doped by Ti (1% by atom) at the inner surface to a depth of ~0.1 μm. The purpose of this dopant is to characterize the shape and physical conditions at the fuel–ablator interface using Ti He b line emission (in the 5.45- to 5.65-keV x-ray band) at the time of hot-spot formation since this line emits at $T_e \geq 1$ keV (Ref. 22).

Figures 149.3(a) and 149.3(b) compare experimental and simulated self-emission images, respectively, from shot 79638 at $t = 2.7$ ns (the TIM-5 viewing direction at $\theta = 100.8^\circ$ and $\phi = 270^\circ$ in the OMEGA coordinates). These images represent the shape of the ablation surface at the end of the acceleration phase. The simulations assume the known illumination nonuniformity seeds: OMEGA beam overlap and measured individual beam power histories (which introduce beam imbalance and mistiming) and mispointing (with $\sigma_{\text{rms}} \approx 16$ μm). The measured and simulated images were post-processed to determine perturbations of the ablation surface. Figure 149.4 shows the evolution of the amplitude and phase of mode-2 perturbations in experiment and simulations. The measured mode-2 amplitude grows in time in good agreement with simulations [see Fig. 149.4(a)]. The mode-2 phases are almost independent in time in both experiment and simulations, but they are different by about 40° [see Fig. 149.4(b)]. The latter discrepancy in the phases suggests that the nonuniformity seeds assumed in simulations do not accurately represent the actual seeds.

Figures 149.5(a) and 149.5(b) compare experimental and simulated self-emission images of shot 79638 at $t = 2.9$ ns (in the same viewing direction as in Fig. 149.3). At this time, emissions from the ablation surface (outer ring) and from the core (center spot) are observed simultaneously. The offset of the core (≈5 μm), which is seen as a directional variation of the gap $\Delta R$ between the core edge and ablation surface edge in Fig. 149.5, indicates significant mode-1 perturbations. The offset and its direction are in good agreement in both experimental and simulated images. Simulations show that this offset corresponds to mode-1 distortion of the implosion shell and fuel volume at bang time, as shown in Fig. 149.6. As a result, the simulated neutron yield $4.49 \times 10^{10}$ is reduced to 43% of the yield of the corresponding uniform (1-D) implosion. This yield is a factor of 3 larger, however, than the measured yield $(1.79 \pm 0.09) \times 10^{10}$. Several factors explain the better-simulated performance: (1) an underestimation of the assumed nonuniformity seeds,
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Figure 149.6
(a) Meridional and (b) equatorial cross sections of the distribution of density from simulations of room-temperature shot 79638 at peak neutron production, \( t = 3.02 \) ns. The dashed line in (a) shows the equatorial plane and in (b) the location of the cross-section plane in (a). The solid line inside the dense shell shows the fuel–ablator (D–CH) interface.

Figure 149.7
Narrowband Ti He\( _\beta \) (from 5.45 to 5.65 keV) images for shot 79972 at (a) \( t = t_{bang} - 100 \) ps and (b) \( t = t_{bang} \). The view is opposite the position of the target-mounting stalk.

The observed mode-1 asymmetry in shot 79972 is likely caused by laser-illumination nonuniformities and can be quantified by comparing it with results of ASTER simulations. Figure 149.8 shows simulated distributions of the density and electron temperature in the equatorial cross section of shot 79972, assuming measured individual beam-power histories and pointing misalignment. The assumed perturbations result in mode-1 asymmetry of the dense CH-ablator shell and wide directional motion of the fuel material, which can be seen in Fig. 149.8 as distortion of the hot, low-density central volume occupied by this material. There is also a narrow, high-velocity jet moving in the same direction as the wide flow. This jet develops in the fuel material during successive bouncing of converging shocks produced by the shell during its deceleration. The yellow arrow in Fig. 149.8(a) indicates the directions of the wide flow and jet and points to a dip in the ablating shell into which the jet “drills.”

The solid (color) line inside the dense shell in Fig. 149.8(a) shows the fuel–ablating interface, at which the Ti-doped material is concentrated [see Fig. 149.2(b)]. Simulated images of Ti He\( _\beta \) line emission from this implosion are presented in Fig. 149.9.

(2) missing effects of small-scale mix that were not included in simulations; and/or (3) an inaccuracy in prescribing 1-D physics effects (laser absorption, CBET, heat transport, preheat, etc.).

Another example of significant mode-1 perturbation in OMEGA implosions is presented by shot 79972. Here, mode 1 was measured at a time near target stagnation. Figure 149.7 compares narrowband Ti He\( _\beta \) emission images from this shot at two times, \( t \approx t_{bang} - 100 \) ps and \( t \approx t_{bang} \). The emission limb, which corresponds to the location of the fuel–ablating (D–CH) interface, is consistently brighter on one side in both images, indicating the presence of dominant mode-1 asymmetry in the implosion core. The imager was located opposite the mounting stalk, so the limb asymmetry is unlikely to be caused by the stalk. There is a bright spot inside the limb, which is clearly observed in Fig. 149.7(a) at the earlier time and less clearly in Fig. 149.7(b) at the later time. This spot can be attributed to a jet that penetrates the hot spot and is introduced by the mounting assembly (stalk and glue spot).
Figure 149.8
Equatorial cross sections of the distribution of (a) density and (b) electron temperature in simulations of shot 79972 at peak neutron production, \( t = 1.785 \) ns. The solid line in (a) shows the fuel–ablator interface where Ti-doped CH material is located. The arrow indicates the direction of a wide flow and jet, which develop in the hot-spot plasma because of the mode-1 perturbation. The solid lines in (b) show linearly spaced contours of the electron number density.

These images are calculated for the polar view and correspond to \( t = t_{\text{bang}} - 80 \) ps and \( t = t_{\text{bang}} \), where \( t_{\text{bang}} = 1.785 \) ns [Figs. 149.9(a) and 149.9(b), respectively], and were produced by applying the same spatial (\( \approx 10-\mu \text{m} \)) and temporal (\( \approx 40-\text{ps} \)) smearing as in the experiment. The arrow in Fig. 149.9(a) shows the direction of the wide flow in the hot spot and corresponds to the same direction as in Fig. 149.8(a).

Simulations indicate that the asymmetry of the limb emission observed in shot 79972 (Fig. 149.7) is related to the wide directional motion of the fuel material caused by the mode-1 asymmetry of the shell. The brighter side of the emission limb develops in the direction of this motion. A detailed analysis shows that this brightening is mainly attributed to a local increase of \( T_e \) in the corresponding part of the fuel–ablator interface, while the role of variation in \( n_e \) is less significant [see Fig. 149.8(b)].

By comparing Figs. 149.7 and 149.9, one finds that while experiment and simulations show good agreement with respect to the amplitude of limb brightening, they disagree in directionality of this brightening. This disagreement is similar to that found in the simulations of shot 79638 (see Fig. 149.4) and confirms the claim that illumination nonuniformity seeds assumed in simulations do not accurately represent the real on-target seeds.

2. Cryogenic Implosions

Figure 149.10 shows a target schematic, pulse shape, and neutron history (from 1-D simulations) for shot 77066—one of the best-performing cryogenic OMEGA implosions—in which
about 56 Gbar of hot-spot pressure was inferred. This shot is characterized by an $\alpha \approx 3.2$ and IFAR $\approx 24$ and should be relatively stable with respect to short-scale RT growth. The neutron yield, neutron-averaged (over DT neutrons) ion temperature ($T_i$)h, and hot-spot pressure from uniform (1-D) ASTER simulations of this shot are $2.06 \times 10^{14}$, 3.39 keV, and 138 Gbar, respectively, and using LILAC they are $1.72 \times 10^{14}$, 3.67 keV, and 115 Gbar, respectively. ASTER simulations result in the absorption fraction of laser energy $f_{\text{abs}} = 0.54$ and bang time $t_{\text{bang}} = 2.66$ ns, while these results from LILAC are 0.60 and 2.68 ns, respectively. Table 149.I summarizes these results as well as the results of measurements. The discrepancies between the 1-D ASTER and LILAC results are relatively small and can be attributed to differences in the hydrodynamic methods used (Eulerian piecewise-parabolic method in ASTER and Lagrangian finite-difference scheme in LILAC) and the physical models (e.g., Spitzer versus nonlocal heat transports, respectively).

Three-dimensional simulations of shot 77066 assume all sources of nonuniformities that can be currently quantified. These include the power history of each individual beam, a target offset of 4 $\mu$m (in the direction of $\theta = 38^\circ$ and $\phi = 315^\circ$), and an ice-shell thickness variation with a mode-1 amplitude of 2 $\mu$m (oriented vertically, where the bottom is thinner), which were all measured in this shot. Simulations also assume beam-power imbalance correction factors and mispointing data (with $\sigma_{\text{rms}} = 8.5$ $\mu$m), which were measured in pointing shot 77059.

Figures 149.11(a) and 149.11(b) show, respectively, the equatorial and meridional (at $\phi = 83^\circ$) cross sections of the distribution of density at peak neutron production, $t_p = 3.572$ ns. Figure 149.12 shows a 3-D view of the hot spot at the same moment, where the hot-spot shape is represented by the isosurface $T_i = 900$ eV. The assumed sources of nonuniformities result in a distortion of the dense shell with the dominant mode 1. This mode can be clearly observed in Figs. 149.11(a) and 149.12 as an ~10-$\mu$m shift of the dense shell and hot-spot centroids in the direction $\theta \approx 30^\circ$ and $\phi \approx 83^\circ$ with respect to the initial target center that was located at the origin. The shell is more dense on the side opposite the direction of the shift because of larger laser drive on that side resulting in higher convergence of the shell mass.

Simulations with the assumed asymmetries predict a yield of $8.07 \times 10^{13}$ neutrons and ($T_i$)h = 3.03 keV, therefore reducing the yield to 39% and ($T_i$)h to 89% of the corresponding values of uniform ASTER simulations. The measured neutron yield is $(3.9 \pm 0.2) \times 10^{13}$, which corresponds to 23% of the yield of LILAC simulations (see Table 149.I).

Neutron-averaged ion temperatures in OMEGA implosions are routinely inferred from DD and DT neutron spectra that include the thermal smearing and bulk motion effects in the hot spot. In the case of cryogenic OMEGA implosions, DT neutron spectra are measured by detectors at three different directions: (1) $\theta = 84.98^\circ$ and $\phi = 311.76^\circ$, (2) $\theta = 87.86^\circ$ and $\phi = 161.24^\circ$, and (3) $\theta = 61.30^\circ$ and $\phi = 47.64^\circ$. These directions are indicated by the white dashed arrows in Fig. 149.11(a). The inferred ion temperatures in shot 77066 in these directions are

| Table 149.I: Simulated and measured performance of OMEGA cryogenic shot 77066. |
|-----------------|--------|-------|--------|--------|--------|
| LILAC           | $1.72 \times 10^{14}$ | 3.67   | 115    | 60     | 2.68   |
| 1-D ASTER       | $2.06 \times 10^{14}$ | 3.39   | 138    | 54     | 2.66   |
| 3-D ASTER       | $8.07 \times 10^{13}$ | 3.03   | 88     | 54     | 2.66   |
| Experiment      | $(3.9 \pm 0.2) \times 10^{13}$ | N/A   | 56$^\pm$7 | 58$^\pm$1 | 2.60$\pm$0.05 |

$^a(T_i)_h$ in the absence of bulk motion cannot be measured.
measurements and simulations show good agreement for the amplitude of directional variation of $T_i$: the measured difference between the minimum and maximum temperatures is 0.6 keV, while the simulated difference is 0.9 keV. The latter agreement indicates that simulations correctly reproduce the actual magnitude of hot-spot asymmetry.

Shifts of the simulated neutron spectra in energy in Fig. 149.13 with respect to the unshifted energy of DT neutrons, $E_n = 14.1$ MeV, show a correlation with the direction of the hot-spot shift (see Fig. 149.11) caused by bulk motions. The spectra in red and green in Fig. 149.13 are shifted by $\Delta E = 40$ keV to smaller and larger energies, respectively. These spectral shifts are explained by negative and positive projection components.
of the hot-spot motion (in the direction $\theta \approx 50^\circ$ and $\phi \approx 83^\circ$) in directions 1 and 3, respectively [see Fig. 149.11(a)]. Direction 2 is more perpendicular to the hot-spot motion and has a relatively small, positive projection component. This explains the relatively small shift of the spectrum shown in red in Fig. 149.13.

The spectral shifts in directions 1 and 3 correspond to the neutron-averaged hot-spot velocity components $\vec{v}_f \sim \Delta E / \sqrt{2E_n m_n} \sim 70$ km/s. Correcting this estimate for an angle of $\sim 50^\circ$ between the hot-spot velocity and these directions [i.e., multiplying $\vec{v}_f$ by a factor of $\sim 1/\cos(50^\circ)$], one obtains an estimate of neutron-averaged velocity of the hot spot, $v_f \sim 110$ km/s. Simulations have found that the local flow velocity in the hot spot can substantially vary, taking the maximum value of about a factor of 5 larger than $v_f$ in the hottest, low-density part of the hot spot. This part produces relatively fewer neutrons, however, and, therefore, insignificantly contributes to $v_f$. The shown example demonstrates the importance of spectral-shift measurements to understanding conditions in hot spots.

Discussion and Conclusions

Three-dimensional hydrodynamic simulations using the code ASTER were conducted to investigate sources of large-scale asymmetries in room-temperature and cryogenic OMEGA implosions. Simulations of room-temperature implosions were focused on studying the effects of laser-induced nonuniformities caused by OMEGA beam overlap, target offset, and beam imbalance, mispointing, and mistiming. It was shown that simulations assuming measured sources of these nonuniformities reproduce the amplitude of modes 1 and 2 observed in experiments at an earlier implosion evolution (up to the end of the laser pulse). The development of modes 1 and 2 was studied using self-emission x-ray radiography in up to three viewing directions. The phases of mode 2, however, were not correctly predicted in simulations. The latter indicates that the measured nonuniformity sources assumed in simulations do not accurately represent the actual sources.

Significant mode-1 asymmetry was observed in room-temperature implosions near the bang time. These implosions used plastic-shell targets, in which the inner edge of the shell was doped with titanium to a depth of $\sim 0.1$ $\mu$m. These targets start producing Ti He line emission from the fuel–ablator interface when the temperature there exceeds $\sim 1$ keV. This emission forms bright limbs on x-ray images. Measurements typically find mode-1 asymmetry of the limb brightening, and this asymmetry is well reproduced in simulations assuming measured sources of illumination nonuniformity. The limb asymmetry is attributed to distortions of the dense shell and hot spot with dominant mode 1, which is induced by laser illumination nonuniformities. Simulations suggest that the brighter limb side is developed in the direction of the hot-spot motion caused by these distortions; however, simulations do not reproduce the measured directionality of the limb brightening. This, again, indicates that the nonuniformity sources assumed in simulations do not accurately represent the actual sources.

To study the effects of large-scale asymmetry on performance degradation of cryogenic implosions, 3-D simulations of cryogenic shot 77066 were performed assuming the best currently known sources of the asymmetry. These sources were quantified and include the above-mentioned laser-illumination nonuniformities and nonuniformities caused by the target offset and variation in ice-shell thickness ($\approx 4$ $\mu$m and $\pm 2$ $\mu$m for mode 1, respectively). Simulations showed the development of dominant mode-1 asymmetry in the implosion shell at the time of maximum compression. This results in bulk motions in the hot spot with the neutron average velocity $\sim 100$ km/s in the direction that coincides with the direction of the mode-1 shell asymmetry. These motions result in a directional variation of the hot-spot temperature that is inferred from DT neutron spectra. The experimental and simulated temperatures show good agreement for the amplitude of this variation, but not for directionality of the maximum and minimum temperature measurements. The large-scale asymmetries result in a reduction of the simulated neutron yield to 39% of that of 1-D ASTER simulations, whereas the experimental yield shows 23% of the yield of LILAC simulations—a factor-of-about-2 overperformance in the simulation yields. This disagreement of the hot-spot temperature asymmetry in experiment and simulations suggests that it can be caused by an inaccuracy of the nonuniformity sources assumed in simulations.

Three-dimensional ASTER simulations of room-temperature and cryogenic OMEGA implosions show that large-scale asymmetries of the magnitudes observed in experiments can explain the measured performance degradation in mid- and high-adiabat implosions. Achieving better agreements between experiments and simulations will require a substantial improvement in the measurements of actual on-target nonuniformity sources that are assumed in simulations. In particular, current simulations assuming measured sources do not accurately reproduce directionality of low-mode perturbations (from modes 1 to 3), which limits the prediction capabilities of 3-D simulations.

A technique to correct the measured implosion shell asymmetry by modifying the power distribution of OMEGA
laser beams is under development. This technique uses a 3-D reconstruction of the shape of implosion shells with the help of self-emission x-ray radiography applied in several (three or more) viewing directions. Modifications of the beam-power distribution, which are based on ASTER predictions, will minimize the shell asymmetry and improve implosion performance.

The present study ignored the possibility that large-scale asymmetries in implosion shells can be affected by small-scale perturbations (with $\ell \geq 50$) through mode coupling at the nonlinear stages of perturbation growth. The importance of this effect is unknown and will be studied in future works.

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REFERENCES


First-Principles Equation-of-State Table of Silicon and Its Effects on High-Energy-Density Plasma Simulations

Introduction
As one of the most-abundant elements on Earth, silicon is important to many different fields ranging from the semiconductor industry,1 geophysics,2 photovoltaics,3 planetary and astrophysics,4–6 to inertial confinement fusion (ICF) physics studies.7–9 For ICF applications, silicon has been used as a dopant to ablators in indirect-drive ICF target designs.10 It has also been applied to mitigate laser-imprint effects11,12 and the two-plasmon–decay instability13,14 for multilayer target designs in direct-drive ICF implosions.15 For these high-energy-density (HED) applications, it is essential to know the properties of silicon under extreme conditions. The equation of state (EOS) of silicon is one of such intrinsic properties that are crucial to both ICF and geophysics applications since it is needed for hydrodynamic simulations of ICF implosions and for understanding the geophysics of the earth’s outer core.2

The EOS studies of silicon under megabar (Mbar) pressures began in the 1960s (Ref. 16) using explosive drive. The principal Hugoniot measurements of silicon were continued in the 1970s and 1980s by different groups.17,18 Many surprises were found in our understanding of the behavior of shocks in silicon. For instance, the elastic behavior of shocks was observed in silicon even at Mbar pressures.19 Namely, the lattice reduction related to shock compression may occur only along the shock-propagation direction, instead of hydrostatical lattice-shrinking in all three dimensions. Furthermore, the measured optical emission from shocked silicon was found to be much lower than expected, which has been hypothesized to be caused by the unusually long electron–ion equilibration time in shocked silicon.20–22 These abnormal phenomena have been observed in shock experiments up to ~6-Mbar pressures. What might occur for silicon pressures >10 Mbar remains to be seen. To the best of our knowledge, these anomalies observed in shocked silicon are not fully understood. To this end, a thorough understanding of silicon properties under HED conditions is necessary.

Theoretical investigations on shock compressions of silicon have been performed by classical molecular-dynamics methods,23–25 quantum molecular dynamics simulations based on the density functional theory (DFT),26–29 and path-integral Monte Carlo (PIMC) modeling.27,29 Most of these studies have been devoted to the moderate-pressure regime of $P < 2$ Mbar, while the two most-recent first-principles calculations27,29 extended the Hugoniot pressures from ~1 Mbar to over ~10 Gbar for the first time. These calculations combined the orbital-based–DFT Kohn–Sham molecular-dynamics (KSMD) method, the orbital-free–DFT molecular-dynamics (OFMD) method, and the PIMC simulation. All three first-principles calculations are in good agreement in predicting the principal Hugoniot of silicon, which was found to be ~20% softer than both the extensively used SESAME-EOS model30 (Table 3810) and the quotidian equation-of-state (QEOS) model.31 The predicted softening of silicon should have important implications for HED simulations of silicon plasmas. However, those calculations are concerned with only the plasma conditions along the principal Hugoniot. To study how such a softening behavior of silicon affects HED plasma simulations, we must expand our first-principles calculations to cover a wide range of off-Hugoniot plasma conditions.

In this article, we calculated the EOS for a wide range of silicon plasma conditions by using DFT-based molecular-dynamics simulations. To be specific, we have sampled silicon densities from $\rho = 0.001 \text{ g/cm}^3$ to $\rho = 500 \text{ g/cm}^3$ and temperatures from $T = 2000 \text{ K}$ to $T = 10^8 \text{ K}$. Based on these ab-initio calculations, we have built a first-principles equation-of-state (FPEOS) table of silicon for ICF and HED applications. For off-Hugoniot conditions, we have investigated the differences in pressure and internal energy between FPEOS and SESAME EOS. Implementing the FPEOS table of silicon into the one-dimensional (1-D) hydrocode LILAC32 and two-dimensional (2-D) hydrocode DRACO, we have tested its effects on HED plasma simulations of ICF implosions using a Si ablator. Comparisons with traditional SESAME-EOS simulations illustrated the need for more-accurate EOS tables to precisely design ICF and HED experiments.

The following sections: (1) describe the details of our first-principles calculations; (2) compare the FPEOS and SESAME
EOS for different isochoric plasma conditions (for completeness, the principal Hugoniot comparison is included, even though it has been reported elsewhere29); (3) present the effects of the FPEOS table on HED plasmas through LILAC simulations of ICF implosions using a silicon layer as the ablator; and (4) present our conclusions.

Molecular-Dynamics Simulations Based on the Density Functional Theory

First-principles methods, such as DFT-based quantum molecular dynamics (QMD),33–36 path-integral Monte Carlo,37 and quantum Monte Carlo (QMC),38,39 have been developed over the past decades to understand the properties of materials under extreme conditions. Two different versions of QMD have been implemented by the condensed-matter and HED physics communities. One uses the orbital-based Kohn–Sham formalism40 with the finite-temperature density functional theory, in conjunction with the molecular-dynamics method for ion motion. The other is the orbital-free molecular-dynamics method,41 which is based on the original DFT idea that the free energy of a many-electron system can be written as a function solely depending on the electron density. For most cases, the KSMD method has been proven to be an accurate and efficient method for calculating material properties under high compression at temperatures generally below the electron Fermi temperature $T_F$. It becomes impractical for high-temperature ($T > T_F$) simulations because thermal excitation of electrons requires a large number of orbitals for convergence. The OFMD method is a natural extension of the KSMD method for high-$T$ material simulations, even though it is not as accurate as KSMD. Nevertheless, the pressure difference between KSMD and OFMD calculations is still within ~1% in the overlapping regime of $T \sim T_F$ (valid for both methods), which is acceptable for general ICF/HED applications.

We have used the Vienna ab initio Simulation Package (VASP)42–44 for KSMD simulations, in which electrons are treated quantum mechanically with a plane-wave finite-temperature DFT description. The electrons and ions of the material are in thermodynamic equilibrium with equal temperature ($T_e = T_i$). The electron–ion Coulomb interaction is represented by a projector augmented-wave (PAW) pseudopotential with “frozen” 1s-core electrons. The electron exchange-correlation potential is described by the generalized-gradient approximation (GGA) with the Perdew–Burke–Ernzerhof (PBE) functional.45 Under the Born–Oppenheimer approximation, the self-consistent electron density is first determined for an ion configuration. Then, the classical ions are moved by the combined electronic and ionic forces, using Newton’s equation. This molecular-dynamics procedure is repeated for thousands of time steps from which the thermodynamic EOS quantities such as pressure and internal energy can be directly calculated.

In our KSMD simulations, we have employed the Γ point ($k = 0$) sampling of the Brillouin zone. We used either 32 or 64 Si atoms (depending on density) in a cubic cell with a periodic boundary condition. The cubic cell size is determined from the mass density. The PAW potential of Si included 12 active electrons; the plane-wave cutoff energy was set to 2000 eV. In all KSMD simulations, a sufficient number of bands (varying from 500 to 4100) were included such that the occupation of the highest band was less than $10^{-5}$. The time step varied from $\delta t = 1.5$ fs to $\delta t = 0.085$ fs, respectively, for the lowest and highest densities ($\rho_{\text{min}} = 0.1$ g/cm$^3$ and $\rho_{\text{max}} = 50$ g/cm$^3$). Good convergence was obtained for these parameter sets. The sampled temperature points varied from $T = 2000$ K to a maximum temperature of $T = 500,000$ K. Outside these density and temperature ranges, we switched to the OFMD calculations since the 1s-core electrons must be included in the EOS calculations.

The OFMD method41 originated from the true spirit of the Hohenberg–Kohn theorem,46 i.e., the free energy of an electron–ion system at any ion configuration can be written as a function of the electron density. The kinetic energy of the electrons is currently represented by the Thomas–Fermi functional plus the von Weizsäcker correction that takes into account the gradient of electron density. These terms were obtained from the semiclassical expansion of the partition function up to the first order. In OFMD simulations, all electrons, both bound and free, are treated equally. The divergence of the electron-nucleus potential is regularized for each thermodynamic condition through a similar procedure of generating the norm-conserving pseudopotential as the PAW treatment. The cutoff radius is chosen to be less than 10% of the Wigner–Seitz radius to avoid an overlap of regularized ion spheres. The exchange-correlation function is expressed in the local density approximation of Perdew and Zunger.47

At each time step of an OFMD simulation, the electron free energy for an ionic configuration is first minimized in terms of the local electron density. Then, the classical ions are moved by the combined electronic and ionic forces, the same as in the KSMD procedure. In our OFMD simulations of silicon plasmas, we used 128 atoms in a cubic cell with periodic boundary conditions. The time step varied from $\delta t = 0.144$ fs to $\delta t = 6 \times 10^{-5}$ fs, respectively, for the lowest-density/lowest-temperature ($\rho = 0.001$ g/cm$^3$ and $T = 125,000$ K) point and
the highest-density/highest-temperature \((\rho = 500 \text{ g/cm}^3\) and \(T = 10^8 \text{ K}\) point. Finally, the thermodynamic EOS quantities were statistically evaluated from the molecular-dynamics (MD) propagation of the system (5000 to 100,000 steps, depending on the density).

For each isochoric curve, we examined the EOS quantities for the overlapping temperature points between the KSMD and OFMD calculations. We made the transition from KSMD to OFMD at the temperature point where their differences were the smallest (within \(\pm 1\%\)). Carrying out these calculations for a wide range of silicon plasma conditions, we obtained both pressure and internal energies for all the sampled density and temperature points \((\rho = 0.001\) to \(500 \text{ g/cm}^3\) and \(T = 2000\) to \(10^8 \text{ K}\)). As an example, in Fig. 149.14 we plot the total pressures as a function of the silicon plasma temperature for each of the sampled isochoric curves.

![Figure 149.14](image)

Silicon pressure as a function of temperature for all densities \((\rho = 0.001\) to \(500 \text{ g/cm}^3)\) scanned by our first-principles (KSMD + OFMD) calculations.

**Comparison Between FPEOS and SESAME EOS**

From the FPEOS table, we can derive the principal Hugoniot curve for silicon shocks by using the Rankine–Hugoniot equation. The initial state is chosen to be solid silicon \((\rho_0 = 2.329 \text{ g/cm}^3)\) in its diamond phase at ambient pressure \((P_0 = 1 \text{ bar})\). We compare the FPEOS Hugoniot with the one derived from the extensively used SESAME-EOS model \((\text{SESAME 3810 table})\) in Fig. 149.15(a), in which the Hugoniot pressure spanning more than five orders of magnitude is plotted as a function of the shock density. The SESAME-EOS model was based on the chemical picture of matter, meaning that the total free energy can be decomposed into the cold curve, the ionic excitation, and the electron thermal excitation. It was typically constructed (constrained) by the best-available experimental data (typically limited). Specifically, for SESAME 3810 (Si) constructed in 1997, the EOS below the solid–liquid phase transition was based on experimental Hugoniot data.\(^{16–18}\) For conditions above the liquid phase transition, the EOS was constructed such that the shock Hugoniot was “similar” to germanium \((\text{SESAME 3950})\) up to \(4.4 \text{ Mbar}\). The ion thermal contribution is based on a Debye model with a correction for the liquid’s specific heat beyond the melt temperature.\(^{48}\) The correction also ensures that in the high-temperature limit, the proper model (ideal gas) that is recovered will give a shock Hugoniot compression ratio \(\rho/\rho_0 = 4\). The Hugoniot comparison in Fig. 149.15(a) indicates that under shock compression, silicon is much softer than predicts by the traditional chemical picture of materials.\(^{29}\) For example, at a constant pressure of \(~20 \text{ Mbar}\), the SESAME 3810 table predicts a shock density of \(\rho \approx 6.3 \text{ g/cm}^3\), while the FPEOS table gives a much-higher...
shock density of $\rho \approx 7.7$ g/cm$^3$. Namely, the FPEOS table
predicts that silicon under 10- to 1000-Mbar pressures is
$\sim 20\%$ softer than SESAME 3810. For the same shock density
at $\rho = 8$ g/cm$^3$, the SESAME 3810 model predicts a shock
pressure of $P \approx 73.4$ Mbar, which is more than $3 \times$ higher than
the FPEOS case ($P \approx 24$ Mbar). Figure 149.15(a) indicates that
the maximum compression ($\rho/\rho_0$) changes from the SESAME-
predicted value of $\sim 4.6$ to 5.0 in FPEOS. Finally, in the same
figure, we have plotted the existing experimental data$^{16-18}$
which are represented by the different symbols. These Hugoniot
data were obtained from explosively driven shock experiments.
To the best of our knowledge, no published data exist for laser-
shock Hugoniot measurements in pressures above 10 Mbar. The
opacity of Si for most velocity interferometer system for any
reflector (VISAR) laser wavelengths$^49$ is one of the hurdles
for accurate shock measurements in silicon. Nevertheless, it
is shown in Fig. 149.15 that the explosively driven shock data
up to $\sim 2$ Mbar agree well with our calculations, which seems
also to indicate the softening of silicon under compression.
It is noted that at the measured highest shock density of $\rho =
4.6$ g/cm$^3$, the SESAME-EOS–predicted pressure is at least $2 \times$
higher than the experimental value of $P \approx 2$ Mbar.

To further examine the properties of shocked silicon, we
have calculated the heat capacity $C_v$ along its principal Hugo-
niot. Because $C_v$ is a measure of the energy change with respect
to temperature at a fixed volume, it can give some indication of
how rapidly the entropy is increasing with temperature in a sil-
icon shock. The obtained $C_v$ results are plotted in Fig. 149.15(b)
as a function of the Hugoniot density for both SESAME 3810
(dashed red line) and FPEOS (solid blue line). In Fig. 149.15(b),
we also plot three horizontal lines to indicate the expected heat
capacities for ideal-gas plasmas of three different ionization
stages of Si$^{4+}$, Si$^{12+}$, and Si$^{14+}$, respectively. For instance, the
lowest dashed black line represents the ideal-gas plasma that
includes only Si$^{4+}$ and free electrons without any interactions.
Since the electron ionization process acts like a “heat sink” for
the system, one expects the heat capacity to increase during
the ionization of bound electrons. This is especially true for
the innermost shell electrons because of the large energy gaps
between the L-shell and K-shell electrons. This is exactly what
can be seen in Fig. 149.15(b), where the FPEOS calculation
(solid blue line) gives a peak of $C_v$ near the peak compression
at $\rho \approx 11.5$ g/cm$^3$ [see Fig. 149.15(a)]. After the 1s-electron ion-
ization is completed, the heat capacity approaches the ideal-gas
limit (horizontal dashed pink line) as a fully ionized Si plasma
is formed. The SESAME 3810–predicted $C_v$ has a similar trend,
but the same value of $C_v$ is reached at a smaller density. In other
words, at the same density the FPEOS-predicted $C_v$ is $\sim 50\%$
lower than the SESAME 3810 case, meaning that less entropy
increase is expected in FPEOS. By referring to the ideal-gas $C_v$,
one can argue that the same ionization stage is first reached at
much-lower densities in SESAME 3810 than in FPEOS. Again,
all of these features are consistent with the higher compress-
ibility of silicon predicted by FPEOS.

Next, we compare the pressure and internal energy of silicon plasmas for off-Hugoniot conditions between FPEOS (solid blue line) and SESAME 3810 (dashed red line) in Figs. 149.16–149.18. Figures 149.16(a) and 149.17(a) show the pressure as a function of plasma temperature, respectively, for silicon densities of $\rho = 5$ g/cm$^3$ and $\rho = 10$ g/cm$^3$, while the internal energy comparisons are made in Figs. 149.16(b) and 149.17(b). One sees in Fig. 149.16(a) that the SESAME pressure is $\sim 10\%$ lower than FPEOS for temperatures $T < 10^4$ K, but it

\begin{figure}[h]
\centering
\includegraphics[width=\textwidth]{figure149.16.png}
\caption{Figure 149.16}
\end{figure}
The off-Hugoniot equation-of-state comparisons between FPEOS and
SESAME 3810. The (a) pressures and (b) internal energies are plotted as
functions of temperature for a silicon density of $\rho = 5$ g/cm$^3$. 

\begin{table}[h]
\centering
\begin{tabular}{|c|c|c|}
\hline
Temperature (K) & Pressure (Mbar) & Internal energy (eV/atom) \\
\hline
10^3 & 10^2 & 10^1 \\
10^4 & 10^3 & 10^2 \\
10^5 & 10^4 & 10^3 \\
10^6 & 10^5 & 10^4 \\
10^7 & 10^6 & 10^5 \\
\hline
\end{tabular}
\caption{Table of Silicon and Its Effects on High-Energy-Density Plasma Simulations}
\end{table}
First-Principles equation-of-state table of silicon and its effects on high-energy-density plasma simulations

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Reverses for $10^4 < T < 10^6$ K with a “crossover” temperature at $T \sim 10^4$ K ($\sim 1$ eV). The pressure difference between FPEOS and SESAME 3810 reaches a maximum of $\sim 50\%$ in the warm dense regime ($T \sim 10^5$ K) at this density ($\rho = 5$ g/cm$^3$). This is the regime in which both electron degeneracy and strong ion–ion coupling play significant roles in determining the EOS. The internal energy comparison in Fig. 149.16(b) shows a similar trend, although the difference is only $\sim 20\%$. For high temperatures of $T > 10^6$ K, both FPEOS and SESAME 3810 are in good agreement with each other as the two EOS tables correctly approach the ideal gas limit. Figure 149.17 shows similar EOS comparisons for $\rho = 10$ g/cm$^3$. At this higher density note that the crossover temperature now moves to near $10^5$ K ($\sim 10$ eV), and the maximum difference in pressure between FPEOS and SESAME 3810 reduces to $\sim 20\%$. The difference in internal energy in Fig. 149.17(b) is also reduced when compared to Fig. 149.16(b).

Finally, we explore two other isochores at high densities of $\rho = 50$ g/cm$^3$ and $\rho = 500$ g/cm$^3$, respectively, in Figs. 149.18(a) and 149.18(b). Again, the two panels compare the pressures of FPEOS with SESAME 3810 at various temperatures. Figure 149.18(b) indicates that both FPEOS and SESAME 3810 are very close to each other at this high density of $\rho = 500$ g/cm$^3$, even though SESAME 3810 gives a slightly higher pressure over the entire temperature range (no more crossover is seen between...
the two EOS’s). Both EOS tables are in better agreement with each other in this electron-degeneracy-dominated regime. For the intermediate density of $\rho = 50$ g/cm$^3$, Fig. 149.18(a) still shows a trend similar to the one seen in Figs. 149.16 and 149.17. Namely, the SESAME 3810 model still underestimates the pressure for the low-$T$ regime ($T < 10^6$ K). With these large EOS differences identified in both on-Hugoniot and off-Hugoniot warm-dense-plasma conditions, we expect to see significant effects on HED plasma simulations between using the newly established FPEOS and using the SESAME 3810 for silicon.

**EOS Effects on HED Plasma Simulations Involving Silicon**

To examine the EOS effects on HED plasma simulations, we have implemented our FPEOS table of silicon into our radiation–hydrodynamics codes LILAC and DRACO. We have extrapolated our EOS results for temperatures outside our calculation range (2000 K to $10^8$ K). With the implementation of the FPEOS table, we can investigate its effects on HED simulations involving silicon plasmas. Since in an ICF implosion the capsule generally undergoes a path sweeping through many different density and temperature conditions, integrated ICF implosion simulations would be more suitable for examining EOS effects. As an example, we consider a NIF (National Ignition Facility)-type direct-drive implosion for the same FPEOS table$^{52,53}$ and the same first-principles opacity table$^{54}$ of DT, so that the EOS tests solely focused on the silicon ablator layer. In Fig. 149.20, we plot the density and temperature profile snapshot at $t = 0.9$ ns as a function of target radius for the two simulations. At this time, the shock is still propagating inside the Si layer (the shock front is located at $R \sim 1180 \mu m$). Figure 149.20 indicates that (1) the shock density in FPEOS is ~20% higher than the SESAME simulation and (2) the shock in the SESAME simulation is ahead of the FPEOS case, giving a shock-speed difference of ~10%. These features can be understood by considering the softening of silicon shock in FPEOS (see Fig. 149.15). Namely, the identical laser drive gives the same ablation pressure in the two simulations; for the same shock pressure ($P_s$), the FPEOS simulation will give ~20%-higher shock density ($\rho_s$) as the Hugoniot curve seen in Fig. 149.15(a). Since the shock speed depends on the shock density through $V_s = \sqrt{\frac{P_s}{\rho_0} \left(1 - \frac{\rho_0}{\rho_s}\right)}$, one can see that for the same $P_s$, the ~20%-higher shock density in FPEOS will give an ~10%-smaller shock speed than the SESAME case. Figure 149.20 also indicates that the shock temperature is ~20% higher in FPEOS.

![Figure 149.19](TC13231JR)

The laser pulse shape and target dimensions for implosion simulations to test the silicon EOS effects. The capsule consists of a 40-$\mu m$ Si layer filled with 3 atm of DT gas. The initially target radius $R = 1200 \mu m$.

![Figure 149.20](TC13232JR)

Comparisons of density and electron temperature profiles predicted by the two LILAC simulations using FPEOS (solid blue lines) and SESAME 3810 (dashed red lines) EOS models. The snapshot was taken at $t = 0.9$ ns, when the first shock was still propagating in the silicon layer.
Figure 149.21
Same as Fig. 149.20 but for different implosion times: (a) $t = 5.4$ ns (in flight of the imploping shell) and (b) $t = 7.9$ ns (the end of shell acceleration).

Figure 149.22
Comparisons of density and ion temperature profiles predicted by the two LILAC simulations using FPEOS (solid blue lines) and SESAME 3810 (dashed red line) EOS models. Peak neutron production is at $t \approx 9.0$ ns.

As the implosion proceeds, Fig. 149.21 shows the density and temperature profiles during the in-flight stage of $t = 5.4$ ns [Fig. 149.21(a)] and at the end of acceleration of $t = 7.9$ ns [Fig. 149.21(b)]. One sees from Fig. 149.21 that the peak density of the shell from the FPEOS simulation is always $\sim 20\%$ higher than the SESAME 3810 case. This can be attributed to the greater compressibility of silicon predicted by FPEOS. Except for the difference in peak density, the two simulations give very similar density and temperature profiles for the imploding shell. Some difference in the back surface of the shell appears only at the end of the acceleration phase, as indicated by Fig. 149.21(b). Note that the coronal plasma conditions are also very similar.

Figure 149.23
Comparisons of (a) the areal density $\rho R$ and (b) the total neutron yield as functions of time for the two LILAC simulations using FPEOS (solid blue lines) and SESAME 3810 (dashed red line) EOS models.
in the two cases, as the EOS difference becomes very small at high temperatures of $T > 10^6$ K. Figure 149.21 also shows an interesting double-ablation-front feature, which can develop in such mid-$Z$–ablator implosions$^{55}$ because of the significant radiation preheat from coronal emissions. The $\sim 20\%$ difference in peak density in the two simulations can have significant consequences when the imploding shell stagnates. Figure 149.22 displays the situation at the time of peak neutron production (near peak compression). Again, the figure shows the density and ion temperature as functions of the target radius. The maximum density reached in the FPEOS simulation is $\rho_p = 271.9$ g/cm$^3$, in contrast to the SESAME 3810–predicted $\rho_p = 185.5$ g/cm$^3$. The Si shell is converged slightly more in FPEOS than SESAME, resulting in a somewhat different hot-spot radius ($R_{hs} = 30.5$ $\mu$m versus $R_{hs} = 33.6$ $\mu$m). Consequently, the maximum ion temperature is increased from $T_i \simeq 3.07$ keV (SESAME) to $T_i \simeq 3.45$ keV (FPEOS).

Finally, we plot the history of the compression areal density ($\rho R$) and neutron yield, respectively, in Figs. 149.23(a) and 149.23(b) for the two implosion simulations. One sees from Fig. 149.23(a) that the peak areal density reaches a value of $\rho R = 1.38$ g/cm$^2$ in FPEOS, which is $\sim 30\%$ higher than the SESAME simulation. The total neutron yield predicted by FPEOS, shown by Fig. 149.23(b), is increased by more than $\sim 70\%$ with respect to the SESAME case [$Y = 5.0 \times 10^{14}$ (FPEOS) versus $Y = 2.9 \times 10^{14}$ (SESAME)]. As a result, the EOS difference can have significant consequences on predicting the 1-D target performance. This illustrates the importance of having a more-accurate EOS table in the 1-D hydrodynamic designs of ICF/HED experiments.

Conclusion

We have applied DFT-based molecular-dynamics simulation methods to investigate the EOS of silicon, spanning a wide range of plasma conditions from $\rho = 0.001$ to 500 g/cm$^3$ and $T = 2000$ to $10^8$ K. The resulting pressures and internal energies have been assembled into a first-principles equation-of-state table, which is studied in detail by comparing it with the extensively used SESAME 3810 table of silicon. We found that the shock Hugoniot of silicon is $\sim 20\%$ softer in FPEOS than SESAME 3810. For off-Hugoniot warm-dense-plasma conditions, the pressure difference can reach $\sim 50\%$ between FPEOS and SESAME 3810, while the internal energy difference is within $\sim 20\%$. After implementing the FPEOS table of silicon into our 1-D radiation–hydrodynamics code LILAC, we tested its effects on HED plasma simulation by carrying out hydro-simulations of an ICF implosion with a Si shell using either FPEOS or SESAME 3810. The simulation results showed (a) the FPEOS-predicted shock density is $\sim 20\%$ higher than the SESAME 3810 case (accordingly, the shock speed is $\sim 10\%$ lower in the former case); (b) the peak density of the imploding Si shell is $\sim 20\%$ larger in FPEOS than in SESAME; (c) the maximum density at peak compression is higher by $\sim 40\%$; and (d) the final areal density and yield predicted by FPEOS are respectively higher by $\sim 30\%$ and $\sim 70\%$, with respect to the SESAME simulation. The observed differences in target performance can be attributed to the different compressibility of silicon predicted by FPEOS. These studies illustrate the importance of having a more-accurate EOS table in order to precisely design ICF/HED experiments. Hopefully these results will facilitate shock-wave experiments in the untested high-pressure (>10-Mbar) regime.

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REFERENCES


In direct-drive inertial confinement fusion (ICF) ignition designs, a cryogenic deuterium–tritium (DT) shell surrounding a vapor and encased in a thin ablator (<10 μm) is symmetrically heated with nominally identical laser beams. In most designs, laser ablation launches single or multiple shocks through the converging shell and into the vapor region. The shock-transit stage of the implosion is followed by a deceleration phase, where the kinetic energy of the converging shell is converted to the internal energy of the hot spot. Thermonuclear fusion reactions are initiated in both the shock phase and the compression phase once sufficiently high temperatures and densities are reached. To achieve conditions relevant for ignition implosion designs, the hot-spot size must exceed the mean free path of fusing ions and the mean free path of the alpha particles.

Previous experiments on OMEGA have reported anomalous $Y_{DT}/Y_{DD}$ values (different by as much as a factor of 4) with the measured pre-shot fuel composition and experimentally inferred ion temperatures in room-temperature implosions. Several studies suggest that species separation of the hydrogen isotope resulting from multifluid effects is likely responsible for the observed discrepancies in the yield ratios. These classes of implosions—for example, exploding pushers that use thin glass (~3-μm SiO₂) or thin CH (<16-μm) shells—are, however, characterized by fusion reactions that occur predominantly during the shock phase at very high temperatures (~10 keV) and relatively low densities (~10 mg/cm³). The mean free path for 90° deflection is given by $\lambda_{ii} \sim T_i^{2/3} / Z_i^2 Z^2 \rho$ (Ref. 4) for ions of charge $Z_i$, average ion temperature $T_i$, ion charge $Z$, and density $\rho$. Conditions during the shock phase result in large mean-free-path lengths of the ions relative to the size of the fusing-plasma region (see Table 149.II). These conditions are also typical of ignition-relevant direct-drive cryogenic implosions during the shock phase; however, cryogenic targets differ from exploding-pusher targets in two respects: First, most of the neutron yield in a cryogenic implosion occurs later in the implosion, during the compression phase, when the kinetic energy is converted to the internal energy of the hot spot. Simulations using the spherically symmetric hydrodynamics code LILAC indicate that nearly 99% of the yield occurs in this compression phase. Second, compression yields occur at significantly higher densities (~20 g/cm³) and lower temperatures (~3 keV), leading to mean free paths of thermal ions that are much shorter than the hot-spot size. Nonlocal transport of energetic ions is therefore not expected to significantly influence yields during compression. Evidence of fuel species separation that persists into the compression phase would suggest a reduction in the number of alpha particles produced from the dominant D–T fusion reactions. In ignition-scalable cryogenic implosions described in this article, however, measurements give the first evidence that species separation does not persist from the shock phase and has an insignificant influence on.

### Table 149.II: Calculated Implosion Parameters

<table>
<thead>
<tr>
<th>Implosion Type</th>
<th>$\rho$ (g/cm³)</th>
<th>$T_i$ (keV)</th>
<th>$\lambda_{ii}$ (μm)</th>
<th>$R_{shell}$ (μm)</th>
</tr>
</thead>
<tbody>
<tr>
<td>Exploding pusher:</td>
<td></td>
<td></td>
<td></td>
<td></td>
</tr>
<tr>
<td>Shock phase</td>
<td>0.03</td>
<td>10</td>
<td>400</td>
<td>100</td>
</tr>
<tr>
<td>Cryogenic implosions:</td>
<td></td>
<td></td>
<td></td>
<td></td>
</tr>
<tr>
<td>Shock phase: vapor</td>
<td>0.1</td>
<td>8</td>
<td>80</td>
<td>100</td>
</tr>
<tr>
<td>Shock phase: cold-fuel layer</td>
<td>6.0</td>
<td>0.02</td>
<td>0.0002</td>
<td>$\Delta R_{shell} \sim 10$</td>
</tr>
<tr>
<td>Compression phase</td>
<td>20.0</td>
<td>3</td>
<td>0.08</td>
<td>25</td>
</tr>
</tbody>
</table>
the yield ratio into the compression phase in direct-drive D–T cryogenic implosions consisting of a near-equimolar mixture of deuterium and tritium.7

Direct-drive ICF targets consisting of a deuterated plastic (ablator) shell with a 460-μm outer radius are imploded at an ignition-scalable, on-target laser intensity with a laser energy of ~25 kJ (Ref. 8). The implosion velocity (Vimp, defined as the velocity of the compressing shell when the kinetic energy of the shell is at a maximum) ranged from 3.5 × 107 cm/s to 4 × 107 cm/s and the adiabat (α, defined as the ratio of the pressure to the Thomas–Fermi pressure at maximum shell density) ranged from 2.4 to 5. The average ion temperature T_i in this class of implosions is varied by adjusting the implosion velocity, T_i ~ V^{1.1}\text{imp} which, in turn, is governed by the thickness of the cryogenic DT layers or the CH (CD) ablator. The capsule is filled by a permeation technique at a temperature of 300 K, where increasing pressure is applied to the outside of the shell, allowing the gas to diffuse inside. Fill rates for a typical cryogenic target are carefully controlled by holding the pressure ramp rate at ~ 1 atm/min to ensure the integrity of the shell is not compromised.9 At the final fill pressure (between 400 and 800 atm) depending on the desired ice thickness, the capsule is cooled to a few mK below the triple point (~19.8 K), producing a DT ice layer ranging from 40 to 90 μm in thickness. The primary nuclear-fusion reactions examined in this study are given by

\[ D + D \rightarrow ^3\text{He} + n + (3.27 \text{ MeV}), \]  
(1)

\[ D + T \rightarrow ^4\text{He} + n + (17.6 \text{ MeV}). \]  
(2)

The neutron yields are measured using the time-of-flight (nTOF) diagnostics positioned around the OMEGA target chamber. The fusion yield is given by

\[ Y_n^{DD/DT} = \int f_D f_T \rho(\vec{r}, t)^2 \langle \sigma v \rangle_{DT/DD} \times d\vec{r}^3 dt / (1 + \delta_{DD}), \]

where f_D and f_T are the atomic fractions of the reactants, ρ is the fuel-mass density, ⟨σv⟩ is the Maxwellian-averaged reactivity for the D–T or D–D fusion reaction (which scales as ~T_3^3/2 for the D–T reaction and ~T_3 for the D–D reaction for the typical temperatures in OMEGA implosions), T_i is the average ion temperature, m is the average reactant mass, and δ_{DD} = 1 for DD and 0 for DT to account for double counting of the identical D–D reaction.

The primary D–T yields observed in cryogenic experiments are always lower relative to radiation–hydrodynamics codes that assume spherical symmetry and include the deposition of the laser energy through collisional absorption and account for laser–plasma interactions such as cross-beam energy transfer (CBET).10 These codes include nonlocal heat conduction10 and multigroup diffusive radiative transport.11 Several multidimensional effects that reduce the overall yield relative to these state-of-the-art spherically symmetric fluid codes have been proposed, including nonuniformity growth caused by beam-to-beam energy imbalance,12 on-target beam misalignment,13 single-laser-beam nonuniformity,13 and isolated defects on the target14 that potentially reduce T_i and/or fuel density. All these mechanisms include only hydrodynamic effects and do not exhibit yield ratio anomalies. More recently, an extension to fluid codes has been proposed. Calculations that include plasma barotropic diffusion,15,16 where hydrogen isotope species separation occurs during the shock phase into the hot spot because gradients in pressure and temperature, have been shown to influence the D–T and D–D fusion yields differently. Two phases of an ICF implosion have been analyzed using this model: the shock phase (when the shock is moving through the vapor toward the center of the capsule) followed by the rebound phase (outward-going shock). It was reported that during the shock phase, up to 5% of the deuterium can leave the fuel volume for an equimolar mixture of deuterium and tritium. During the subsequent shock-rebound phase, the barotropic diffusion rate decreases to zero and the ability for fuel to leave the volume is significantly reduced if not eliminated. Since the D–D fusion and D–T fusion reactivity are well-known17 and the composition of the fuel is measured prior to the implosion, the ratio of the neutron yields (Y_{DT}/Y_{DD}) from these reactions should follow a calculable trend with the measured ion temperature with the exclusion of diffusive effects. Table 149.II summarizes the mass-fuel density (ρ) and the key implosion parameters to calculate the ion–ion mean free path (λ_{ij}) for the plasma conditions across the class of implosions discussed earlier in this article. The radius of the shell (R_{shell}) is calculated from simulations for the different phases of the implosion.

As shown in Table 149.II, the mean free path during the shock phase for the ions at the relevant average ion temperature approaches the radius of the shell. At this time, however, the vapor region is surrounded by a relatively cold (~20-eV) and highly dense DT-fuel layer. The energetic and thermal ions that escape the vapor phase do not leave the target and instead are stopped in the cold dense DT shell. At peak neutron production, the mean free path is several orders of magnitude smaller (~10^{-2}) than the boundary of the cold-fuel shell.
Cryogenic implosions are additionally different from shock-driven implosions that have been studied previously since the shell material is also made of DT fuel. When the shell decelerates in the compression stage of any ICF implosion, the cold fuel ablates into the hot spot. Simulations using the code LILAC indicate that, in the case of cryogenic layered DT implosions, nearly $5 \times$ the mass of the original vapor is injected into the hot spot through the ablation process, which is the primary source of the fusion neutrons during compression. Therefore, it would be expected that the ions that are stopped in the cold-fuel shell would be restored into the hot core during the compression phase, compensating for any loss of particles that may have occurred earlier in the implosion.

For this analysis, the yields ($Y_{DT}$ and $Y_{DD}$) for the different reactions are measured along the same diagnostic line of sight using the 13.4-m high-resolution nTOF spectrometer. This diagnostic uses several microchannel-plate–based phototubes to increase the dynamic range required to measure the primary DT and DD signal in a single line of sight. The yield is inferred by fitting the recorded signal with a forward-fit approach using a relativistic model of the neutron distribution. Cross-calibration of the neutron diagnostics with standard measurements on OMEGA give an uncertainty in the D–T and D–D yields of 5% and 9%, respectively. In ignition-scalable implosions, the neutron yield is attenuated by the compressed fuel at peak neutron production (see Table 149.II). To recover the fusion birth yield, a correction to the measured yields must be included as a function of the areal density from the compressed fuel. The elastic scattering is proportional to the areal density of the implosion, which is inferred from separate measurements. The transmission factors ($\eta_{DT}$ and $\eta_{DD}$) for the neutrons from the two fusion reactions are calculated using the well-known total scattering cross sections and the measured areal density. Typical values of these transmission factors for an areal density of 220 mg/cm² are 4% and 10% for the DT and DD neutrons, respectively. With the areal densities achieved on OMEGA, multiple scattering can be neglected, thereby providing an ideal platform to study the effects of fuel-species separation in ignition-scalable implosions. By adding the uncertainty of the D–T and D–D yields, the attenuation of the yield from the compressed fuel and the reaction rate for both of the primary reactions in quadrature, an error of 10% for the $Y_{DT}/Y_{DD}$ ratio can be inferred.

As indicated earlier, it is important to know the ion temperature in the implosion and the fuel composition. The energy spread of the primary neutron distribution provides a good measure of the ion temperature characteristics of peak neutron production. If mass flow within the reaction region is present, this effect can lead to a broadening of peak distribution and an incorrect interpretation of ion temperature. On OMEGA, several nTOF detectors measure the width of the DT neutron spectrum temperature from various lines of sight around the target chamber. The ion temperature inferred from the width of the neutron spectrum in ignition-scalable implosions can vary up to ~1 keV across the three different detectors. Simulations indicate that this variation in the temperature is caused by bulk fluid motion of the fusing plasma. The uncertainty in the inferred ion temperature, excluding effects caused by bulk fluid motion, is ±0.2 keV for implosions between 2 keV and 5 keV. To minimize the effect of bulk motion, the minimum ion temperature will be used in this analysis as an approximation of the thermal temperature. It should be noted that the implosions that can vary up to 1 keV are only 3% of the data points. The remaining 2σ that vary up to 0.8 keV account for 90% of the implosion analyzed. Using this variation in the ion temperature, the calculated fuel fraction has an uncertainty of less than 7%.

The observed reaction yield ratio is plotted as a function of the minimum ion temperature in Fig. 149.24 for each cryogenic shot on OMEGA (35 experimental campaigns with 120 implosions taken over a period of three years). The composition of the DT inventory in the assay volume is periodically measured on OMEGA to within an accuracy of 1.5%. In this case, the gas used to fill the targets was taken at various stages during the pressurization of the fuel so that the deuterium-to-tritium (D:T) concentration could be calculated. Over time, the tritium supply in the system gradually changes as a result of beta decay of the hydrogen isotope. Figure 149.24 also shows the calculated ratios using the measured fuel fraction and the minimum ion temperature. The measured ratios show good agreement with the calculated ratios expected from the DT inventory and experimentally inferred ion temperatures. It should be noted that while the accuracy of the fuel composition in the both the assay volume and the pressurized system are well understood, an extrapolation of the fuel fraction is required of the gas composition during the fill process in the permeation cell that is used to fill cryogenic capsules. A project is underway to better characterize the fuel composition of the gas as it is sent into the permeation cell used to fill the capsules. Presently, this effect is known to change the composition between 3% and 5%. The calculated reaction yield ratios follow the form $Y_{DT}/Y_{DD} \sim 2T^{0.4} (f_T/f_D)$ using the assumption that hydrodynamic models of an ICF implosion predict that the reactant density ratio ($f_T/f_D$) is spatially and temporally constant during all phases. This indicates that additional effects that change this ratio or the volume over
which each of the D–T and D–D reactions are produced do not significantly influence yields from the hot-spot stagnation. Pre-shot fuel fractions are measured during each fill process for every campaign. Variations in the yield ratio measurements resulting from the fuel composition are reflected in Fig. 149.24 with the solid and dashed lines representing the initial and final measurement, respectively, before the inventory underwent a scheduled refinement.

The measured D–T and D–D yield ratios and the ion temperature are used to instead infer a fuel fraction ($f_{D}$ and $f_{T}$) for each of these shots. The measured fuel fraction is compared against values inferred from nuclear measurements in Fig. 149.25. The average of the ratio of the inferred fuel fraction from the nuclear measurement over the composition obtained from the permeation cell is 1.07 with a standard deviation of 0.09. Although error on the mean is small with 1% for 120 implosions used for this study, given the 10% systematic error on the $Y_{DT}/Y_{DD}$ ratio, both measurements of the fuel fractions are consistent within the experimental uncertainties.

In summary, nuclear measurements of the D–T to D–D yield ratio from OMEGA cryogenic implosions scale predictably with the known composition of the fuel and experimentally inferred ion temperatures with a calculated 7% systematic offset. These observations indicate that multifluid effects that may take place during the shock phase of the implosion (and potentially influence species profiles in the compressing target) do not persist into the subsequent compression phase of the implosion. A plausible explanation for this rests on the composition of the target; the shell is also DT fuel. During the deceleration phase of cryogenic DT implosions, the fuel from the inner DT wall is ablated into the hot spot. Simula-
tions indicate that nearly $5 \times$ the mass of the neutron-emitting region is from the ablation of the cold DT shell. Therefore, the energetic ions that may be lost because of their long mean free paths earlier in the implosion return to the hot spot during peak neutron production, leading to an unchanged fusion yield ratio. These observations indicate that multifluid effects have an insignificant influence on the yield ratio in ignition-scalable cryogenic implosions.

The largest contribution to the uncertainty in the yield ratio measurement is caused by the D–D yield. Upcoming experiments are designed to increase the accuracy of this measurement to 5%. These experiments will reduce the uncertainty in the $Y_{DT}/Y_{DD}$ ratio to 7%, which, in turn, will also increase the accuracy of the inferred fuel fractions obtained from this measurement.

Presently, there is no measurement available of the true temperature of the plasma, which is very important for this measurement. Several projects are being considered that will provide a true thermal temperature that is not influenced by the bulk motion of the plasma.

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REFERENCES


Generating strong shocks of up to several hundred megabars makes it possible (1) to explore plasma and material properties at the most-extreme conditions of energy density and (2) to develop two-step inertial confinement fusion (ICF) schemes, where ignition is separated from the main compression of the thermonuclear fuel. A promising two-step ignition scheme is shock ignition (SI), where ignition is triggered by a strong shock launched at the end of the implosion and driven by a pressure above \( \sim 300 \) Mbar. Detailed reviews of the current status and physics issues for SI are found in Refs. 5–7. One of the most-critical issues is that the ignitor spike pulse requires a laser intensity of \( 5 \times 10^{15} \) to \( 1 \times 10^{16} \) W/cm\(^2\), which will excite parametric laser–plasma instabilities (LPI's) in the hot plasma corona surrounding the imploding capsule, thereby transferring a significant amount of the laser energy to the hot electrons. Recent work demonstrated that hot electrons can enhance the shock pressure. It is still an open question whether they might preheat a SI target or if the benefits will prevail because the areal density is large enough to stop them in the shell and augment the shock. Another concern pertains to the energy coupling. The spike pulse must couple sufficient energy into the target in order to generate a strong-enough shock. LPI's may reduce the coupling efficiency and prevent the seed shock pressure from reaching the required magnitude.

Measuring the pressure at these high intensities directly is nearly impossible, so it must be instead inferred indirectly. Experiments in planar geometry at the Laboratoire pour l’Utilisation des Lasers Intenses (LULI), Omega, and Prague Asterix Laser System (PALS) laser facilities have inferred ablation pressures in the range of \( \sim 40 \) to 90 Mbar, which were limited by lateral heat flow from the laser spots in the planar geometries. The lateral transport was suppressed with the development of a new platform that applies spherical targets and x-ray diagnostics. It allows one to evaluate the pressure at shock-ignition–relevant laser intensities. The laser launches an inwardly propagating shock wave that converges at the center, heating a small volume and generating a short x-ray flash that is measured with a time-resolved diagnostic. The shock-launching conditions are inferred by constraining radiation–hydrodynamic simulations to the experimental observables. Several experiments established this scheme as a reliable platform using a variety of laser energies, pulse shapes, and target diameters.

There is a continuing interest in exploring new ablator materials in direct-drive ICF research to improve the hydrodynamic efficiency, mitigate the hot-electron production, and suppress the Rayleigh–Taylor instability. Recent theoretical work demonstrated an overall better performance with mid-Z ablators than plastic (CH) ablators by suppressing the threshold of detrimental LPI while preserving the hydrodynamic stability properties. All of this work has been performed, however, at laser intensities of up to \( \sim 1 \times 10^{15} \) W/cm\(^2\), which is relevant for the standard hot-spot–ignition concept but not for the spike interaction in shock ignition. No work has been performed so far to study how the ablator material affects the spike interaction.

This article describes for the first time the important role that the ablator material plays in the interaction physics at shock-ignition–relevant laser intensities. We discovered that CH ablators produce significantly more hot electrons than the other materials and show that differences in the hot-electron production influence the shock formation. Instantaneous conversion efficiencies (CE’s) of laser energy into hot-electron energy reach \( \sim 13\% \) in CH and \( \sim 4\% \) in C. According to simulations, hot electrons increase the effective maximum ablation pressure by \( \sim 77\% \) in CH and by \( \sim 45\% \) in C. This important finding sheds light on the LPI physics in an intensity and plasma regime that is insufficiently explored and might provide a path to higher-energy-density states in direct-drive geometry.

The experiment used 60 UV (\( \lambda = 351 \) nm) beams from the OMEGA laser with a total energy of 22 to 26 kJ that were focused to an overlapping beam intensity of up to \( \sim 5 \times 10^{15} \) W/cm\(^2\) on the surface of a spherical solid target. The beams were equipped with small-spot phase plates, polarization smoothing, and smoothing by spectral dispersion (SSD). Details on the phase-plate configuration can be found in Ref. 18.
The targets with an outer diameter of 412 to 496 μm consist of an inner CH core that is doped with Ti with an atomic concentration of 5% and an outer ablator layer with a thickness of 20 to 46 μm of a different material [Fig. 149.26(a)]. The outer layer is irradiated with the laser pulse shown in Fig. 149.26(b). A low-power prepulse of ~1-ns duration produces a plasma corona with which the high-power part of the pulse interacts to generate the shock and the hot electrons.

Four different ablator materials (CH, Be, C, and SiO2) with different atomic numbers (Z) were used. Table 149.III summarizes the parameters of the ablators. The shock wave converges in the center, which results in a short burst of x-ray radiation that is detected spatially and temporally resolved with multiple x-ray framing cameras. Each framing camera was absolutely timed through dedicated timing shots19,30 with an accuracy of 30 ps. Time-resolved and time-integrated hard x-ray measurements provide a characterization of the hot-electron population (hot-electron temperature and total energy). Optical backscatter diagnostics measure the amount of absorbed laser energy and the back-reflected laser light.

Figure 149.27(a) shows the measured flash time, which is defined as the occurrence of the x-ray flash relative to the start of the laser pulse, for the different ablators with SSD on (squares) and SSD off (circles) in sequence of increasing Z. The measured flash times were adjusted to account for differences in target size, laser energy, and ablator thickness. One-dimensional (1-D) radiation–hydrodynamic simulations were performed with the code LILAC31 to analyze the dependence of the flash time on these variables for each material using the actual measured mass densities. The flash times were then adjusted for an ablator thickness that results in a constant ablator mass, a laser energy of 24 kJ, and a target outer diameter of 430 μm in order to obtain a valid comparison for the different targets. The data show the general trend of an earlier flash with increasing Z except for CH, which produced the earliest flash. Turning SSD off advances the flash in CH by ~70 ps, while no significant effect is observed in the other materials. Figure 149.27(b) shows the measured time-integrated CE. Plastic stands out by producing by far the most hot electrons with up to ~2 kJ of total hot-electron energy (time-integrated CE ~8%) deposited in the target when SSD was turned off. Nine and seven shots were performed for CH with SSD on and off, respectively, to prove that the observed difference is

Table 149.III: Ablator materials along with the ratio of average mass number and average ionization degree (assuming full ionization), average outer target diameter (OD), average ablator layer thickness, and measured mass density.

<table>
<thead>
<tr>
<th>Ablator</th>
<th>⟨A⟩/⟨Z⟩</th>
<th>⟨OD⟩ (μm)</th>
<th>⟨Thickness⟩ (μm)</th>
<th>Density (g/cm³)</th>
</tr>
</thead>
<tbody>
<tr>
<td>CH</td>
<td>1.86</td>
<td>454</td>
<td>40</td>
<td>1.04±0.01</td>
</tr>
<tr>
<td>Be</td>
<td>2.25</td>
<td>430</td>
<td>20</td>
<td>1.84±0.01</td>
</tr>
<tr>
<td>C</td>
<td>2.00</td>
<td>444</td>
<td>28</td>
<td>1.4±0.4</td>
</tr>
<tr>
<td>SiO2</td>
<td>2.00</td>
<td>433</td>
<td>20</td>
<td>1.75±0.2</td>
</tr>
</tbody>
</table>
not an artifact. If CH is treated as an exception, there is the general trend of a slight increase in hot-electron production with higher $Z$. The inferred hot-electron temperatures lie between 60 and 80 keV and are independent of the ablator and SSD. A high hot-electron fraction corresponds to an earlier flash time, which indicates that hot electrons play a role in the shock formation and augment its strength. The experimental data provide information about the dominant mechanism of hot-electron generation. A clear correlation between hot-electron production and the stimulated Raman scattering (SRS) backscatter signal is observed [Fig. 149.27(c)]. Switching SSD on significantly decreases the SRS signal in all ablators, potentially caused by the suppression of beam filamentation. In contrast, the two-plasmon-decay (TPD) instability, which is the other important hot-electron–generation mechanism, is unaffected by SSD and seems to be far less important than SRS in producing hot electrons. The optical emission generated by electron plasma waves (EPW’s) with half the laser frequency ($\omega/2$) is much weaker than the SRS emission and monotonically increases [Fig. 149.27(d)] with $Z$.

An effective maximum ablation pressure has been inferred (see Fig. 149.28) from simulations. The effect of hot electrons was taken into account by increasing the flux limiter so that the flash time was recovered in the simulations for each ablator material. Although it has been shown in Ref. 17 that the pressure increase from hot electrons may be described by an increased flux limiter, this simplified description does not capture important details such as slowing down, preheat, and local energy deposition. Additional simulations were performed for the CH target that included a detailed hot-electron transport model, which confirmed the pressures shown in Fig. 149.28.

Figure 149.29 shows the inferred time-resolved CE (red) for two shots with CH (solid) and C (dashed). The blue curves represent the corresponding laser pulse shapes. The onset of hot-electron production lags by ~0.2 ns with respect to the

1.2 1.4
with a scale length below
A simulation with CH was compared to one
promotes a high SBS gain during this time. After the
in the implosion),
time-resolved hard x-ray emission in the photon energy range
Time-resolved measurements of the SRS backscattering appear
to be closely correlated with the hot-electron production. The
time-resolved conversion efficiency is based on the measured
time-resolved hard x-ray emission \( n_e \) in the photon energy range
between 50 and 100 keV. It is assumed that the instantaneous
amount of hot electrons is proportional to the instantaneous
hard x-ray emission. The conversion efficiencies reached
13±2% and 4±1% in CH and C, respectively, during the second
half of the high-intensity pulse, while the time-integrated CE
over the whole pulse, including the laser energy when no hot
electrons were generated, yielded 9±1% and 3±1% for these
shots, respectively.

The amount of hot-electron energy coupled into the target
core can be estimated with the technique described in Ref. 34
by using two target types that provide the same corona condi-
tion and therefore the same hot-electron source but different
core conditions. The differences in hard x-ray emission from a
target containing a pure CH core and ablator and the Ti-doped
core with CH ablator were compared. About 25% of the hot-
electron energy was deposited beyond the ablator layer into
the unablated dense target, emphasizing the importance of the
energy transport by hot electrons.

The experiments demonstrated significant differences
between CH and C ablators, indicating that the H species plays
an important role in the LPI. To elucidate the SRS physics, 2-D
decay or collapse of
Figures 149.30(a) and 149.30(b) show the calculated longitudi-
na field strength from EPW as a function of time
and distance along the direction of laser propagation. Distinct
differences in the fields are observed. The electromagnetic
wave excites strong EPW over a large region in CH compared
to C. The wave modes survive longer in CH and couple bet-
ter with thermal electrons because of a larger \( k \) vector. As a
result, more hot electrons are generated. Figures 149.30(c) and
149.30(d) compare the calculated signal level of ion-acoustic
waves (IAW’s), showing a stronger damping in CH compared
to C because of the presence of light H ions. The calculated
CE’s into electrons with kinetic energy exceeding 50 keV from
the PIC simulations were 12% and 2% for CH and C, respect-
ively. A possible explanation is that the SRS saturation level
is controlled by the secondary parametric decay or collapse of
the driven plasma wave. The secondary parametric decay has
been discussed in many papers; the experimental demonstration was reported in Ref. 36. The threshold of the parametric decay is proportional to the IAW damping rate. In the case of high IAW damping (with H), the threshold is higher and the plasma-wave amplitude can grow to a higher level, producing a stronger SRS signal and a larger number of hot electrons. Conversely, for a small IAW damping, the SRS is saturated by the EPW collapse at a lower level, producing large-scale density modulations and fewer hot electrons. It has been shown theoretically for a fixed $T_e$ and density scale length that a high IAW damping rate promotes higher hot-electron generation;20 also, theoretical work that studied the nonlinear saturation of SRS in laser hot spots linked an increased SRS reflectivity with a higher IAW damping rate.37 The observed close correlation between SRS and hot-electron production indicates that IAW damping plays a major role in the CH plasma.

It is expected that the ablator material affects the ablation pressure in various ways. In general, thermal electron-heat conduction is lower in higher-Z materials, and we would expect a reduced mass ablation rate and lower ablation pressure. Based on a simple stationary laser ablation model38 that neglects radiation and hot electrons, the ablation pressure from thermal transport is given by $p_a = \rho_c^{1/3} I_{abs}^{2/3}$, where $\rho_c$ is the critical mass density and $I_{abs}$ is the absorbed laser intensity. Therefore, the ablation pressure $p_a \sim (A/(Z))^{1/3}$ depends only weakly on the ratio of mass number and ionization degree for fixed laser wavelength and fixed $I_{abs}$. The expected increase in $p_a$ from CH to Be is only $\sim 7\%$ and even less with respect to the other materials.19 This experiment demonstrates higher ablation pressures for CH and SiO$_2$, however, indicating that other factors such as hot electrons and potentially radiation transport are more important. Higher-Z materials result in increased collisional absorption and a higher production of x-ray radiation. The radiation impinges deeper into the ablator layer than the thermal electrons, creating a double-ablation front for medium- and high-Z materials.22,39

In conclusion, the experiments demonstrate peculiar differences in hot-electron production in the various ablator materials—especially for CH, which generates the most electrons. PIC simulations using input parameters from radiation–hydrodynamic simulations reproduce the higher hot-electron production in CH. This is likely caused by a stronger damping of IAW’s in the CH plasma because of the presence of light H ions.
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REFERENCES


Introduction
Layered cryogenic DT targets are the baseline approach to achieving ignition in direct-drive inertial confinement fusion (ICF) experiments.\textsuperscript{1,2} Steady progress has been made in experiments with hydrodynamically equivalent,\textsuperscript{3} energy-scaled implosions\textsuperscript{4–9} on OMEGA.\textsuperscript{10}

These implosions are designed to achieve similar peak shell velocities ($v_{imp}$), hot-spot convergence ratios (CR, the ratio of initial ice radius to hot-spot radius), and in-flight aspect ratios (IFAR’s) as ignition designs. The IFAR is defined as the ratio of shell radius to shell thickness, given by the full width at 1/e density of the shell, when the shell has reached 2/3 of its initial radius. Recent direct-drive experiments on OMEGA\textsuperscript{9} achieved record performance parameters that when scaled to the laser energy available at the National Ignition Facility (NIF), would achieve a Lawson parameter $P_x \geq 60\%$ of the value required for ignition,\textsuperscript{11} where $P$ is the hot-spot pressure and $x$ is the confinement time. This scaled $P_x$ is similar to the values achieved in indirect-drive implosions on the NIF.\textsuperscript{12,13}

In these experiments the inferred hot-spot pressure $P$ is $\sim 40\%$ lower than one-dimensional (1-D) simulations,\textsuperscript{9} indicating that the experimental performance is significantly degraded. The current hypothesis to explain this performance degradation for implosions with an adiabat ($\alpha$) $> 3.5$ (ratio of shell pressure to the Fermi pressure) is based on low-mode hydrodynamic instabilities.\textsuperscript{8,9,14} These hydro-instabilities can be seeded by long-wavelength nonuniformities in the initial conditions, like ice-thickness variations,\textsuperscript{5} target offset,\textsuperscript{5,15} and laser-drive nonuniformity in space and time (target placement, beam pointing, power balance, and beam timing). Isolated defects like the target stalk,\textsuperscript{16} debris on the target surface, or short-wavelength structures like target-surface roughness\textsuperscript{17} or laser imprint,\textsuperscript{15,18} especially for low-adiabat implosions ($\alpha < 3.5$), can also seed these instabilities.

The performance of experiments with layered cryogenic DT targets has been measured using nuclear and x-ray self-emission diagnostics.\textsuperscript{8,9} Recent three-dimensional (3-D) hydro simulations\textsuperscript{14} have indicated that the x-ray self-emission images show the influence of long-wavelength nonuniformities on the hot core and do not observe the assembly of the cold shell. Figure 149.31 shows an equatorial density map from 3-D ASTER simulations\textsuperscript{14} (a) at peak neutron production compared to (b) a simulated self-emission image from an orthogonal polar view in the 4- to 8-keV x-ray band at the same time. The comparison between the density map and a simulated x-ray image demonstrates that the shape of the x-ray image does not follow the density distribution in the shell.

X-ray backlighting can be used to observe the flow of the dense and relatively cold shell material in these cryo DT implosions. This technique has been used successfully in both direct-drive room-temperature experiments with gas-filled plastic (CH) targets\textsuperscript{19} and in surrogate indirect-drive\textsuperscript{20} ICF implosion experiments to measure the velocity and uniformity of the imploding shell. Figure 149.31(c) shows a simulated backlit image 50 ps before peak neutron production at CR $= 12$. The image is oriented so that the vertical is along the target offset direction. The image shows the absorption of the dense shell as a white ring and the self-emission of the core, which is seen as a darker central feature. The dominant effect from the offset, which will grow into a 5:1 density perturbation at peak compression, is clearly visible in the image and measurable in the lineout [Fig. 149.31(d)], even at this relatively modest convergence.

Direct-drive cryogenic DT implosions on OMEGA are difficult to radiograph because of the low opacity of the DT shell, the high shell velocity, the small size of the stagnating shell, and the very bright self-emission of the hot core. A shaped crystal imaging system with a Si backlighter driven by short (10- to 20-ps) laser pulses from OMEGA EP\textsuperscript{21} was used to radiograph the OMEGA cryogenic implosions. It has the benefits of a narrow spectral width, high photon throughput, and a backlighter with a short emission time and high brightness. Processes with features below the spatial resolution of the imaging system, like mix, can be detected through the opacity effects from the carbon of the ablator material, which will significantly increase
the absorption of the DT shell  if mixing between the ablator and DT shell occurs.

The following sections (1) present the setup of the experiments, including a description of the narrowband crystal imaging system; (2) describe the experimental results in three subsections: (a) low-order modes, (b) stalk effects, and (c) mix; and (3) present our conclusions.

Experimental Setup

The cryogenic targets used in these experiments had an outer radius of \( \approx 430 \) to \( 480 \) \( \mu \)m. An \( \approx 8 \) to \( 12 \) \( \mu \)m-thick ablator shell of either plastic (CH), deuterated plastic (CD), or CD doped with 0.7% germanium encased a 50- to 75-\( \mu \)m-thick cryogenic DT ice layer [see Fig. 149.32(a)]. All targets were characterized using optical shadowgraphy and showed ice thickness variations of typically \( \leq 1\mu \)m root mean square (rms).

Triple-picket pulses of \( \approx 22 \) to \( 25 \) kJ laser energy were used to irradiate the targets, with smoothing by distributed phase plates (DPP's);\(^{22}\) polarization smoothing (PS) with birefringent wedges;\(^{23}\) two-dimensional (2-D), three-color-cycle, 0.33-THz smoothing by spectral dispersion (SSD);\(^{24,25}\) optimized energy balance (<4% beam-to-beam);\(^{26}\) and optimized beam-to-beam timing of \( \approx 10\) -ps rms (Ref. 14). The targets were placed within \( \approx 10 \) \( \mu \)m of target chamber center.\(^{14}\) The shape of the laser pulse was designed to put the shell on a specific adiabat that ranged from \( \approx 2 \) to \( 4 \) in these experiments. Figure 149.32(b) shows examples of both a lower- and a higher-adiabat pulse at comparable total laser energies. The high-adiabat pulses are shorter and have larger picket energies than the low-adiabat pulses. The total laser energy and the total shell mass determine the peak implosion velocity, which ranged from \( v_{\text{imp}} \approx 2.4 \) to \( 3.7 \times 10^7 \) cm/s. The IFAR ranged from 10 to 20 in these experiments. The IFAR is predominantly controlled by the shell thickness and shell adiabat.

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Figure 149.31
(a) Equatorial distributions of the density from a 3-D radiation–hydrodynamic simulation at peak neutron production taken from Ref. 14. (b) Simulated self-emission image from a polar view in the 4- to 8-keV x-ray band at the same time. The direction of the 20-\( \mu \)m target offset is indicated by an arrow. The thin black line in (b) shows the 17% contour of the maximum x-ray fluence. (c) A simulated backlit image 50 ps before peak neutron production at a convergence ratio (CR) \( \approx 12 \). The image is oriented so that the vertical is along the target offset direction. (d) Vertical lineout through the backlit image.
A shaped Bragg crystal-imaging system was used to obtain radiographs of the imploping targets at various convergence ratios. The OMEGA crystal-imaging radiography system\textsuperscript{21} (see Fig. 149.33) uses a Si backlighter driven by the OMEGA EP laser to backlight implosion targets driven by the 60 beams of the OMEGA laser (not shown for clarity). A quartz crystal, cut along the 10\text{1}1 planes for a 2\textit{d} spacing of 0.6687 nm, was used for the Si He\textsubscript{\textalpha} line at \textasciitilde1.865 keV (0.664 nm). The Bragg angle for this configuration was 83.9\textdegree. The crystal was mounted by direct optical contact on an aspheric glass substrate by INRAD.\textsuperscript{27} The crystal has a major radius of curvature of 500 mm and is placed 267 mm from the implosion target. The image is recorded on a detector located \textasciitilde3.65 m from the target, for a magnification of \textasciitilde15\times. The quartz crystal is rectangular with a size of 25 \times 10 mm, resulting in \textit{f} numbers of \textit{f} = 10 in the horizontal and \textit{f} = 25 in the vertical direction. The spectral bandwidth of the imager is of the order of 10 eV, which matches the typical broadened linewidth of the resonance line from the backlighter driven by a short-pulse laser.

The available solid angle for the backlighter foil is quite limited since the backlighter target must not intercept any of the 60 beams pointed at the implosion target. Because the backlighter laser intensity must be kept as high as possible, the 500-\textmu m-sq backlighter was placed 5 mm from the implosion target. A fast target insertion system (FASTPOS) inserts the backlighter target 100 ms after the shroud that protects the layered cryogenic target from ambient thermal radiation has been removed. FASTPOS also acts as the direct line-of-sight (LOS) block. Two additional collimators are placed on the mounting structure for the FASTPOS to suppress background from Compton scattering and fluorescence from structures in the target chamber. To reduce the impact of the self-emission of the hot core of the cryo DT implosion, an x-ray framing-camera (XRFC) head\textsuperscript{28} is used as a detector. The XRFC head is run with either a single-strip microchannel-plate (MCP) detector, with a 300- to 500-ps-long exposure, or a four-strip MCP with an exposure time of \textasciitilde40 ps, where the backlit image is placed in the center of one of the four strips. The spatial resolution of the XRFC recording system is typically \textasciitilde50 \textmu m (Ref. 29). Experiments with resolution grids show an \textasciitilde15-\textmu m, 10\% to 90\% edge response for the crystal-imaging system. This spatial resolution is adequate for these initial experiments. Work is underway to improve the resolution to \textasciitilde10 \textmu m. The XRFC is triggered by an ultrastable electro-optical trigger system with

\begin{figure}[h]
\centering
\includegraphics[width=\textwidth]{Fig149.32.png}
\caption{(a) The cryogenic DT capsules consist of a thin, 8- to 12-\textmu m-thick CH, CD, or doped-CD ablator filled with several hundred atm of DT gas to create a 60- to 75-\textmu m-thick ice layer at cryogenic temperatures below the triple point of DT (\textasciitilde19 K). (b) The laser drive pulse consists of a series of three pickets to establish the shell adiabat and control shock coalescence and a high-intensity main drive with a total energy of 22 to 25 kJ.}
\end{figure}

\begin{figure}[h]
\centering
\includegraphics[width=\textwidth]{Fig149.33.png}
\caption{Schematic of the spherical-crystal-imager backlighting setup from Ref. 21 (not to scale). The short-pulse laser illuminates a backlighter foil behind the primary target, which is heated by 60 beams from the OMEGA laser (not shown). A direct line-of-sight (LOS) block and a collimator protect the detector [an x-ray framing camera (XRFC)] from background x rays emitted by the backlighter and primary targets.}
\end{figure}
a jitter of ~1.5-ps rms. Experiments using only the backlighter foil showed that the XRFC system has a jitter of <10-ps rms with respect to the arrival of the OMEGA EP laser on the backlighter target. The timing of the OMEGA EP pulse to the OMEGA laser was measured to ~10-ps rms using the neutron temporal diagnostic (P11NTD),30 which is also sensitive to the high-energy x rays produced during the interaction of the OMEGA EP laser with the backlighter target.

Figure 149.34(a) shows the temporal evolution of the implosion from 1-D LILAC simulations close to peak compression compared to the laser pulse shape (blue line) for a typical backlit cryogenic implosion. All LILAC simulations shown in this article include the effects of cross-beam energy transfer (CBET) and use a nonlocal thermal-conduction model.6 The trajectory of the shell radius (peak density: green; 1/e of peak density: black) starts at the ~430-μm outer radius of the target and shows the shell moving toward the center until peak compression at ~3.5 ns. The neutron-production rate (orange) peaks ~40 ps before the calculated areal density (magenta). The exposure time of the XRFC is indicated by the gray-shaded area and the arrival time of the OMEGA EP short-pulse laser by the red vertical line. A time-gated image of a backlit DT cryogenic implosion with an exposure time of ~40 ps is shown in Fig. 149.34(b). The dashed white line indicates the original shell diameter, and the white line at the bottom of the image shows the location of the target stalk. The backlighter emission is shown in the center of the image. It is clipped at the top of the XRFC slit because of a misalignment caused by repeatability issues in the crystal insertion mechanism. The absorption from the compressed shell is seen in the image as a ring-like feature around an emission feature from the central bright core of the implosion.

To measure the absorption in the compressed shell and to quantitatively compare the signal recorded by the crystal imager with simulations, the data must be corrected for the backlighter shape. A simple first-order physical model was constructed21 to describe the shape of the backlighter by assuming a constant brightness source. This source was convolved with a Gaussian point-spread function (PSF), representing the spatial resolution of the imaging system at a 5-mm defocus. The brightness and extent of the source and the width of the PSF were varied to obtain a best fit to the shape of the measured signal outside the area affected by the absorption of the target. These uncertainties associated with correction are taken into account in the errors reported on the measured absorption.

**Experimental Results**

1. Low-Order Modes

Long-wavelength nonuniformity can be seeded in an implosion by a number of processes including nonuniformities in the laser illumination, target placement, and thickness variations in both the ablator and the DT ice layer. To study the impact of these long-wavelength nonuniformities on the assembly of the compressed high-density shell close to stagnation, a series of
experiments were performed with preimposed initial-thickness perturbations in the CH shell.

Figure 149.35 illustrates a shaped target with preimposed initial-thickness perturbations in the CH shell. The amplitude of the variation in shell thickness was 2 to 4 \( \mu \text{m} \) peak to peak. This variation caused an \( \sim 2-\mu \text{m} \)–rms inner ice radius nonuniformity in the layering process. A fiducial glue spot of \( \sim 30-\mu \text{m} \) diameter was used to orient the target horizontally, i.e., perpendicular to the stalk that is mounted vertically in the target chamber. Standard-quality targets with an ablator-thickness nonuniformity of \(<0.1-\mu \text{m} \) rms in all modes and a DT ice layer nonuniformity of \(<1.0-\mu \text{m} \) rms were used in separate experiments to establish a reference.

The radiograph from the reference experiment with a standard-quality target (shot 81590) from Fig. 149.34(b) is shown on an expanded scale in Fig. 149.36(a). The image was recorded at \( \sim 100 \) ps before peak neutron production at a CR of 7, with an exposure time of \( \sim 40 \) ps. The absorption of the backlighter by the compressed shell is seen in the image as a ringlike feature around a central emission feature from the bright core of the implosion. The initial CH ablator thickness of the target was \( \sim 12 \mu \text{m} \), with an outer diameter of \( \sim 890 \mu \text{m} \). The measured nonuniformity of the outer surface was \( 0.24-\mu \text{m} \) rms. The thickness of the DT ice layer was measured at \( \sim 61 \mu \text{m} \) with a \( 0.5-\mu \text{m} \)–rms thickness variation. The target was imploded with a triple-picket pulse of 24-kJ energy at a calculated adiabat of \( \sim 2.5 \), which led to a calculated IFAR = 10. The measured offset from target chamber center at shot time was \(<10 \mu \text{m} \). The recorded yield was 20\% of the 1-D calculations [yield over clean (YOC)] and the measured areal density was \( \sim 80\% \) of the predictions.

Figure 149.36(b) shows the backlighter shape–corrected horizontal lineout compared to \textit{Spect3D} post-processed, 1-D \textit{LILAC} simulations.
To further analyze the radiographs and to obtain quantitative data on the shape of the compressed shell, radial lineouts were taken from the center of the self-emission peak and the radius of peak absorption and its magnitude were evaluated as a function of azimuthal angle (see Fig. 149.37). The contour at 1/e of the peak of the core emission is also determined and plotted in Fig. 149.37(a) for comparison. The errors shown in the graph are estimates of the uncertainty determining the peak absorption location or the 1/e of the emission given the signal/noise on the experimental signal. The radius of peak absorption shows predominantly an \( \ell = 1 \) feature of \( \pm 10\)\( \mu \)m amplitude, with a small extra feature at 180° azimuthal angle, which is associated with the stalk (see *Stalk Effects*, p. 42). Within the errors of the evaluation, the 1/e contour of the core self-emission is observed to be circular. The magnitude of peak absorption shows a small \( \pm 5\% \) peak-to-peak variation as a function of angle.

A radiograph obtained in an experiment using a shaped target with a 4-\( \mu \)m peak-to-peak variation in the CH ablator wall thickness (shot 82717) is shown in Fig. 149.38(a). The image was recorded at a CR = 10, \( \pm 50 \) ps before peak neutron production. The gate time of the XRFC was \( \pm 40 \) ps. Because of drifts in the OMEGA EP beam pointing, the registration between the backlighter emission and the implosion is not as good as it was for shot 81590. Nevertheless, the absorption feature from the compressed shell is clearly visible. Since the image was recorded \( \pm 50 \) ps closer to peak neutron production and at peak x-ray emission, the emission of the central core is brighter than in the shot shown in Fig. 149.36. The target had an outer diameter of \( \pm 960 \) \( \mu \)m with an initial CH ablator thickness of \( \pm 11 \) \( \mu \)m. The measured total variation in the radius of the inner DT ice layer was \( \pm 2-\mu \)m rms and its thickness was \( \pm 63 \) \( \mu \)m. The nonuniformity of the outer surface radius was \( \pm 0.21-\mu \)m rms. The target was irradiated with a triple-picket

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**Figure 149.37**
(a) Radius of peak absorption as a function of angle obtained by evaluating lineouts taken from the center of the self-emission peak. The 1/e radius contour from the self-emission is shown for comparison. (b) Peak absorption as a function of angle.

**Figure 149.38**
(a) Backlit image of cryogenic implosion shot 82717 using a shaped target with a 4-\( \mu \)m variation in peak-to-peak CH shell thickness. (b) Backlighter shape–corrected horizontal lineout compared to *Spect3D* post-processed 1-D *LILAC* simulations.
pulse of 25-kJ energy at a calculated adiabat of ~2.0. The calculated IFAR was 14. The measured offset from target chamber center at shot time was <10 μm. The observed YOC was 8% and the measured areal density was ~40% of the calculated value.

Figure 149.38(b) shows the backlighter shape–corrected horizontal lineout of the radiograph in Fig. 149.38(a) compared to Spect3D post-processed, 1-D LILAC simulations, where the backlighter intensity was adjusted to match the observed self-emission of the core. The simulated lineout does not match the experiment quite as well as it did for the comparison shot 81590. While the shape of the self-emission peak is reproduced quite well, the absorption feature from the compressed shell is significantly underestimated. The experimental lineout shows a significant left–right asymmetry, which is consistent with the initial placement of the target, where the thicker side of the CH ablator is placed on the left side of the image shown in Fig. 149.34(a).

The radius of peak absorption and its magnitude are evaluated again as a function of azimuthal angle, together with the contour at 1/e of the peak of the core emission (see Fig. 149.39). The radius of peak absorption shows a feature of ~20-μm amplitude. Clipping on the XRFC strip caused by the pointing instability of the crystal-insertion mechanism made it impossible to extract data in the stalk region around the 180° azimuthal angle. The 1/e contour of the core self-emission shows a measurable ℓ = 2 variation with an amplitude of ~7 μm. A much larger perturbation in the magnitude of peak absorption as a function of an angle of ~±20% is observed with the shaped shell compared to the reference shell.

The backlit images show that even for the reference implosion without any preimposed nonuniformity, deviations from a spherical shell assembly can be seen. Additionally, the fact that the interfaces between shell and core and at the outside of the shell are significantly steeper in the simulation indicates that there is probably small-scale mixing occurring in the deceleration phase that cannot be spatially resolved with the imager and therefore is visible only in the change of the gradients compared to the 1-D simulations. The radiograph for the reference implosion also shows that the shape of the dense shell where a significant ℓ = 1 perturbation is visible, does not necessarily correspond to the shape of the hot spot, which is seen to be round.

The images from the experiments using targets with preimposed CH ablator thickness variations show much larger perturbations than the reference implosion, both in the radius and magnitude of peak absorptions. The lineouts show significantly more absorption over a larger radius than the post-processed 1-D simulation, indicating more mixing between the ablator CH and the DT ice layer. Even though the targets and laser pulses are quite similar, the small differences in both the adiabat and the IFAR lead to significant differences in the shape of the absorption features as compared to simulations.

2. Stalk Effects

The impact of the target stalk and the glue spot, with which the stalk is attached to the shell, on the implosion symmetry has been observed previously using the crystal-imaging system in an implosion experiment with a mass-equivalent CH target fielded from the cryo target insertion system. At a conver-
gence of 2.5, the image revealed a cusp-like feature in the shell radius at the location of the stalk. Figure 149.40(a) shows the shell radius as a function of azimuthal angle, evaluated at the 50% point on the absorption feature seen in the backlit image of the mass-equivalent CH target (shot 69789). The target had a shell thickness of 24 μm and was irradiated with 23 kJ of laser energy. The evaluation shows a narrow feature of ~25-μm amplitude at the stalk location at 180° azimuthal angle. At the stalk feature, the shell radius is larger than the average shell radius.

The change in direction of the stalk perturbation from being larger than the average radius at CR = 2.5 during the acceleration phase to being smaller than the average radius at CR = 7 during the deceleration phase is expected based on an analysis of multidimensional hydro simulations. During acceleration, the stalk area lags behind the rest of the shell because the extra mass of the glue and the shadowing of the laser drive by the stalk reduce the shell velocity. During deceleration, the extra mass at the stalk location causes it to decelerate more slowly against the growing pressure of the core, allowing it to push farther in compared to the rest of the shell.

3. Imprint and Mix

The images from most cryogenic DT target experiments show significantly more contrast than expected from Spect3D post-processed LILAC simulations, which indicates that carbon from the ablator mixes into the DT ice layer. Figure 149.41 shows a lineout through the image of shot 70535 corrected for the backlighter shape. A 300-ps gate was used in these experiments and was timed to start ~500 ps before the calculated time of peak core emission, according to 1-D LILAC hydrocode simulations. The OMEGA EP short-pulse laser was fired ~100 ps before the end of the gate at a time when the shell assembly was compressed to an inner radius of ~90 μm, which translates to a convergence of ~4, given an inner ice shell radius of ~380 μm. The calculated areal density of the DT at this convergence was ~14 mg/cm² with an adiabat of 2.5. The simulations show an IFAR = 12 for this implosion.

The result from a 1-D LILAC simulation, post-processed with the radiation-transport code Spect3D, is plotted for comparison on the left side of the experimental lineout (green line). The backlighting timing had to be shifted ~50 ps earlier to match the measured size of the absorption feature, indicating that the implosion was slightly delayed compared to the simulations. The timing of the OMEGA EP laser during these experiments was not as well controlled as it was for the shots with the 40-ps-exposure-time framing camera and had a jitter of the order of 20-ps rms. The measured absorption was much higher than the absorption calculated from the simulations. One possible explanation for this discrepancy is Rayleigh–Taylor mixing of carbon from the outer CD shell into the DT ice during the shell acceleration. Adding a small amount of carbon
uniformly into the shell in the Spect3D postprocessor [0.1% C (blue line), 0.2% C (red line)] significantly increases the absorption in the model and brings the simulation much closer to the experimental data, especially in the areas of highest absorption corresponding to the dense shell. In the center of the image, the calculated absorption with carbon mixing is higher than observed. This is probably caused by a small amount of self-emission, which is not fully suppressed by the gating.

To estimate the depth of the mixing of the carbon into the DT, the DT shell was split into five regions of equal thickness in the 1-D LILAC simulations. The two-layer simulation (cyan) shows significantly more absorption but still does not match the experiment. Even the four-layer absorption does not compare as well to the experiment as the fully mixed data, indicating that the carbon is most probably fully mixed throughout the DT shell.

Figure 149.42 shows backlighter shape–corrected lineouts through the radiographs from two additional cryogenic target experiments compared to Spect3D post-processed LILAC simulations. The lineouts show only one side of the implosion because they could not be corrected for the backlighter shape resulting from a significant misalignment of the backlighter.

![Figure 149.41](image1)

Backlighter shape–corrected lineout through the radiograph of a cryogenic target shown in Ref. 21 (black line) compared to a Spect3D post-processed LILAC simulation (colored lines). In the simulation the DT ice was split into five layers and C was uniformly mixed into these layers. The left side of the image shows that simulations with a uniform mix of 0.2% C into the DT match the experimental data (green, blue, and red lines). The right side of the image shows results from simulations where the same mass of C is added to the DT, penetrating into more and more layers (magenta, cyan, orange), showing that at least four layers must be mixed for and adequately matched to the experimental data.

![Figure 149.42](image2)

Backlighter shape–corrected lineouts through the radiographs from two cryogenic target experiments compared to Spect3D post-processed LILAC simulations. Mixing of ablator material is required to match the experimental data for a Ge-doped low-adiabat (α = 2.5) shot (80543) but not for a high-adiabat (α = 4.0) implosion (75372).
The target in shot 80543 had an 8-µm CD shell doped with 0.7% Ge (atomic) and a 50-µm-thick DT ice layer. It was imploled with 25 kJ of laser energy using a pulse that set the calculated adiabat of the shell to 2.5. Preheat from the Ge dopant caused the adiabat to rise to 3.5 at the end of the laser pulse. The IFAR of the shell was calculated to be 20. The radiograph was taken with a 40-ps-wide gate, ~150 ps before peak neutron production at a CR = 5 and a predicted areal density of ~40 mg/cm². Shot 75372 used a target with a 7-µm pure CD shell without any dopant and a 75-µm-thick DT ice layer. It was imploled with 23 kJ of laser energy with a calculated shell adiabat of 4. The calculated IFAR was 20. The radiograph was recorded with a 200-ps XRFC gate, 150 ps before bang time at a CR = 7 and a predicted areal density of ~40 mg/cm².

Mixing of ablator material at a level of ~0.2% is required to match the experimental data for the low-adiabat, Ge-doped shot (80543), similar to the mix observed in the low-adiabat, pure-CD shot (70535). No indication of mixing is observed in the higher-adiabat implosion (75372). In both radiographs, strong self-emission from the core is observed.

The radiography data show that the most important parameter controlling the mix from the CH/CD outer shell into the ice seems to be the adiabat since even a stable, very low IFAR = 10 implosion (70535) shows significant mix throughout the DT quite early in the implosion at the end of the acceleration phase, well before the onset of deceleration of the shell. Two similar IFAR = 20 implosions show a mix threshold in adiabat at around α = 4. The magnitude of the mixing appears to be quite small (~0.2%), which is most likely due to the fact that the DT is starting to be ablated quite early in the implosion. The analysis using five layers for shot 70535 shows that at least the outer 20% of the DT shell gets ablated before the end of the acceleration phase. This ablated DT could serve as a buffer between the CD and the dense DT shell that limits the mix.

Conclusions

X-ray backlighting has been used to radiograph the compressed shell in implosion experiments with layered cryogenic DT targets on OMEGA at convergence ratios from 4 to 10. A shaped-crystal-imaging system with a Si backlighter driven by short laser pulses from OMEGA EP has been set up for this challenging radiography configuration.

The effects of long-wavelength nonuniformities on the shell assembly close to stagnation have been studied in an experiment with preimposed initial thickness perturbations in the CH shell. The radiograph from the reference implosion without any preimposed modulations shows a significant ℓ = 1 perturbation in the shape of the dense shell, which does not match the shape of the hot spot. Additionally, indications of small-scale mixing are observed at the interfaces between ablator, DT shell, and the hot core. The images from targets with preimposed thickness variations show much larger perturbations than the reference implosion, in both the radius and magnitude of peak absorptions and significantly more mixing between the ablator CH and the DT ice layer.

The impact of the target stalk and the glue spot—with which the shell is attached to the stalk—on the implosion symmetry has been observed in both mass-equivalent CH targets and layered DT cryo targets. As expected from simulations, the stalk area lags behind the rest of the shell in the acceleration phase because the extra mass of the glue and the shadow from the stalk reduce the shell velocity and push in farther during the deceleration phase because of the extra mass at the stalk location.

The experimental data show that the most important parameter controlling the mix from the CH/CD outer shell into the ice is the adiabat. A threshold in adiabat at around α = 4 has been observed, where mix is below the detection threshold of 0.02%. The magnitude of the mixing appears to be quite small at ~0.2%, which is most likely caused by the fact that the DT is starting to be ablated quite early in the implosion, thereby serving as a buffer between the CD and the dense DT shell, which could limit the amount of mix.

Future experiments will use this radiography technique to separate the performance degradation from different sources of nonuniformity, such as target offset and laser energy imbalance, and the experimental data will be compared with detailed multidimensional hydrocode calculations. A project has been started that will improve the spatial resolution of the shaped crystal imager and increase the brightness of the backlighter in order to radiograph the implosions at a higher convergence closer to peak neutron production. To illustrate the benefit from higher spatial resolution, radial lineouts from Spect3D post-processed LILAC simulations of cryogenic implosions at a convergence ratio of CR = 16 are shown in Fig. 149.43 using either (a) a measured spatial resolution of ~15 µm or (b) an improved resolution of 8 µm. The green lines show the absorption of the DT and CH shell, the red lines show the self-emission from the core, and the black lines show the combination of both effects. With the lower resolution of ~15 µm, the location of the minimum absorption feature from the DT shell with self-emission, indicated by the black arrows in (a) and (b), is seen at a significantly different radius than the mini-
minimum absorption without self-emission, indicated by the green arrows. This discrepancy is reduced at the higher resolution of 8 \( \mu m \), which will allow one to more-accurately determine the location of the dense DT in the presence of self-emission.

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REFERENCES


An Improved Method for Characterizing Plasma Density Profiles Using Angular Filter Refractometry

Introduction
The measurement of plasma density profiles is important to many areas of high-energy-density (HED) laser–plasma interactions. Quantitative analysis of large HED plasmas has historically been challenging in the range of electron densities near $10^{20}$ to $10^{21}$ cm$^{-3}$. This density range is too low for x-ray probing techniques and too high for most optical techniques. The large integrated phase obtained with optical probes makes it difficult to quantitatively measure the density profile when using typical interferometric techniques. A variety of techniques do exist by which one can attempt to measure this region, but each technique has limiting drawbacks.

A novel diagnostic called angular filter refractometry (AFR) can enable one to characterize plasma density profiles up to densities of $10^{21}$ cm$^{-3}$ by measuring the refraction angle of a probe beam passing through the plasma. The refractive information can be analyzed to characterize the density profile of the plasma. Previously used methods of reducing this experimental data to produce a plasma density profile were cumbersome and at times oversimplified the density profile, both resulting in higher uncertainties.

A new method of analysis has been developed that involves simulating the AFR diagnostic response. A density profile described by seven parameters is used to generate synthetic AFR data, and a quantitative method for defining the degree of similarity between synthetic and observed AFR data provides feedback for subsequent iterations. The synthetic density profile is altered using an intelligent annealing algorithm to iteratively converge upon a solution whose resulting synthetic AFR data closely matches observed AFR data.

This approach has multiple advantages over other methods of analysis: It requires minimal user interface, which eliminates human error that exists from direct manipulation of the observed data. It lends itself to a statistical uncertainty calculation based on $\chi^2$ statistics, allowing one to assess quantitative uncertainties. The resultant density profile, by nature of it being analytic, provides smooth gradients free of noise for scale-length calculations.

In this article, the process by which the synthetic density can be matched to observed AFR data is explained in detail. The different causes of uncertainty specific to this approach and to the diagnostic in general will be described.

Angular Filter Refractometry
The AFR diagnostic is part of the fourth-harmonic ($4\omega$) probe system on LLE’s OMEGA EP laser. The $4\omega$ probe is created from the conversion of a Nd:glass laser pulse to its fourth harmonic ($\lambda_p = 263$ nm) and has a pulse width of 10 ps with 20 mJ of energy. Figure 149.44 shows a conceptual schematic for the AFR diagnostic. The red lines represent the path of the undisturbed probe beam. The beam has a diameter of 3.5 mm and passes through the target chamber center (TCC), where the plasma will be created. The probe is collected at $f/4$ and transported more than 4 m to the diagnostic table.

Figure 149.44
A simplified schematic of the angular filter refractometry (AFR) diagnostic. Unrefracted probe rays (red lines) are blocked by the opaque center of the angular filter. Refracted probe rays (dashed blue lines) hit or miss the filter based on their refraction angles through the plasma. The filter casts shadows corresponding to specific refraction angles that are observed on the image plane. TCC: target chamber center
The TCC plane is image relayed to a charge-coupled–device camera with a resolution of 5 μm over a 5-mm field of view in the object plane.\textsuperscript{8}

An angular filter is placed at the focus of the unrefracted probe beam, or the Fourier plane.\textsuperscript{10} The filter consists of a central opaque dot and oscillating transmissive and opaque rings (Fig. 149.44). The unrefracted probe is stopped by the central dot. In the presence of a plasma, refracted rays (dashed blue lines) will fill a larger area of the angular filter. The opaque regions of the filter block bands of refraction angles, thereby casting shadows in the image plane. Because the angle of refraction of a probe ray is directly proportional to its radial location in the Fourier plane, the shadows on the image plane have contours of constant refraction. To calibrate the specific angular cutoffs, a plano-concave lens was placed at TCC, allowing one to deduce the refraction angle (\(\theta\)) as a function of displacement (\(r\)) from the optical axis in the Fourier plane. For a more-detailed description of the diagnostic, see Ref. 6.

Figure 149.45 shows an example of an AFR image measured from a 250-μm-thick CH target that was ablated by four UV (\(\lambda = 351\) nm) laser beams with a total of 9 kJ of energy in a 2.5-ns square pulse focused to an 800-μm-diam spot that contained 95\% of the energy. The target surface was set to \(y = 0\), and the plasma expanded in the positive \(y\) direction away from the target. The refraction bands produced by the AFR diagnostic show the shape of the expanding plasma plume.

**Analysis**

1. Creating a Synthetic AFR Image

The AFR images are analyzed by comparing them to a synthetic image generated with a model 3-D plasma density \(n_e (x,y,z)\). (Note: All following references to \(n_e\) assume a dependence on \(x, y, z\).) A typical HED laser-plasma plume from a planar target can be modeled by a super-Gaussian parallel to the target and exponential normal to the target.\textsuperscript{11} The 3-D density profile was assumed to be axisymmetric along the target normal. The behavior along the target normal at the plasma’s center is modeled as

\[
n_e (0,y,0) = n_0 \left[ A \cdot \exp \left( \frac{-y}{L_{y1}} \right) + (1-A) \cdot \exp \left( \frac{-y}{L_{y2}} \right) \right],
\]

(1)

where \(n_0\) is the peak density and \(A\) assigns relative strength to two exponential profiles with scale lengths \(L_{y1}\) and \(L_{y2}\). This allows the profile to adapt to a decreasing scale length as typically occurs close to the target surface. The full density profile, including the transverse dimension, is

\[
n_e (x,y,z) = n_e (0,y,0) \times \exp \left\{ -\left( \frac{x^2 + z^2}{L_{xz}^2} \right) \right\}.
\]

(2)

where \(L_{xz}\) is the scale length in both \(x\) and \(z\) and the parameters \(c_1, c_2,\) and \(c_3\) (representing two independent parameters) are used to define the order term for the super-Gaussian profile. The order term was empirically found so one could accurately match experimentally observed AFR contours. Equations (1) and (2) together form the seven-parameter function that constitutes the synthetic 3-D density. Figure 149.46(a) shows the density solution to the experimental AFR map shown in Fig. 149.45, where the deduced density spans two orders of magnitude (\(10^{19}\) to \(10^{21}\) cm\(^{-3}\)).

With the probe propagating in the \(z\) direction, the accumulated phase of the probe ray passing through the plasma is related to the 3-D plasma density according to
A n Improved Method for Characterizing Plasma Density Profiles Using Angular Filter Refractometry

\[ \phi(x,y) \approx \frac{\pi}{\lambda_p \lambda_c} \int_{-\infty}^{+\infty} n_e(x,y,z) \, dz \text{[rad]}, \]

where \( \lambda_p \) is the probe-laser wavelength (263 nm) and \( n_e < n_c \), where \( n_c = 1.1 \times 10^{21} / \lambda_p \lambda_{\mu m} = 1.6 \times 10^{22} \text{ cm}^{-3} \) is the critical plasma density for the probe laser. Figure 149.46(b) shows the integrated phase of the density profile in Fig. 149.46(a), where changes in \( x \) and \( y \) along the ray path are ignored. The angle of refraction of a probe ray exiting the plasma is calculated from the transverse gradient of the accrued phase:

\[ \theta(x,y) = \frac{\lambda_p}{2\pi} \left[ \frac{\partial \phi}{\partial x} \right]^2 + \left[ \frac{\partial \phi}{\partial y} \right]^2 \text{[rad]}. \]

From the calculated 2-D refraction map, a filter function based on the calibration is applied that creates a synthetic AFR image comparable to that measured in the experiment [Fig. 149.46(c)].

2. Simulated Annealing

An iterative solver alters the parameters of the synthetic density function to optimize the match between the synthetic and experimental AFR images. The quality of the match is based on the location of the edges of the bands. This was accomplished by taking many radial lineouts of the synthetic and experimental AFR images and finding the position of the edges of the bands at each angle [Fig. 149.47(a)]. The edge locations in the experimental images were found by applying a 20% intensity threshold to the normalized data, which eliminated most noise without notable alterations to the bands’ behavior and size. Figure 149.47(b) compares the thresholded experimental edges to the synthetic AFR edges. The squared differences of the locations between the synthetic and experimental AFR band edges were averaged over the entire image:

\[ m = \sum_{n=1}^{14} \sum_{r=1}^{100} (E_{n,r} - S_{n,r})^2, \]

Figure 149.46
(a) A 2-D slice of a 3-D axisymmetric synthetic density profile generated from the density function described by Eq. (2). (b) The integrated phase accrued by the probe passing through the synthetic plasma. (c) The synthetic AFR image made by extracting refraction information from the phase map following Eq. (4) and eliminating angles blocked by the filter. (d) A visual overlay of the synthetic (orange) image and experimental (red) image. Blue indicates where the profiles overlap.

Figure 149.47
(a) The algorithm draws lineouts from the center of the target surface and finds the edges of the experimental and synthetic bands on those lines. (b) A lineout of the sixth band (counting from the bottom) taken along \( x = 0 \). A threshold was applied to the experimental data to eliminate noise.
where $E_{n,r}$ and $S_{n,r}$ are the edge locations of the experimental and synthetic bands, respectively, at an edge $n$ and an angle $r$. All 14 edges were examined at a large number of angles so that slight fluctuations at some points in the experimental AFR image do not dominate the match (Fig. 149.47).

The solver incorporates the simulated annealing (SA) algorithm to systematically alter the variables of the density function until a global minimum for $m$ in the parameter space is found. An SA algorithm was chosen because of its ability to find a global minimum in a large parameter space [seven parameters; see Eq. (1)], where many local minima exist. Figure 149.48 displays the logical flowchart that the SA algorithm follows. The SA algorithm takes an initial user-defined density guess $n_i$, generates the synthetic AFR image, and calculates $m_i$. The density is then perturbed by $\Delta n$ and a new match profile $m_{i+\Delta n}$ is calculated. If $m_{i+\Delta n} - m_i < 0$, the new profile is accepted and $n_i + \Delta n$ becomes the new $n_i$. If $m_{i+\Delta n} - m_i > 0$, the new profile is considered for rejection, although there is a chance that it may be accepted.

Accepting a poorer match over a superior one allows the profile to escape from local minima enroute to the final solution. The range over which parameters’ values are generated and the likelihood of accepting a poorer match decrease at later iterations so that the algorithm focuses around a nearby solution. After a set number of runs, the SA algorithm resets the range of parameters in the search and the likelihood of accepting a poor match. Repeating this process numerous times makes it unlikely for the algorithm to get stuck in a local minimum. The simulated annealing algorithm terminates when a chosen number of iterations pass without a new best match being discovered (stop criterion). At this point the profile corresponding to the lowest calculated match is returned by the algorithm. Figure 149.49(a) shows how a single simulating annealing algorithm closes in on its results. Occasionally it escapes its local minimum and finds a new one, trending toward the optimal value. The process of escaping a local minimum can be seen more clearly in Fig. 149.49(b).

![Flowchart of simulated annealing](image)

**Figure 149.48**
A flowchart describing simulated annealing.
represents the degree of noise for a given experimental AFR image, so $\sigma^2 = m_{\text{min}}$. The uncertainty in each parameter is related to the way it alters the behavior of $\chi^2$ around $\chi^2_{\text{min}}$. Altering the parameters around their best-fit values increases $\chi^2$, indicating a lower probability that these parameter values are correct. A confidence interval $\Delta S$ is defined to describe the increase to $\chi^2_{\text{min}}$ that would result in an $N$-percent certainty that the solution lies within $\Delta S$ (Ref. 5). Each parameter is individually altered until $\chi^2 = \chi^2_{\text{min}} + \Delta S$; the boundaries of this window represent the uncertainty in the parameter, $\sigma$. This was factored back into Eq. (6) to find the uncertainty in density. Because $\chi^2$ is inversely proportional to $m_{\text{min}}$, $\Delta n_{\text{stat}}$ will be larger for profiles whose best matches are not as strong. The calculated uncertainty map from $\chi^2$ statistics for the case in Fig. 149.46 is shown in Fig. 149.50(a).

### Figure 149.50

(a) The statistical uncertainty map corresponding to the density function.
(b) The degenerative uncertainty. Note that the color bars cut off at 30% but at one point are as high as 100%. (c) The left–right uncertainty. Note that the color bars cut off at 30%, but uncertainties at the outer regions get higher. (d) The combined uncertainty.

#### 2. Degenerative Uncertainty

The next source of uncertainty is related to the fact that the AFR diagnostic measures refraction, which is proportional to the gradient of the plasma density, or phase. The phase is proportional to the integration of refraction plus an integration constant $c$ allowing for degenerate solutions. The value of $c$ is a source of uncertainty in the density since changing $c$ does not change the AFR image; therefore, boundary conditions must be established.
The main physical boundary condition on the density function is that density must fall to zero away from the target surface. The lowest density that contributes to the AFR image lies somewhere along the outer band’s edge. This value must be non-negative, which gives a lower bound for \( c \) (negative value). Positive values for \( c \) violate the boundary condition of density going to zero without the introduction of additional gradients that exist outside the outermost band that are smaller than measurable by the AFR diagnostic.

There is no way to define the upper bound for \( c \), so for testing purposes, the maximum shift to the density in either direction was taken to be the largest downward shift possible. A density function with over 20 parameters was used to create test AFR images. Those AFR images were treated as experimental ones and run through the iterative solver. The statistical uncertainty and degenerative uncertainty combined were always able to encompass the percent difference between the test cases and the corresponding optimized synthetic densities, proving the validity of these uncertainty calculations.

The corresponding uncertainty can be described by [Fig. 149.50(b)]

\[
\Delta n_{\text{deg}} = \pm \frac{c}{n_e} \%.
\]  

3. Asymmetry Uncertainty

There is a consistent left–right asymmetry in all AFR images. In theory the plasmas should be approximately axisymmetric due to nearly axisymmetric illumination, so it is believed that this asymmetry is symptomatic of an aberration in the probe beam. Efforts to model the presumed aberration were unsuccessful; therefore, it is accounted for as a source of uncertainty. Optimizations are run separately on the left and right sides of each shot and the solution is taken to exist somewhere within the percent difference between the resulting densities. This percent difference contributes to the uncertainty [Fig. 149.50(c)]:

\[
\Delta n_{\text{lr}} = \left| \frac{n_{\text{left}} - n_{\text{right}}}{1/2 \times (n_{\text{left}} + n_{\text{right}})} \right| \%.
\]

Discussion

The total uncertainty is generated by adding the three separate uncertainty sources in quadrature.\(^{16}\) The uncertainty calculations for the far left and right edges exceed 100% because of the asymmetry. Over a large region of interest, the central two-thirds of the profile has an uncertainty of under 20%. The uncertainty along the target normal is under 10%.

The use of an analytic density function is an additional benefit to this analysis method. It facilitates an accurate calculation of the density scale lengths caused by the smoothness of the density derivative. This results in a low uncertainty for scale length. The scale length can be calculated as

\[
L_y = \pm n_e \left[ \frac{dn_e}{dy} \right]^{-1}.
\]

Figure 149.51 shows the density and scale length of the plasma analyzed throughout this article along the target normal. Note that the uncertainty increases as the plasma is farther from the target surface but does not exceed 10%.

![Figure 149.51](image)

The blue line is the plasma density profile along the target normal at the center of the plasma profile (\( x = 0 \)) measured from the AFR data shown in Fig. 149.45. The original target surface is located at \( y = 0 \). The green curve is the corresponding scale length. The uncertainty in scale length increases with \( y \) but is under 10% at all points.

Conclusion

A new method of analyzing data from the AFR diagnostic has been developed. A seven-parameter density profile was used to produce synthetic AFR images, and an iterative solver was developed that could successfully match synthetic data to experimental AFR images. A 2-D uncertainty map for the 3-D density was presented that has an uncertainty of less than 10% in the region of interest.
Several future improvements could increase the accuracy of this analysis. By adding more variables to the density function, it will have more flexibility to match the experimental AFR images, therefore improving the model fit. This would, however, be gained at the cost of computer run time. The degeneracy uncertainty could be erased completely if a boundary condition was known. This could be accomplished, for example, by measuring phase in the low-density regions using simultaneous interferometry. If the asymmetry was caused by an aberration, it could be largely reduced or eliminated by successful modeling the aberration on the probe beam, or experimentally fixing the aberration.

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REFERENCES


Laser-Driven Magnetized Liner Inertial Fusion on OMEGA

Introduction
Magnetized liner inertial fusion (MagLIF) is a concept that utilizes pulsed-power–driven Z pinches of metal liners to compress deuterium–tritium (DT) gas to fusion-relevant temperatures and pressures. For cylindrical compression and pulsed-power time scales (~100 ns), it is required that the fuel be preheated to ~100 eV and axially magnetized to suppress radial conduction losses to achieve near-adiabatic compression, reducing the radial convergence required to reach the temperatures and pressures needed for thermonuclear fusion. At stagnation, the axial magnetic field is compressed to the point where it is strong enough to magnetize alpha particles, allowing self-heating to occur at low areal densities.

Cylindrical implosions can be achieved with 40 beams of the OMEGA Laser System, and in fact, magnetized cylindrical implosions have been carried out on OMEGA, leaving only the laser preheating of the gas. A single beam has been redirected down a symmetry axis of OMEGA to heat the gas prior to compression.

MagLIF is being scaled down from a pulsed-power–driven device to a laser-driven device for several reasons: pulsed-power devices, like the Z machine at Sandia National Laboratories, are very violent environments in terms of debris and electromagnetic noise, making it very difficult to field diagnostics and maintain a high shot rate. Furthermore, diagnostic access around the target chamber in Z is limited by the installation of the axial magnetic field coils and the geometry of the current delivery system. OMEGA can perform roughly 10^7 more shots per day than Z and can provide better statistics and wider scans of the MagLIF parameter space. Furthermore, OMEGA has the capability to perform measurements that cannot be done on the Z machine, such as proton radiography of the compressed axial magnetic field, low-yield neutron measurements, and time-resolved x-ray measurements of the liner trajectory. Experiments at the OMEGA scale can provide another experimental data point for the energy scaling of the MagLIF concept and will ultimately give us the confidence in extrapolating MagLIF to ignition-scale designs.

Point Design
The OMEGA point design^5 is energy scaled from the Z machine’s 29-MA point design. The conserved quantity is energy per unit volume, which mandates that the linear dimensions be scaled down by a factor of 10 to match the factor-of-1000 difference in driver energy between OMEGA and Z. Other factors such as fuel preheat and initial axial magnetic field can be the same. A higher implosion velocity will be needed at the OMEGA scale to compensate for increased thermal losses at the smaller scale. The preheat temperature, liner aspect ratio, and fuel density can be changed to achieve different implosion energetics for a complete scan of the MagLIF parameter space. An ensemble of 1-D MHD simulations that include electrothermal terms in Ohm’s Law^6 was used to determine the optimal laser pulse length, taking into account the drop in on-target energy for pulses longer than 1 ns, and fuel density for shell thicknesses from 20 μm to 50 μm for a fixed 10-T initial axial magnetic field and 200-eV preheat temperature, the objective being to maximize neutron yield at a fuel convergence ratio close to the 25 chosen for the Z point design. Only a square-shaped laser pulse was considered. The optimal design for a 30-μm shell is a 1.5-ns pulse length with an initial fuel density >1.5 mg/cm^3 as shown in Fig. 150.1. Thicker shells did not give adequate final fuel temperatures.

This optimal point is for a fixed magnetic field and preheat temperature, which is easily achievable. If the magnetic-field capabilities of OMEGA are expanded to values above 10 T,
the optimal point may change. Increasing magnetic field and preheat reduces convergence ratio and implosion speed, providing a more stable cylindrical implosion. Higher core pressures are achieved for higher magnetic fields because of the suppression of radial conduction losses (seen in Fig. 150.2). A higher preheat temperature leads to a lower final pressure for a fixed energy implosion, which is consistent with a simple model for adiabatic compression. For a given energy in a piston $E$ and initial pressure $P_0$ and volume $V_0$, the final pressure increase is

$$\frac{P_f}{P_0} = \left( \frac{E}{P_0 V_0} \right)^{\gamma/\gamma-1} \quad (1)$$

and is therefore inversely proportional to the starting temperature for a fixed initial mass piston. Substituting this back into the energy balance equation and solving for the convergence ratio, we obtain

$$\text{CR} = \frac{R_0}{\sqrt{E/P_0 \pi}}, \quad (2)$$

which is proportional to the final pressure or inversely proportional to the initial temperature. Higher starting temperatures, therefore, give a lower final pressure and lower convergence ratio, which is the trend highlighted in Fig. 150.2.

Figure 150.1
(a) The $D_2$ fuel convergence ratio as a function of initial gas density for three different pulse lengths; (b) the neutron yields from each of these designs. These plots show that a 1.5-ns pulse is optimal and that the design requires an initial fuel density higher than that which optimizes neutron yield to maintain a fuel convergence ratio <30, indicated by the red dashed line.

Figure 150.2
(a) As the magnetic field increases, the volume-averaged thermal pressure of the fuel increases, resulting in (b) a lower convergence ratio at the end of the implosion. This is mostly caused by the magnetic field suppressing radial conduction losses. The red circled region is the point design considering the capabilities of the OMEGA Laser System.
Regardless of the starting magnetic field or shell thickness, 1-D calculations show that a minimum preheat temperature of ~100 eV is required for neutron yield increases larger than a factor of 2 from the magnetic field above the implosion-only baseline of ~10^{10} mm^{-1}. Once above the threshold preheat, neutron yields and ion temperatures do not increase with initial temperature, but convergence ratio decreases, increasing the stability of the imploding shell. With a sufficient preheat temperature, increasing the initial magnetic field from 10 T to 30 T increases the neutron yield as shown in Fig. 150.3. Above 30 T, heat loss is ion diffusion dominated since radial electron conduction is essentially zero. Therefore, there is no further benefit from increasing the initial field. The magnetic field required to suppress ion heat flow introduces too much magnetic-field pressure, making it difficult to compress the target.

**Preheat Experiments**

The focused preheat experiments determined the laser entrance hole (LEH) window transmission, backscatter, and sidescatter; a gas-filled cylindrical target was illuminated with up to 200 J of 3ω light in a 2.5-ns square-shaped pulse to study gas preheating in situ. The beam was focused on the LEH window and a 200-μm phase plate with smoothing by spectral dispersion and distributed polarization rotators were used. To study the window behavior in detail, window-only assemblies that consisted of the same 1.84-μm-thick polyimide foils used for the gas cylinder targets were studied. Using calorimeters, Raman and Brillouin spectrometers in different ports around the laser axis, soft x-ray emission from the LEH window and the gas, and optical emission from the surface of the gas-filled cylinder, we characterized the LEH window disassembly and the energy that propagates into the gas and determined a lower bound on the preheat temperature.

From the calorimeter measurements and backscatter diagnostics of the window-only shots, we determined that 64.5±2.0% of the laser energy incident on the LEH window is transmitted, with only 0.72±0.22% scattered outside of a 28° cone and 0.59±0.16% backscattered. It should be noted that the backscatter measurements herein are Brillouin measurements since the Raman measurement was below detectable threshold. We can then infer that 34±2.0% of the laser energy is absorbed in the window material as it disassembles. We have calculated the absorbed fraction using the 2-D hydrocodes DRACO and FLASH, both of which give an absorbed energy of ~30%. Furthermore, we can post-process the output from these hydrocodes to model the soft x-ray spectra of the LEH window disassembly. The results of this spectral analysis compared directly to measurements from an array of differentially filtered x-ray diodes are in good agreement, as illustrated in Fig. 150.4.
Analysis of soft x rays from a diagnostic side window in the gas-filled cylindrical targets infer that a minimum possible gas temperature of 100 eV was achieved 1.3 ns into the laser pulse. This minimum value is determined from one shot, with other shots showing solutions above this minimum temperature. A parylene-AF4 fluorinated plastic cylinder was filled with 2-at. % neon-doped deuterium gas. The gas was then heated using the same laser beam that illuminated the LEH window assemblies. A side-on diagnostic window was imaged using a differentially filtered three-channel soft x-ray imager (SXR). Since the SXR is not absolutely calibrated and we have limited spectral information, the ratios of the spatially integrated channel signals are compared to a grid of possible temperatures and densities for the gas and the wall generated by a simulation. Comparing the channel ratios with this grid gives an infinite set of solutions, but if we constrain the solution space by insisting that the wall temperature cannot exceed the gas temperature, we establish the lowest possible value of the gas temperature to be 100 eV (shown in Fig. 150.5). Unfortunately, because of the limited dynamic range of the SXR and the quick increase in emission from increasing $T_{\text{gas}}$, temperatures above this 100-eV lower limit cannot be determined. We also obtain information about the gas heating from the x-ray diode array by looking at the LEH region. Much of the data is heavily encoded because of spectral integration, so we will rely primarily on comparison with hydrocodes to get a good idea of the heating process of the gas. A more-detailed paper on this experiment is expected to be submitted to Physics of Plasmas in the near future.

### Implosion Experiments

Implosion-only experiments were used to optimize the beam pointing and balance between normal and oblique beams; normal beams, referring to two rings of ten beams at an incidence angle of ±9°, and oblique beams, referring to two rings of ten beams at an incidence angle of ±31°. Both the separation and the intensity of the beams determine the uniformity and length of the cylindrical implosions. Using time-resolved x-ray images of the shell in flight, a shape can be determined by fitting the inner surface with a fourth-order polynomial function as shown in Fig. 150.6:

$$R(z) = a + b(z-z_0)^2 + c(z-z_0)^4.$$  \hspace{1cm} (3)

The coefficients of this polynomial give a numerical measure that indicates if the shell has been overdriven at the ends or middle or uniformly imploded. Lineouts from time-integrated x-ray pinhole camera images also show the uniformity of the core and the length of the imploded region (as seen in Fig. 150.7). The illumination pattern that gives the most-uniform implosion empirically is an overlap of the oblique-angled beams at the center with the normal beams at the end and a reduction in energy of the normal beams to 83% of the maximum energy of the oblique beams. The result is the relatively uniform axial intensity profile shown in Fig. 150.8.

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**Figure 150.5**

The ratio of Channel 2 to Channel 3 of the soft x-ray imager versus the wall temperature of the cylinder shows that the lower bound solution is 100 eV. Other solutions from different channel ratios either violate the condition of $T_{\text{gas}} > T_{\text{wall}}$ or give higher values.

**Figure 150.6**

Each side of the shell was fit with a fourth-order polynomial to determine the shape of the shell in flight. The laser energy was then tuned to get the shape as flat as possible over the longest region.
Yield enhancement from both preheat and magnetic field and preheat only matched with 1-D and 2-D simulation predictions based on the point design. We have many shots with just the implosion from the beam-balancing campaign. In Fig. 150.10, the three best-quality implosion-only targets are shown along with a preheat and implosion shot, the two successful integrated MagLIF shots from the preheat beam-timing campaign, and the predicted performance of the point design from 2-D HYDRA MHD simulations. An implosion with magnetic field and no preheat has yet to be successfully completed. One- and two-dimensional MHD simulations replicating the implosion dynamics and magnetic-field compression are also under development. This is the first demonstration of yield enhancement in a magnetized cylindrical implosion.

Integrated MagLIF Experiments

The first integrated MagLIF experiments on OMEGA were used to scan preheat beam timing relative to the drive beams. Simulations and experiments both indicate the optimal time to fire the preheat laser was \( \sim 1.0 \) ns before the start of the drive beams, which corresponds to preheat finishing as the shell starts to implode. This made it possible for preheat to occur without introducing too much mix of wall material into the gas. Three times were scanned and the results can be seen in Fig. 150.9.

Figure 150.7
Thinner shells that have faster implosion velocities give broader and flatter emission from the core. A raw lineout of charge-injection device counts plotted versus the axial position from a pinhole camera demonstrates this fact. This suggests that a faster implosion is needed at the OMEGA scale to mitigate increased thermal losses. Pinhole images also provide a second metric for implosion uniformity.

Figure 150.8
The irradiation pattern that gives the most cylindrically uniform compression.

Figure 150.9
Neutron yield as a function of preheat beam time. The start of the drive beams is at \( t = 0 \). Some implosion beams were partially obstructed by the magnetic-field coils for both cases at the \(-1.3\)-ns timing, so these shots must be repeated, but the data follow the expected trend predicted by the simulations.

Figure 150.10
A summary of neutron yields for each configuration of the OMEGA MagLIF platform. Integrated MagLIF shots are compared to a 2-D magnetohydrodynamic (MHD) prediction made by the code HYDRA.
on OMEGA and is a very promising result that confirms the utility of laser-driven MagLIF in the development of a viable path toward ignition.

**Future Work**

At the time of publication of this article, the first experiments that will probe the evolution of the axial magnetic field will have been performed. Understanding the dynamics of the magnetic field within the fuel is directly related to the confidence interval of yield predictions from MHD simulations. If the magnetic-field advection within the gas is poorly understood, yields can vary in simulation by whole orders of magnitude. The dominant contribution to magnetic-field advection within the gas is from the Nernst effect, which is an additional advection velocity to the magnetic field proportional to the electron heat flow. Therefore, the heat flow itself can push against the magnetic field, thereby negating any benefit of the reduced thermal conduction. Proton radiography of the implosion can provide a direct indication of the rate of this additional advection.

Experiments to explore the MagLIF parameter space are scheduled to occur over the next year. A magnetic-field scan will explore the dependence on the magnetic field and help understand the scaling of the Nernst effect with the magnetic field. A scan of the initial fuel density will determine the highest achievable convergence ratio before a decrease in performance. Laser-driven MagLIF has the ability to use thinner shells with higher implosion velocities than pulsed-power-driven MagLIF because of ablative stabilization of the Rayleigh–Taylor instability. A scan will be made to determine the minimum shell thickness that can be imploded without a significant decrease in neutron yield.

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**REFERENCES**


Mitigation of Cross-Beam Energy Transfer in Symmetric Implosions on OMEGA Using Wavelength Detuning

Introduction
In the direct-drive\textsuperscript{1,2} approach to inertial confinement fusion (ICF), laser beams directly illuminate a spherical target, depositing most of their energy in the coronal plasma. This energy is transported by electron thermal conduction through a conduction zone to higher densities, where ablation occurs. At the ablation surface, material rapidly expands, producing pressure that drives the shell of the target and thermonuclear fuel \[\text{usually deuterium–tritium (DT)}\] toward the center of the capsule, compressing the target to \[\sim 400 \text{g/cm}^3\] (Ref. 2). To achieve this compression, the laser pulses are precisely shaped to launch a series of synchronized spherical shocks that cause the fuel to compress quasi-adiabatically. As the capsule converges, its kinetic energy is converted to internal energy, creating a hot dense core in which fusion reactions initiate, surrounded by a cold, dense, nearly Fermi-degenerate shell.\textsuperscript{1–3}

Successful direct-drive ignition requires both efficient deposition of the laser energy in the coronal plasma and uniform target illumination to produce the spherically symmetric drive required to avoid hydrodynamic instabilities and low-mode-number asymmetries that can quench the implosion.\textsuperscript{4,5} The target is illuminated by a number of beams, distributed symmetrically around the target, with diameters that are selected by the trade-off between increased drive uniformity and decreased drive efficiency as the laser spot size increases.\textsuperscript{6} When neglecting laser–plasma instabilities, a laser focal-spot radius approximately equal to the target radius provides the best compromise.\textsuperscript{7}

The total laser drive pressure and its uniformity can be significantly degraded by the transfer of energy between laser beams crossing in the coronal plasma.\textsuperscript{8–12} Cross-beam energy transfer (CBET) is a three-wave process that occurs when the beat wave created by the interference between two electromagnetic waves resonantly excites a plasma ion-acoustic wave (IAW) as shown in Fig. 150.11. When two lasers with frequencies \(\omega_1, \omega_2\) and wave vectors \(\vec{k}_1, \vec{k}_2\) cross in a plasma, the ponderomotive force of their beat wave can drive a plasma density perturbation. These density perturbations form a grating and cause Bragg diffraction, facilitating the transfer via stimulated Brillouin scattering (SBS). The coupling is maximized when the driven wave satisfies the dispersion relation for the IAW:

\[
\omega_{\text{IAW}} - \vec{u}_f \cdot \vec{k}_{\text{IAW}} = \pm c_s \left| k_{\text{IAW}} \right|, \\
\omega_{\text{IAW}} = \omega_1 - \omega_2, \tag{1}
\]

where \(\vec{u}_f\) is the local plasma hydrodynamic flow velocity, \(c_s\) is the local sound speed, and \(\omega_{\text{IAW}}\) and \(k_{\text{IAW}}\) are the frequency and wave vector, respectively, of the ion-acoustic wave. The two branches \(\pm c_s \left| k_{\text{IAW}} \right|\) of the dispersion relation correspond to the direction of power flow from the higher-frequency (in the plasma reference frame) “pump” wave to the lower-frequency “seed” wave. Since CBET is seeded by a laser beam rather than small-amplitude thermal noise, significant energy can be exchanged even when the SBS gain is small.

![Figure 150.11](E26609HR)

A \(k\)-space diagram of cross-beam energy transfer (CBET). Energy is transferred from the pump beam to the seed beam as indicated by the magenta arrow labeled “Energy transfer.”
Experiments have demonstrated the existence of CBET between frequency-mismatched beams and beams with the same frequency but crossing in a flowing plasma. CBET has been modeled many times for a pair of crossing beams. Indirect-drive hohlraum experiments at the National Ignition Facility (NIF) have identified CBET as a mechanism responsible for transferring significant amounts of energy between laser beams. In these experiments, the angle between crossing beams was small enough that steady-state CBET models could use a 3-D paraxial approximation or neglect small 3-D effects.

These models showed that significant CBET occurred between NIF beams when they were at the same wavelength and that energy was forward scattered from beams pointed toward the hohlraum equator to those directed nearer to the ends of the hohlraum, affecting the implosion symmetry on indirect-drive hohlraum experiments. It was also shown that CBET can distort the effective beam profile even when the net transfer between beams is zero. These models predicted that relatively small wavelength shifts of the order of 1 Å could tune the shape of an indirect-drive hohlraum implosion by transferring energy between beam rings. Independently varying the wavelength of the NIF beams to control CBET is now used as a tool to tune the implosion symmetry on the NIF and to reduce stimulated Raman scattering (SRS) backscatter. Recently, an in-line CBET model was incorporated into the main 3-D radiation–hydrodynamics code, known as HYDRA, for the NIF.

In direct drive, the presence of CBET was first inferred from the experimental observation of the scattered-light spectra and the implosion velocity. Early direct-drive CBET modeling typically used a 1-D linear geometry, however, to properly model a direct-drive implosion, the crossings of many beams must be calculated simultaneously. The complex beam paths caused by refraction through the coronal plasma invalidate the paraxial approximation, and CBET models for direct-drive implosions typically use 3-D ray tracing to calculate the crossing beam trajectories. Initial CBET modeling suggested that in direct drive, CBET could backscatter energy out of ingoing rays from the hydrodynamically efficient small-impact parameter inner portion of the laser beam spot to outgoing large-impact parameter rays near the edge of the beam spot. This would allow significant amounts of the incident energy from the central portion of the laser beams to bypass the highest absorption region near the critical surface, reducing the ablation pressure and hydrodynamic efficiency of the implosion. The redistribution of power modifies the effective beam profile identically for each beam in a symmetric implosion and can have a large effect on a target’s illumination uniformity. The details and orientation of the redistribution depend on the 3-D positions of the beams with respect to each other and should be modeled in 3-D. In-line ray-based CBET models have now been added to the direct-drive codes LILAC (1-D) and DRACO (2-D), which allow one to hydrodynamically self-consistently model CBET in direct-drive implosions. CBET redistributes ~30% of the incident energy on OMEGA at intensities of $5 \times 10^{14}$ W/cm² and is responsible for a 10% to 20% reduction in laser absorption according to the LILAC model.

Several different schemes have been proposed to mitigate CBET in direct-drive implosions, including doped ablators, narrow beams, and multicolor lasers. The predictions in direct drive that outgoing light from the edge of the beam was taking energy out of the incoming light from the beam center led to the proposal that shrinking the beam radius would reduce CBET. Studies have shown that reducing the diameter of the laser beams by 30% can restore nearly all of the kinetic energy lost to CBET, but at a cost of increased low-mode perturbations. Low-mode uniformity might be maintained by using two-state beam “zooming,” where the implosion is initiated using full-sized beams that are then reduced in radius after the corona has developed a sufficient conduction zone to smooth out perturbations. Implementing zooming on OMEGA would require new phase plates, referred to as zooming phase plates (ZPP’s) and co-propagating dual driver lines. Using laser beams with multiple wavelengths has been proposed to mitigate CBET in direct-drive implosions. Color-splitting the beams into two or more co-propagating wavelengths with $\Delta \lambda > 5$ Å reduces CBET by ~50% in 1-D modeling. Instead of each beam containing multiple wavelengths, the beams could be grouped into subsets of monochromatic beams with distinct wavelengths. The current in-line models will not capture the full effect of the 3-D beam distribution because of their respective 1-D and 2-D approximations.

In this article, the effects of frequency detuning laser beams in direct-drive symmetric implosions are studied using a 3-D CBET model. To our knowledge, this is the first fully 3-D modeling of CBET for direct-drive implosions. The 3-D ray-based CBET model was benchmarked against full-field calculations, providing confidence in the implementation of the model to calculate the effects of CBET in full-scale implosion experiments. These calculations show that interactions between beams with relative angles between 45° and 90° are the most significant for CBET in OMEGA direct-drive implosions. Redistribution of laser power because of CBET can increase the rms absorption nonuniformity by an order
of magnitude. Shifting the relative wavelengths of three groups of laser beams by ~10 Å maximized the total absorption, and the rms absorption nonuniformity was near minimum for the implosion conditions studied in this article.

The following sections discuss the model equations and gridding; report on model results for two-beam and many-beam CBET coupling when all the beams are launched with the same wavelength; and present predictions for a CBET mitigation scheme in 60-beam symmetric OMEGA direct-drive implosions using wavelength detuning.

The 3-D CBET Model

The CBET model used here (BeamCrosser) was originally developed as a MATLAB-based hydrodynamics code postprocessor to simulate scattered-light spectra from OMEGA implosions and provided the first evidence that CBET was significantly degrading implosion performance relative to 1-D hydrodynamic predictions. As a hydrodynamics code postprocessor, the CBET model relies on time-varying coronal plasma parameters calculated independently by a hydrodynamics code such as LILAC (1-D) or DRACO (2-D). The CBET model is used to gain insight into 3-D effects during an implosion, even though its calculations are not fully self-consistent with the plasma hydrodynamics.

1. Ray Tracing and Model Gridding

The model is a ray-based CBET model and therefore does not solve the full electromagnetic Maxwell equations. The reduced ray equation for geometric optics is used to determine the laser beam propagation through the coronal plasma of an implosion. The ray equation is solved by a fourth-order Runge–Kutta method similar to that of Sharma et al. A single geometric optics-based ray is simply a path through space that by itself carries no inherent information about the local light intensity. The laser intensity along a ray is calculated using the intensity law of geometrical optics and the spacing between points of neighboring rays on the same wavefront (having equal optical path lengths) along with the change in intensity caused by absorption and CBET, as discussed in the next section.

An example of the ray paths for a single laser beam propagating through a spherically symmetric coronal plasma of a direct-drive implosion is shown in Fig. 150.11. Refraction of the rays produces a paraboloid-like shadow behind the target inside of which rays cannot reach. The envelope of tangential rays defining the boundary surface of the shadow volume is a caustic of the “fold catastrophe” type, where rays fold upon themselves and all points on the unshadowed side of the caustic are intersected by two distinct rays. The points where the rays graze the caustic are sometimes referred to as their “turning points” but that is not true in general. The turning point is best defined as the point of deepest radial penetration into the coronal plasma. It is clear from the outermost rays shown in Fig. 150.12 that these rays cross well away from their closest approach to the target. The intensity law of geometrical optics breaks down near the caustic, where the spacing between points on the same wavefront becomes very small, giving unphysically high intensities. The actual intensity where two rays cross is limited by diffraction. The intensity near the caustic is similar to an Airy pattern superimposed on the intensity from the geometric law. The CBET model limits the intensity along the rays from reaching unphysically high intensities by using either a fixed limiting factor or a field swelling limit based on Krue; however, this limit has only a small effect on the results of the code since it is applied only in a small volume.

\[ \Delta \omega = \omega_L \frac{\partial \tau}{\partial t}, \]  

(2)
ized Doppler effect is essential because the change in frequency between incoming and outgoing rays in a typical direct-drive implosion is of the same order as the difference in the frequency needed for CBET (approximately a few angstroms).

The model “gridding” follows the beam trajectories determined by the ray trace. Each on-target laser beam profile is discretized into many square “beamlets” with flat intensities on a 2-D grid, as shown in Fig. 150.13(a). The distance along the path of each beamlet provides the third dimension for the gridding of each beam. Figure 150.13(b) shows this non-orthogonal overlapping grid in which more than one cell for a single beam can occupy the same physical space. CBET at these intrabeam crossings between beamlets from the same beam is calculated by the model in addition to crossing between beamlets from different beams. The significant refraction of the laser light in a direct-drive implosion plasma is a major difference from indirect drive, where the refraction of the laser beams can typically be ignored and the paraxial approximation can be employed in the volume where the beams cross.

\[
dP_b(s) = P_b(s) \left\{ \frac{1}{L_{\text{abs}}} + \sum_{\text{all beam crossings}} \frac{C_{\text{CBET}}}{L_{\text{CBET}}} \right\} ds, \tag{3}
\]

where \(L_{\text{abs}}\) is the scale length of inverse bremsstrahlung absorption, \(C_{\text{CBET}}\) is a multiplier, typically of the order of 2 for implosion modeling, applied to the calculated CBET coupling to better match experimental measurements (discussed in detail below). \(L_{\text{CBET}}\) is the local spatial rate of energy gain/loss because of CBET in the strong damping limit:

\[
L_{\text{CBET}}^{-1} = 5.85 \times 10^{-2} \frac{1}{n_a} \frac{n_e/n_c}{(1-n_e/n_c)} \times I_{14} \lambda_{0,\mu m} \times T_{e,\text{keV}} \left( 1 + 3T_{i,\text{keV}}/ZT_{e,\text{keV}} \right) \times \psi \times R(\eta)(\mu m^{-1}), \tag{4}
\]

where \(\lambda_{0,\mu m}\) is the laser wavelength in microns, \(I_{14}\) is the crossing laser intensity in \(10^{14}\) W/cm\(^2\), \(T_{e,\text{keV}}\) and \(T_{i,\text{keV}}\) are the electron and ion temperatures, respectively, in keV, \(Z\) is the average ionization, \(n_a\) is the dimensionless amplitude damping rate for the IAW, \(n_e\) is the electron density, and \(n_c\) is the critical density. Since \(L_{\text{CBET}}\) depends on the intensity of the crossing beamlets, Eqs. (3) and (4) form a set of coupled nonlinear equations.

The factor \(R(\eta)\) is the resonance function accounting for how closely the driven wave satisfies the IAW dispersion relation [Eq. (1)],

\[
R(\eta) = \frac{\nu_a^2 \eta}{(\eta^2 - 1)^2 + \nu_a^2 \eta^2}. \tag{5}
\]

The factor \(\psi\) accounts for the effect of polarization on the coupling of the crossing beams. For random polarization or when the beams have their polarization evenly distributed in two orthogonal components,

\[
\psi = \frac{1}{4} \left( 1 + \cos^2 \theta_k \right), \tag{6}
\]
where $\theta_x$ is the beam crossing angle. This is appropriate for most implosions on OMEGA, where distributed polarization rotators (DPR’s) are used to split the beams into orthogonal polarizations, or on the NIF, where the beams are arranged in quads in such a way that the polarizations of two beams are orthogonal to the polarization of the other two beams in the quad. When DPR’s are not used on OMEGA, the beams are linearly polarized and the coupling between the beams affects only the shared polarization component.

As mentioned above, outside the beam shadow, all points are crossed twice by rays from each beam, so there are a total of $2N-1$ possible crossings to be considered at each point along a beamlet, where $N$ is the total number of beams. Equations (4)–(6) are applied to all beamlet crossings to determine the total CBET coupling at each grid point along all beamlets. Since pump depletion is inherent in CBET, the system is solved using fixed-point iteration. Energy is conserved by balancing the power exchanged between beamlets such that the power gain (loss) calculated for beamlet A where it is crossed by beamlet B is identical to the power loss (gain) for beamlet B where it is crossed by beamlet A.

3. Benchmarking the Model

The CBET model was benchmarked by comparing it with the predictions of a full-wave code LPSE (laser-plasma simulation environment). LPSE solves the time-enveloped Maxwell equations coupled to a linearized time-dependent fluid plasma response to calculate the enveloped electric-field vector and the ponderomotively driven ion-density perturbations. LPSE is impractical for full-scale 3-D implosion modeling because of its computational costs, but full-scale 2-D and reduced-scale 3-D runs provide good benchmarks for a ray-based model.

Figure 150.14 shows an LPSE calculation of two lasers crossing in a constant-density plasma with a linearly varying plasma velocity profile that places the maximum of the resonance function [Eq. (2)] at $x = 8.6 \mu m$. Both beams are polarized 45° out of the plane. CBET affects only the components of the polarization that are shared by the beams, so the polarization of each beam is expected to rotate. The intensities of the beams after undergoing CBET predicted by the ray-based model (with $C_{\text{CBET}} = 1$) are an excellent match to those predicted by LPSE [Fig. 150.14(b)]. The predicted rotation in the polarization caused by CBET is very good over the region where the beam power is significant, but some divergence between the calculations is observed where the beam intensities are small. Overall, the comparison with the full-field calculations of LPSE provides confidence on the validity of the approximations made in the ray-based code.
ent in direct-drive implosions are missing from the ray-based models. Possible candidates for the missing phenomena include diffraction, polarization details, and nonlinear multibeam effects. All predictions presented here, unless mentioned otherwise, use a factor of \( C_{\text{CBET}} = 2 \).

**Beam Coupling with No Wavelength Shift**

In this section, the coupling between OMEGA beams during a direct-drive implosion is modeled when all beams are launched with the same wavelength (351 nm). It is important to note that although all beamlets enter the plasma with the same wavelength, the Doppler effect changes the wavelength of each beamlet as it passes through the plasma. This wavelength shift varies across the beam profile depending on the path each beamlet takes through the coronal plasma. The magnitude of the Doppler shift is of the order of a few angstroms and must be included when calculating the CBET coupling along a beamlet.

All of the simulations use 1-D LILAC predictions with a nonlocal electron heat transport model of the coronal plasma conditions for a typical OMEGA symmetric direct-drive implosion of a CH target (shot 60,000). The plasma profiles were taken from a single time late in the pulse when CBET is predicted to be largest. Figure 150.15 shows the distribution of the 60 OMEGA laser beams. All beams use a super-Gaussian of the order of 4 (SG4) intensity profile measured for the SG4 distributed phase plates (DPP’s) used in the implosion. All beams entered the plasma with 0.35 TW of power, which was the nominal power of all the beams in the implosion late in the pulse.

1. Two-Beam CBET Calculations

   In a direct-drive implosion, each beam can interact with all other beams. In a nominally symmetric implosion, all beams have identical beam powers, intensity profiles, and relative geometries (i.e., the “view” from each beam looks the same with respect to the relative positions of the other beams). There is zero net exchange of total power between the beams, but there will still be a redistribution of power because of CBET. It is useful to determine which of the other beams has the strongest exchange with any single beam and the effect of that exchange on the effective beam intensity profile. How CBET affects the exchange between any two specific pairs of beams is mainly dependent on the angle between the two beams.

   Figure 150.16 shows the laser absorption for two-beam simulations, where the angle \( \theta \) between the beams was varied; \( \theta = 180^\circ \) indicates beams launched on opposing sides of the target. Because the coupling of any two beams is small compared to the total interaction between a set of 60 beams in an OMEGA symmetric implosion, a CBET multiplier of \( C_{\text{CBET}} = 5 \) was used for these two-beam interactions to accentuate the effects of CBET. The degradation in absorbed power is caused by redistribution of the beam power and is identical for both beams because of symmetry. The absorbed power is degraded most strongly by beams that are separated by 45° to 90°. Beams separated by more than 135° are practically decoupled. The absorption for these nearly opposite beams is essentially the same as when only intrabeam CBET interaction between a single beam and itself is considered. For \( 0^\circ \) the beams are copropagating and the laser absorption is the same as the intrabeam CBET of a single beam with twice the intensity. Figure 150.16(b) shows the effective importance of CBET between beams at different angles. The effective importance of a direct-drive implosion depends on the number of beams at that angle. For an infinite number of beams, the importance of CBET for beams at a specific angle is determined by the change in absorption of the beam from Fig. 150.16(a) weighted by the differential surface area of a sphere for that angle \( \sin \theta \cdot d\theta \). The normalized change in absorption because of CBET weighted by \( \sin \theta \) is shown by the solid red line in Fig. 150.16(b). Compared to Fig. 150.16(a), the importance of different beams is skewed toward the equator where the differential area is maximum. For a finite set of beams, the effect of CBET from one beam at a specific angle is weighted by the actual number of beams at that angle. The importance of different beams weighted for the symmetric 60-beam OMEGA geometry is shown by the
Because of CBET integrated along the path of each beamlet (as described in Ray Tracing and Model Gridding, p. 63), the beamlets near the horizontal axis experience a net loss and those on the beam edge closest to the other beam experience a net gain. The total absorbed power in each beamlet is shown in Fig. 150.17(b). Near the beam center there is a region of lower absorbed power caused by the CBET losses. The overall absorption profile is radially asymmetric, and there is a shift in the centroid of the absorption away from the other beam compared to the no-CBET absorption profile [Fig. 150.17(c)].

Figure 150.17 shows the redistribution of power in the beam profile for two laser beams at a relative angle of 90°. Although CBET was known to shift the centroid of the outgoing beam profiles for simple beam geometries,\textsuperscript{20,25} the redistribution is complicated for beams refracting through a spherical plasma where not all beamlets encounter a resonance with the other beam. Figure 150.17(a) shows the power gain and loss because of CBET integrated along the path of each beamlet.

Figure 150.18 shows the CBET exchange and absorption profiles calculated for the coronal plasma conditions modeled with the full 60 beams on OMEGA. In a direct-drive implosion, CBET’s total effect on a beam is the sum of its interactions with all other beams. The absorbed power is significantly less in magnitude and shows a more-complicated profile structure than the two-beam case. As in the two-beam case, there is...
no net exchange in power between beams because of beam symmetry but there is significant redistribution of power. The effect of many beams extends the ingoing losses over the hydrodynamically efficient small impact parameter beamlets and distributes the net outgoing gain to a ring of less hydrodynamically efficient high-impact parameter beamlets. The change in the absorbed power profile from the single-beam, no-CBET profile [Fig. 150.16(c)] has a significant effect on the absorption uniformity over the target surface.

Figure 150.19 shows an absorption surface map calculated by radially integrating the 3-D absorption at all points over the target surface. When CBET is ignored, the absorption is very uniform with an rms variation of ~0.2%. When CBET is included, the nonuniformity is an order-of-magnitude larger with an rms of 2.0%. Since there is no azimuthal symmetry in the absorption surface map, the effects of this nonuniformity cannot be captured by a 1-D or 2-D hydrodynamics code.

### CBET Mitigation Using Wavelength Detuning

One possible scheme for CBET mitigation during symmetric direct-drive implosions is wavelength detuning between groups of beams. The 60 beams of the OMEGA Laser System originate from a single seed-pulse driver that is split three ways into “legs” and amplified separately to feed 20 beams each. The beams from each leg are distributed around the target chamber (Fig. 150.15). By shifting the wavelength of two of the legs in such a way that all three legs have different wavelengths, the CBET coupling between the groups of beams fed by the legs could be altered. With sufficient wavelength shifting, the groups of beams could be effectively decoupled from each other.

#### 1. Single-Beam Wavelength Shift

Figure 150.20 shows the effect on absorbed laser power of shifting the wavelength of a single beam in a 60-beam symmetric implosion while keeping the other 59 beams fixed at 351 nm. The modeling predicts that small wavelength shifts can significantly increase or decrease the power absorbed in the single wavelength-shifted beam and that it takes a wavelength shift of $|\Delta \lambda| > 30 \, \text{Å}$ to completely decouple the beam from the other 59 beams. The behavior shown in Fig. 150.20

![Figure 150.20](image)

**Figure 150.20**

Absorbed power per beam in the OMEGA 60-beam geometry when one beam is wavelength shifted. The plot with red dots is the absorbed power in the single wavelength-shifted beam. The solid blue curve is the average absorbed power for the other 59 beams. The dotted black line is the absorbed power for a beam completely decoupled from all other beams but still experiencing CBET with itself.
is complicated because of the complex 3-D crossings with 59 other beams, but the behavior can be broken down into a few general phenomena.

When $\Delta \lambda$ is increased negatively, less Doppler shift is required to satisfy the IAW dispersion relation [Eq. (1)] for power loss in the central beamlets entering the plasma. This moves their resonance to smaller radii, where the plasma velocity is lower and the density is higher. Here the CBET coupling is stronger [Eq. (4)], increasing the losses compared to the $\Delta \lambda = 0$ case. At the same negative wavelength shift, the resonance for the power gain of the large-impact parameter beamlets moves radially outward, where the coupling parameter is weaker, thereby decreasing their power gain. Both effects reduce the total power absorption for the shifted beam, but the increased losses of the central beamlets are the primary source of the sharp drop in absorbed power shown in Fig. 150.20 as $\Delta \lambda$ is increased negatively. The absorbed power reaches a minimum at a shift of $\Delta \lambda \approx -2$ Å and then rises again as the resonance location for the power loss moves inside the beam shadow (Fig. 150.12) and CBET losses decrease until the beams decouple at $\Delta \lambda < -30$ Å. When $\Delta \lambda$ is increased positively, the resonance shifts are reversed, reducing the losses of the central beamlets, increasing the gains in the outer beamlets, and producing a sharp rise in the absorbed power that peaks near $\Delta \lambda \approx 3$ Å. Because the beam can gain energy from the 59 other beams, the total power absorbed from the shifted beam can exceed the original power in that beam (0.35 TW). For positive $\Delta \lambda$, a second maximum occurs near $\Delta \lambda \approx 18$ Å. This broad peak occurs because the wavelength shift is large enough to change the direction of CBET for the incoming central beamlets such that they gain energy while entering into the plasma from the other 59 beams.

2. Three-Leg Wavelength Shifts

Figure 150.21 shows the effect that shifting the wavelength of the three OMEGA legs has on the absorbed power and its uniformity over the implosion target for the coronal plasma conditions modeled. When $\Delta \lambda$ is given as the wavelength shift, it means that the beams in leg 1 are wavelength shifted by $-\Delta \lambda$ and those in leg 3 are shifted by $+\Delta \lambda$, while the beams in leg 2 remain unshifted at 351 nm. As $\Delta \lambda$ is increased from zero, CBET losses in leg 1 increase and the absorbed power in leg 1 beams drops, while the opposite occurs in leg 3, whose gains increase from CBET [Fig. 150.21(a)]. This loss/gain grows sharply until about $\Delta \lambda \approx 2$ Å. Here, the difference in absorbed power between legs 1 and 3 is maximum. As $\Delta \lambda$ is increased further, the CBET coupling between the legs decreases, and, as a result, the difference between their absorbed power decreases until $\Delta \lambda > 30$ Å, where the legs are essentially decoupled. If $\Delta \lambda$ is negative, these effects remain the same except the roles of leg 1 and leg 3 are reversed. Of particular interest is the region where $\Delta \lambda \approx 10$ Å. Here, the absorbed power averaged over all 60 beams is higher than the decoupled case ($\Delta \lambda > 30$ Å), indicating that CBET may work in favor of increased implosion drive.

Figure 150.21(b) shows that as $\Delta \lambda$ is increased from zero, the absorption nonuniformity (rms of the absorbed energy over the target surface) increases sharply to a maximum.
around $\Delta \lambda \equiv 2$ Å, then falls as $\Delta \lambda$ continues to increase and the legs decouple. For $\Delta \lambda > 8$ Å, the absorption nonuniformity changes only weakly, but at the same location as the maximum total absorption ($\Delta \lambda \approx 10$ Å), there is a local minimum in the absorption nonuniformity of 1.3%, which is almost as small as the value when the legs are completely decoupled (1.2%). Figure 150.22 shows absorption surface maps calculated for wavelength shifts between the legs of 2 Å, 6 Å, and 10 Å. Not only does the total absorption rms change with $\Delta \lambda$, but the surface pattern of the nonuniformity varies as well.

**Summary**

A fully 3-D modeling of CBET for direct-drive symmetric implosions has been used to investigate the effects of wavelength detuning on CBET. The 3-D ray-based CBET model was benchmarked to full-wave calculations, providing confidence in the implementation of the model. For this study, coronal plasma conditions from late in the drive pulse of a typical warm OMEGA implosion were modeled. The model calculations show that beams with relative angles between 45° to 90° are most significant for CBET in OMEGA direct-drive implosions. The redistribution of laser power by CBET increases the absorption rms nonuniformity by an order of magnitude. Implosion degradation effects resulting from this increase in absorption nonuniformity from CBET should be studied by 3-D hydrodynamic modeling. By shifting the relative wavelengths of three groups of laser beams by $\pm 10$ Å, the total laser absorption was maximized and the rms absorption nonuniformity nearly minimized for these plasma conditions.

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**REFERENCES**


Picosecond Time-Resolved Measurements of Dense Plasma Line Shifts

Understanding the time-averaged and time-dependent response of ions in dense plasma is important for correctly interpreting and modeling atomic structure and radiation transport in extreme environments, including stellar atmospheres and imploding inertial fusion capsules. The potential in and near an ion immersed in a dense plasma is influenced by its bound electrons, free electrons, and neighboring ions. These influences can change the radiative and thermodynamic properties of the plasma by modifying the energy levels that are available to an ion and causing energy-level shifts.

Dense plasma line shifts originate from free-electron modification of the ionic potential. Free electrons in the ion sphere screen the nuclear charge and shift bound energy levels toward the continuum. The energy separation between levels is decreased for a given bound-electron configuration, and emission lines are shifted to lower photon energies. Correct identification of this effect helped resolve the disparity between spectroscopically inferred white dwarf masses and the results from other measurement methods and the predictions of general relativity.

While dense plasma line shifts have been described mathematically and confirmed experimentally, few measurements have tested line-shift model predictions at high energy density, leading to an incomplete picture of how this process is understood and modeled in extreme conditions. Data comparisons to line-shift model predictions in this regime have been hampered by the difficulty of obtaining uniform, well-characterized, and high-energy-density (HED) plasmas. Line-shift measurements are especially needed in hot dense plasmas to provide a stringent test for analytic and detailed atomic kinetics and radiative-transfer calculations. Equally important, line-shift measurements can provide a sensitive benchmark for free-electron distribution models within the ion sphere—an important application of the N-body problem.

This article reports the first picosecond time-resolved measurements of dense plasma line shifts of the 1s2p−1s2 transition in He-like Al ions as a function of the instantaneous plasma conditions. Line shifts were measured with picosecond time resolution for free-electron densities of 1 to $5 \times 10^{23}$ cm$^{-3}$ and temperatures of 250 to 375 eV. The plasma conditions were inferred with well-quantified errors from spectroscopic measurements of the Al He$_x$ complex. The data are compared to a generalized analytic model proposed by Li et al. based on a parameterization of numerical ion-sphere model calculations. The predicted line shifts show broad agreement with the data over the full range of densities and temperatures studied, with evidence for deviation from the experimental data at the most-extreme densities. This work provides an experimental test of a simplified method that calculates quantum-number-dependent energy level shifts for ions in dense, finite-temperature plasma.

The experiments were carried out at LLE’s Multi-Terawatt (MTW) Laser Facility. Figure 150.23(a) shows the experimental setup. The laser directly irradiated small-mass targets with 0.7-ps, ≤16-J pulses at the laser’s fundamental wavelength ($\lambda = 1054$ nm) or second harmonic. The laser was focused to an ~5-μm full-width-at-half-maximum (FWHM) focal spot by an f/3 off-axis parabolic mirror at normal incidence to the target at intensities greater than $10^{18}$ W/cm$^2$. The targets were thin plastic foils with a buried Al microdot. The microdot was vacuum deposited in a 0.2-μm layer on a 1-μm parylene-N (CH) support. The thickness of the front parylene-N overcoat was varied between 0 and 2 μm to access different plasma densities. The Al microdot was kept purposefully thin to limit spatial gradients, while the outer plastic layers constrain sample expansion to achieve near-solid-density conditions.

High-intensity laser pulses with low contrast reduce the maximum electron density that can be achieved in buried-layer target interactions by causing the target to prematurely heat and decompress. In the experiments reported here, free-electron densities of up to $2 \times 10^{23}$ cm$^{-3}$ were achieved at the laser’s fundamental wavelength with a measured temporal contrast of the order of $10^9$ up to 100 ps prior to the main pulse. Experiments with high-contrast, frequency-doubled pulses achieved free-electron densities of the order of $5 \times 10^{23}$ cm$^{-3}$. Based on work by previous authors and the measured contrast of the 1ω
beam, the temporal contrast of the frequency-doubled beam is estimated to be $10^{12}$ (Ref. 15); however, no on-shot contrast measurement was available for this particular experiment. Residual $1\omega$ light was rejected at a spectral contrast of the order of $10^{12}$ by six transport mirrors with 99% $1\omega$ extinction coatings.

Picosecond streaked x-ray spectroscopy was used to infer the density and temperature of the Al layer. For this measurement, a conically curved potassium acid phthalate (KAP) streaked x-ray spectrometer was used in combination with a time-integrating flat pentaerythritol (PET) crystal spectrometer. The streaked spectrometer was configured to study Al He$_\alpha$ ($1s2p - 1s^2$) thermal line emission with spectral resolving power $E/\Delta E \sim 1000$ and 2-ps temporal resolution.\textsuperscript{16} Time-integrated spectra were measured on each shot and used to correct the streaked spectra for variations in spectral sensitivity introduced by the photocathode.\textsuperscript{17}

It is noted that the plasma-induced line shifts measured in the experiment could be exaggerated or disguised by streak camera charge-coupled–device (CCD) clocking errors. The CCD clocking was measured offline using a structured photocathode illuminated by a static x-ray source. The tests identified a 0.46±0.01° correction that was applied to the experimental data. Experimental tests with low-density, laser-driven Al plasmas confirm that the clocking offset was properly corrected.

Figure 150.23(b) shows example streaked data where the dispersion of the streaked x-ray spectrometer was determined in situ from the emission lines of He- and Li-like Al ions at low plasma density. A well-resolved emission spectrum was selected after the plasma was allowed to expand for 12 ps after the high-intensity pulse [Fig. 150.24(a)]. The initial He$_\alpha$ and intercombination line positions cannot be directly identified with tabulated transition energies because the plasma environment modifies the ionic energy-level structure. The Li-like satellite lines are not expected to shift measurably because of the screening effect of the $n = 2$ spectator electron.\textsuperscript{18} The
green trace in Fig. 150.24(a) corresponds to the latest time in the plasma evolution that could be reliably measured, corresponding to \( t_0 + 275 \text{ ps} \), where \( t_0 \) is the arrival time of the high-intensity laser pulse at the target. At this time, the plasma consists of approximately isolated radiators and the He\( \alpha \) resonance and intercombination lines are observed at a higher photon energy. The observed positions of these lines are constant in time and provide an absolute energy fiducial to register the calibration. Additional Al K-edge measurements were carried out to verify the calibration. A 2-\( \mu \)m Al filter placed in situ over the detector aperture was backlit by laser-produced x rays and the K-shell absorption edge at 1559.6 eV was recorded. The measured edge location is free from plasma effects and provides an absolute energy fiducial to confirm the dispersion slope and offset.

The measured fiducial position \( P \) was related to the photon energy \( E \) via Bragg’s law for the spectrometer geometry:

\[
P (\text{pixel}) = A \cdot \tan \left( \frac{hc}{2dE} \right) + B, \tag{1}
\]

where \( A \) and \( B \) are fitting parameters, \( h \) is Planck’s constant, \( c \) is the speed of light, and \( 2d = 26.64 \) \( \text{Å} \) is the Bragg spacing of the KAP crystal. The results of the calibration are shown in Fig. 150.24(b). A conservative estimate for the uncertainty in peak position yields two pixels, or 0.25 eV. Uncertainty in the location of the K edge is slightly larger because of degraded spectral resolution at the edge of the streak camera field of view, where the K edge was measured. This uncertainty provides the dominant contribution to the calculated error in the measured shifts. The overall uncertainty in the dispersion is within the width of the data points in Fig. 150.24(b). This dispersion was applied to all data collected in this work. An important point for these measurements is that the dispersion was established self-consistently without reference to plasma-dependent fiducials. Previous work\(^{19,20}\) relied on time-averaged measurements of the K\( \alpha \) line shape to set the dispersion. The technique presented here provides the dispersion over the full spectral range of interest, registered to the cold Al K-shell absorption edge.

Figure 150.25(a) shows the streaked data gathered from an Al layer heated by a 16-J, 0.7-ps pulse focused to \( -1 \times 10^{19} \text{ W/cm}^2 \). The Al was enclosed on both sides by a 1-\( \mu \)m parylene-N tamper. The data were corrected for the streak tube’s geometric curvature and variations in photocathode spectral sensitivity. The time axis was calibrated in separate tests. The Stark-broadened resonance line and the commensurate strong satellite emission indicate high plasma density over the duration of the experiment. Spectra were averaged over a five-pixel temporal window (shaded region) corresponding to the streak camera’s temporal impulse response. A linear background was removed from the data.

The He\( \alpha \) FWHM and He\( \alpha \)-to-satellite intensity ratio was interpreted using a nonlocal-thermodynamic-equilibrium (NLTE) collisional-radiative atomic model\(^{21}\) to infer the Al density and temperature as a function of time.\(^{2,22}\) The model calculated synthetic spectra for Al IX–XIV ions over a regular density and temperature grid of 0.1 to 6 g/cm\(^3\) and 100 to 600 eV, respectively. The simulation was carried out in 1-D for a 0.2-\( \mu \)m Al slab. Satellite contributions to the line shape were treated in detail by including transitions from all ions with non-negligible populations. Satellite blending with the

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![Figure 150.25](E25516JR)

(a) Streaked He\( \alpha \) emission from a buried Al layer (1 \( \mu \)m CH). (b) The plasma conditions are inferred from the He\( \alpha \) intensity ratio (red) and full width at half maximum (FWHM) (blue). (c) The synthetic spectrum corresponding to \( T_e = 330 \text{ eV} \) and \( \rho = 0.9 \text{ g/cm}^3 \) (red) is compared to the data (blue).
resonance line may otherwise be misinterpreted as spurious Stark broadening or red shift. The effects of radiation transfer were included using an escape probability approach based on local escape factors to calculate photoexcitation rates. The line profiles were calculated with the effects of Doppler, Stark, natural, Auger, and opacity broadening. The synthetic spectra were convolved with the detector resolving power, and the He shifts caused by line broadening and satellite blending were calculated from the model to verify that the observed shifts were caused by true plasma effects.

The experimental He-to-satellite intensity ratio was formed over the same spectral bands as the synthetic data. The error was determined from a Monte Carlo study that produced a ratio distribution from uncertainties introduced by photon statistics and the analysis procedure. The distribution variance characterizes the ratio error with coupled sources of uncertainty. For the lineout selected in Fig. 150.25(a), the ratio was 2.5±0.5. Figure 150.25(b) shows the temperature and density contour specified by the measured ratio (red). The He FWHM provides a second measurement to constrain the inferred conditions. The FWHM was measured from a spline fit through the data to minimize random error introduced by statistical signal fluctuations. Noise was considered separately as a source of uncertainty by calculating the probability distribution for the measured width based on the likelihood that statistical signal fluctuations could be spuriously detected as FWHM crossing points. For the data shown, the width was 5.3±0.6 eV. The measured FWHM specifies a second contour in temperature and density space (blue) that constrains the inferred temperature and density. The width of the two contours and the size of the overlap region are related to the uncertainties in the measured quantities.

An estimate for the true temperature and density is calculated from the mean temperature and density in the overlap region. The error in the estimated temperature and density corresponds to the extent of the overlap region along each axis. The conditions were inferred to be 330±56 eV and 0.9±0.3 g/cm³ (nₑ = 2.2±0.8 × 10²³ cm⁻³). Figure 150.25(c) shows the unfiltered spectrum and the model prediction for the measured conditions. The model considers the instrument resolving power and reproduces the experimental data well. It is noted that the data have been uniformly shifted to higher photon energies by 2.4 eV for comparison with the synthetic spectrum since the atomic kinetics model used here does not include dense plasma line shifts.

The spectral shifts were quantified by the difference between the first moment of the measured He line shape and the predicted (unshifted) He line shape. The limits of integration were selected to fully encompass the He line without contamination from the intercombination line. The uncertainty in the measured shift was calculated assuming independent contributions from dispersion calibration error and statistical signal fluctuations. The measured shift for the spectrum shown in Fig. 150.25(c) was 2.4±0.3 eV.

Figure 150.26(a) shows the measured He line shifts for inferred electron densities from 1 to 5 × 10²³ cm⁻³. The dataset is composed of well-resolved spectra with no self-reversal. The error bars are shown for a few representative data points at low, medium, and high densities. The asymmetric vertical error bars reflect the uncertainty in the measured location of the Al K-shell absorption edge used to register the absolute dispersion calibration. This uncertainty does not propagate to the inferred temperature and density since those quantities are sensitive only to the relative dispersion. The data show a nearly linear shift with increasing electron density. The highest electron densities measured were near 80% of solid and were achieved with the high-contrast 2ω drive. The data were selected over a small range of temperatures between 250 and 375 eV.

For each data point, an apparent shift was calculated from the atomic kinetics model to confirm that the observed shifts were not spurious. These data are plotted in red and the typical error is within the data point. The magnitude of the apparent shift appears to decrease with electron density because the 2ω drive for high-density studies produced proportionately higher temperatures that suppress satellite enhancement of the red wings of the line profile.

Predictions from a generalized analytic ion-sphere model proposed by Li et al. are compared to the data in Fig. 150.26(a). The analytic approach relies on the self-consistent field ion-sphere model (SCFISM) to obtain the self-consistent density distribution of bound and free electrons within the ion sphere. Relativistic atomic structure calculations of the bound wave functions are carried out in a screened nuclear potential determined from the electron density distribution. Detailed scaling studies were performed to obtain a generalized density- and temperature-dependent formula for the energy level structure in the plasma.
For each data point in Fig. 150.26(a), the analytic formulation was used to predict the line shift at the inferred temperature and density. The calculation is monoenergetic and considers only the $1s2p-1s^2$ transition. The calculation agrees well with the data at low density but diverges at higher densities, likely because the calculation neglects unresolved satellites and other contributions to the line shape. An attempt was made to recover this information by adding the apparent shifts to the ion-sphere model predictions. This addition produces better agreement at low and moderate densities, as shown in Fig. 150.26(b). The error bars indicate uncertainty in the calculated shifts based on the density uncertainty in the corresponding data point. The temperature uncertainty was neglected since the model exhibits a weak dependence on electron temperature ($\sim T_e^{0.25}$).

Sensitive spectral measurements of this nature may prove to be a valuable test of electron screening models in extreme conditions. For example, improved agreement between the analytic model and the experimental data is obtained for an ad hoc 10% reduction of the ion-sphere radius [Fig. 150.26(b)]. The optimum reduction was determined by a single-parameter maximum-likelihood least squares fit to the data that considered uncertainties in the inferred densities and line shifts. The quality of the match is surprisingly sensitive to the scale factor. It is noted that a 40% increase in the inferred densities can reproduce the improved agreement. It is unlikely, however, that the analysis overestimates the inferred densities by that amount, and it is improbable that densities above solid were achieved in the experiment.

Moreover, spectroscopic temperature and density measurements from K-shell ions can be sensitive to the choice of atomic model. Recent work has demonstrated that model discrepancies can contribute up to 30% uncertainty in the inferred conditions, mostly from uncertainty in Stark-line–shape calculations. Model-dependent temperature and density measurements will become more reliable as theoretical and experimental work further validates line-shape models.

In summary, spectral line shifts of the $1s2p-1s^2$ transition in He-like Al ions have been studied as a function of the instantaneous plasma conditions at high energy density. The line shifts were measured using a picosecond time-resolved x-ray spectrometer with an absolutely calibrated spectral dispersion. Buried-layer targets and a high-contrast $2\omega$ laser driver provided access to densities near 80% of solid. The plasma conditions were inferred by comparing the measured spectra to calculations from a NLTE collisional-radiative atomic physics model. A generalized analytic line-shift model was found to be broadly consistent with the experimental data for all but...
the highest densities studied. These findings are important to understanding plasma-dependent atomic structure and radiation transport in high-energy-density environments.

ACKNOWLEDGMENT
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REFERENCES
A Framed, 16-Image Kirkpatrick–Baez X-Ray Microscope

Introduction

Kirkpatrick–Baez (KB)–type x-ray microscopes are one of the principal methods of imaging x-ray emission from laser-generated plasmas. They typically have a larger collecting solid angle, better spatial resolution, and larger standoff distance than the simpler method of pinhole imaging. They have been used on both the 60-beam OMEGA Laser System and the previous 24-beam OMEGA Laser System. High spatial resolution (~3 μm) has been demonstrated using four-image KB mirror assemblies, which, when framed, achieved a resolution of ~5 μm.

An advantage of using pinholes to image the plasma x-ray emission is that when coupled to a multistrip, high-speed framing camera, many images can be obtained with a time interval as short as ~30 ps (dictated by the separation of the pinhole images and the voltage propagation speed across the strip). Until recently, KB microscopes have been limited to just four images with larger image separation (52 mm) and corresponding longer time separations (~350 ps) when two images are coupled to a single-strip framing camera. Pickworth et al. have recently developed a KB mirror assembly for use at the National Ignition Facility capable of being coupled to a four-strip, high-speed framing camera. Additionally, Yi et al. have implemented an eight-image KB mirror assembly also coupled to a four-strip framing camera. These previous limits have been removed by the use of compact KB microscope mirrors whose design has increased the number of images to 16, which, when properly aligned, can be coupled to a four-strip, high-speed framing camera having strip separations of 9 mm. The assembly of compact KB mirrors that makes this image alignment possible has been accomplished for the first time, as described in this article. For image separations of 9 mm, along the strip, the corresponding image-to-image time separation is 60 ps. The sampling time interval can be decreased to 15 ps by using cables that delay the pulses to the strips by 15-ps intervals. This has been achieved in the instrument described in this work known as KBFRAMED.

The 16-Image KB Optic

The design of a 16-image KB microscope was originally put forth by Marshall, Oertel, and Walsh. In this design, mirrors were cut so they would fit together in a perfect 16-sided polygon, i.e., a hexadecagon. The resulting array of image locations falls on a circle; therefore, a framing camera with circular photocathode strips is needed to frame these images. Subsequently, Marshall proposed a modification to the ideal hexadecagon arrangement of the mirrors that would allow images to be relocated to fall on the rectangular strips of the modern high-speed framing-camera design. The KB mirror focus [Eq. (1)] is given by

$$\frac{1}{p} + \frac{1}{q} = \frac{2}{R \sin i},$$  \hspace{1cm} (1)

where $p$ is the distance from object to mirror, $q$ is the mirror image distance, $R$ is the mirror radius of curvature, and $i$ is the angle of incidence of x rays at the mirror center. The basic concept is to simultaneously move and tilt the mirror, maintaining the focus condition while repositioning the image (Fig. 150.27). The pattern of 16 images can in this way be repositioned to fall on the cathode strips of a high-speed framing camera that are nominally 9.0 mm apart [see Fig. 150.28(a)]. For a mirror pair, each mirror obeys a separate focus equation with small differences for small mirrors. That effect will be neglected in...
this work, and the focus equation will be assumed to apply to
the mirror pair, with the center of the pair taken as its location
along the optic axis. For a given magnification $M = q/p$, the
KB focus equation can be re-expressed as

$$ q = \sqrt{M(M+1)} Rdx, \quad (2) $$

where $dx$ is the offset of a single mirror pair perpendicular to
the $z$ axis.

For untilted mirror pairs, the images fall on a circle $r_{\text{circle}}$
given by

$$ r_{\text{circle}} = \sqrt{2} dx(M-1). \quad (3) $$

The images of the ideal framing-camera pattern have three
different offsets from the center of the pattern [Fig. 150.28(a)].
Four images are at the corners, eight images are on the sides,
and four images are at the center of the pattern. The amounts
that a mirror pair must be moved, $\Delta r_{\text{mirror}}$, and tilted in pitch,
$\Delta \sigma_{\text{mirror}}$, to move the image by $\Delta r_{\text{image}}$ are given by

$$ \Delta r_{\text{mirror}} = \Delta r_{\text{image}}/(M+1), \quad \text{(4)} $$

$$ \Delta \sigma_{\text{mirror}} = \Delta r_{\text{image}}/p. \quad \text{(5)} $$

The parameters of the compact KB mirrors used in this work
are given in Table 150.I. The angles $\phi$ that the mirror pairs
make with the axis of the framing camera and the mirror-pair
positions and tilts that generate the pattern of image positions
shown in Fig. 150.28(b) are provided in Table 150.II. Note that
to move the inner images sideways, the mirrors must be tilted
in roll $\Delta \beta_{\text{mirror}}$ by an amount

$$ \Delta \beta_{\text{mirror}} = \Delta x_{\text{image}}/2q, \quad \text{(6)} $$

where $\Delta x_{\text{image}}$ is the perpendicular amount to move the image. As an example, $\Delta x_1$ is shown in Fig. 150.28(a).

The mirror-pair alignment is accomplished by placing the
mirror-pair vertex at the offset positions given by the values
in Table 150.II with preimposed tilts in pitch and roll. This
was accomplished by using precision positioning stages and

Table 150.I: Parameters of compact mirrors used in
the assembly of the KBFRAMED optic.

| $R_{\text{KB}}$ | 27.5 m |
| $\Delta t$ | 4.5 mm |
| $\Delta x$ | 2.2 mm |
| $M$ | 12 |
| $q$ | 2173.2 mm |
| $p$ | 181.0 mm |
| $r_{\text{circle}}$ | 34.22 mm |
| $i$ | 0.696° |

Table 150.II: Mirror-pair offsets and tilts needed to generate the image locations in Fig. 150.28(a) with the pair assignments shown in Fig. 150.28(b). The remaining 12 pairs have common positions and tilts depending on image location as described in the text.

<table>
<thead>
<tr>
<th>Mirror Pair</th>
<th>$\phi$ (°)</th>
<th>$r_{\text{image}}$ (mm)</th>
<th>$\Delta r_{\text{image}}$ (mm)</th>
<th>$\Delta r_{\text{mirror}}$ (mm)</th>
<th>$\Delta \sigma_{\text{pitch}}$ (°)</th>
<th>$\Delta \beta_{\text{roll}}$ (°)</th>
</tr>
</thead>
<tbody>
<tr>
<td>16</td>
<td>−22.5</td>
<td>14.61</td>
<td>19.61</td>
<td>1.51</td>
<td>0.478</td>
<td>0</td>
</tr>
<tr>
<td>1</td>
<td>0</td>
<td>4.5</td>
<td>29.72</td>
<td>2.29</td>
<td>0.725</td>
<td>0.059</td>
</tr>
<tr>
<td>2</td>
<td>+22.5</td>
<td>14.61</td>
<td>19.61</td>
<td>1.51</td>
<td>0.478</td>
<td>0</td>
</tr>
<tr>
<td>3</td>
<td>+45</td>
<td>19.09</td>
<td>15.13</td>
<td>1.16</td>
<td>0.367</td>
<td>0</td>
</tr>
</tbody>
</table>
a rotary stage to position the base under a fixed, magnified viewing system (157× on a video display). Assembled mirror pairs with pre-applied, UV-curable epoxy on the optic base side were held in place over the base with a vacuum chuck that was positioned by a six-axis positioner (three axes of position and three of tilt). In this fashion the mirrors were cured into place with the UV epoxy acting as the tilted interface to the flat optic base. Positioner accuracies were 1/10,000th of an inch (2.54 µm), 0.01° in rotation of the optic base, and 2.36 arcsec in pitch and roll of the mirror pair.

All mirror-pair image positions were measured by placing the optic assembly in a vacuum system with a microscope chassis identical to that used with the framing camera and back-illuminating a grid co-aligned with the axis of the microscope (z axis) and at the focus distance for \( M = 12 \) (181 mm) with an e-beam–generated x-ray source. Exposures were taken using a Fuji image plate and image positions determined to 0.1 mm. Any inaccuracies in image positions were minimized by removing the mirror pair and correcting the tilt angles in pitch and roll. Final accuracies of mirror-pair alignments were \( \approx 5 \) µm in position and 20 arcsec in pitch and roll. This resulted in all images being within 1 mm of the center of the ideal framing-camera cathode strip pattern (i.e., spaced by 9 mm vertically).

The resolution of the mirror pairs at best focus and the off-axis aberrations are discussed in detail in Ref. 13, and the resolution is calculated ideally to be better than \( \approx 5 \) µm over a 400-µm-diam region around best focus. Tilting and repositioning the mirrors, ideally, avoids any additional blurring caused by misalignment from best focus; whereas, in practice, exact alignment is not possible and the framing camera will add additional blurring to the images. Therefore, it is better to determine the resolution by measurement. The inferred point-spread function (PSF), including blurring by the framing camera, is discussed in the next section.

The fused-silica compact KB mirror components are coated with 500 Å of Ir on top of a 150-Å Cr sticking layer as detailed in Ref. 13. The mean radius of curvature of the set of 32 mirrors used to assemble the 16 mirror pairs is 27.2 m, with a range from 25.6 to 28.6 m. The mirror pairs have radii of curvature that are typically within 0.1 m of each other. The x-ray reflectivity\textsuperscript{13} of the mirrors has been measured to approach an ideal reflector at the grazing angle of 0.7°. The typical sensitive energy band of the 16-image KB, calculated from the Henke-scattering factors,\textsuperscript{15} is shown in Fig. 150.29, including the transmission of the blast shield, vacuum window, and example filters. The sensitive band extends from \( \approx 2 \) keV to 8 keV.

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**The KBFRAMED Instrument**

Figure 150.30 shows a schematic of the KBFRAMED instrument. It consists of a chassis fixed in the OMEGA target chamber, the 16-image KB mirror assembly, and the vacuum interface to a high-speed framing camera.\textsuperscript{6} The mirror assembly is held such that the mirror-pair centers are 181 mm from target chamber center, having been set to a precision of 10 µm by a pointer placed on the optic cover when it was installed. A blast cover with holes aligned with the mirror-pair centers contains an x-ray–transparent Be foil that protects the mirrors from exposure to laser-generated target debris. A vacuum Be window separates the chamber vacuum from the path to the image plane, so a separate vacuum system provides a high vacuum (\( \approx 10^{-6} \) Torr) to the framing-camera active-detector.
region. This also isolates the camera from contaminants such as tritium from the targets. At present, the images are recorded on film that is not in the vacuum region of the framing camera, making it easy to exchange.

Figure 150.31 shows example images of a resolution grid taken by backlighting a grid placed at target chamber center by an Au foil placed 5 mm behind the grid. The foil is illuminated with 2 kJ of 351-nm UV light in a 1-ns pulse from six OMEGA beams. The grid (25.4-μm-diam Cu wires, spaced by 50.8 μm) is placed on a Ta foil with a 500-μm-diam hole, thereby producing 16 clearly separated images. The framing-camera images were recorded with Kodak T-MAX 3200 film and digitized on a calibrated PerkinElmer photo microdensitometer using 20-× 20-μm scan pixels. A step wedge was imposed on the film before exposure in the framing camera, which allowed the scanned film density to be converted to intensity. The framing camera adds blurring to the images with a scale of ∼50-μm full width at half maximum (FWHM) at the image plane, (i.e., ∼5 μm at the target plane). To estimate the effective blurring, a step pattern with the width and spacing of the Cu wires is convolved with a 2-D Gaussian blur function and then compared with the observed blurring. Figure 150.32 shows a lineout through a single intensity-corrected image taken through the central 200-μm-wide region, averaged 10 μm vertically to reduce noise. The measured pattern is compared to the Gaussian-blurred step pattern (dashed red curve) whose FWHM is 6 μm. The close agreement indicates that the Gaussian blur function is a good approximation to the net blurring of the framed, KB mirror-pair images.

Figure 150.32
Lineout through a single framed, backlit image obtained with KBFRAMED (image 10 of Fig. 150.31). The lineout (solid curve) through the central 200 μm is compared with an ideal grid pattern convolved with a 6-μm FWHM Gaussian blur function (dashed red curve).

Hot-Spot Evolution Imaged by KBFRAMED
KBFRAMED was developed principally to acquire time-resolved x-ray images of the cryogenic target implosion’s stagnation region (i.e., hot spot). Triggering of the framing camera is accomplished by electrical delay using a reference to the master oscillator of the OMEGA laser that is accurate to the picosecond level. Since the hot spot evolves very quickly in time (∼100 ps), the framing-camera strip times are set to differ by 15 ps from strip to strip by using timed cables whose pulse propagation time differs by this amount (to within ±2 ps, measured to ±1 ps). The relative time of an image is determined from these delays and from the distance of the image from the beginning of the strip, assuming a pulse propagation speed of c/2. Deviations from the above assumptions caused by cross talk between neighboring strips are assumed to be small for these small offsets in pulse arrival times. Absolute times can be assigned to data where the simultaneously measured time history of the neutron emission is measured by the neutron temporal diagnostic (NTD), it is assumed that the x-ray and neutron emissions peak at the same time. Figure 150.33 shows example images of a cryogenic target’s stagnation recorded by KBFRAMED with times so assigned from the beginning to the end of measurable core emission (the relative times are accurate to ∼2 ps, whereas the absolute time may be in error.
A FRAMED, 16-IMAGE KIRKPATRICK–BAEZ X-RAY MICROSCOPE

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by as much as ~50 ps because of uncertainties in the time of the peak of the measured x-ray flux and the absolute timing of NTD). Image signal levels were adjusted for gain as a function of position on the strip determined from measurements of a uniformly illuminated x-ray–emitting foil observed with the same framing camera and the same strip timings. In this experiment an 8.8-µm-thick deuterated polystyrene (CD) shell, 960 µm in diam, filled with DT cryogenically cooled to form a 57-µm-thick DT ice layer, was imploded with 29 kJ of UV (351 nm) from the 60 beams of the OMEGA Laser System, using a triple-picket pulse, having a 1.5-ns-long main pulse.¹⁹

The data were recorded with a 2-mil (50.8-µm) Al filter in front of the framing camera, so the energy band was ~4 to 8 keV (see Fig. 150.29). The emission is seen to start as a low-intensity diffuse emission in a region of ~50-µm diameter, brighten to a maximum in ~70 ps, and then decrease over the next 70 ps. Inferences of hot-spot pressures are made from the size of the hot spot measured by KBFRAMED, the time of fusion burn, the measured ion temperature, and the measured neutron yield.²⁰ Without every one of these measurements, including the high-spatial-resolution framed images provided by KBFRAMED, the inferences of hot-spot pressure would not be possible. Additionally, the structure evident in the images at scales comparable to the PSF (6-µm FWHM, as indicated by a circle of that size in Fig. 150.33) would not be observable without the resolution provided by KBFRAMED.

An example shape analysis of the hot-spot x-ray emission near the peak of the signal is shown in Fig. 150.34. The hot spot is first fit to a super-Gaussian ellipse convolved with the Gaussian point-spread function given by

\[ I'(x,y) = \text{PSF}(x,y) \otimes I_0 \exp \left\{ -\left[\left(\frac{x-x_c}{a}\right)^2 + \left(\frac{y-y_c}{b}\right)^2\right]^{\eta/2}\right\}, \]  

(7)

\[ x^* = (x-x_c)\cos\alpha + (y-y_c)\sin\alpha, \]

\[ y^* = (x-x_c)\sin\alpha - (y-y_c)\cos\alpha, \]

(8)

where \( \otimes \) denotes convolution, \( a \) and \( b \) are the lengths of the semi-major and semi-minor axes, respectively, \( I_0 \) is the peak value, and \( \eta \) is the super-Gaussian order. The values \( x^* \) and \( y^* \) are the coordinates lying along the major and minor axes of the ellipse, given by

Figure 150.34(a) shows the KBFRAMED image at the peak of the hot-spot emission. Figure 150.34(b) shows the best-fit, convolved super-Gaussian ellipse with \( a = 23.5 \mu m, b = 20.9 \mu m, \eta = 2.96, I_0 = 0.32, \) and \( \alpha = 91.4^\circ \). Figure 150.34(c) shows the difference, demonstrating that the fit accurately determines the size of the image with only small-scale structure and noise remaining. An example lineout through the image is shown in

**Figure 150.33**

KBFRAMED images of hot-spot x-ray emission from a cryogenic target implosion. The approximate point-spread function (PSF) (6-µm FWHM Gaussian) is indicated by a circle of that size in the first image.

**Figure 150.34**

A single KBFRAMED cryogenic target hot-spot image at x-ray maximum: (a) image with dashed line indicating direction of lineout, (b) convolved, super-Gaussian–ellipse fit to image, and (c) difference between (a) and (b).
A framed, 16-image Kirkpatrick–Baez X-ray microscope

Fig. 150.35 with the direction of the lineout indicated by the dashed line in Fig. 150.34(a). The need to use a fit is exemplified by the lineout, where it is evident that in order to estimate the average peak of the hot spot in the presence of noise in the image, it is necessary to use the best-fit value rather than a single peak value. The minimal difference in the convolved fitting function and the inferred super-Gaussian ellipse is because the emission is well resolved by the given resolution of KBFRAMED for this hot-spot size. However, since this method makes it possible to compare sizes when measured with differing resolutions, it is the preferred procedure. With the peak of the hot spot so defined, the size of the hot spot is then defined by the convention that the hot-spot radius is given by the average radius where the emission is 17% of the maximum. 21 With this definition, \( r_{17} \) is given by

\[
r_{17} = \left( \ln 0.17 \right)^{1/\eta} r_0,
\]

where \( r_0 \) is the geometric mean of \( a \) and \( b \left( r_0 = \sqrt{ab} \right) \). For the image above, \( r_{17} \) is found to be 26.9 \( \mu \text{m} \).

A more-detailed fit to the hot-spot envelope is determined by fitting the contour of the image at 17% of the fit peak to a Legendre polynomial with the axis of the fit taken as the semi-major axis of the super-Gaussian fit. Figure 150.36 shows the 17% contour, the Legendre fit to the contour (the two sides of the image are separately fit with the major axis of the super-Gaussian fit defining the sides), and the super-Gaussian–fit 17% contour on the image of Fig. 150.34(a). The fractional-radial deviation (departure from a circle) of the contours as a function of angle from the semi-major axis of the super-Gaussian–ellipse fit is plotted in Fig. 150.37. The Legendre modes of the fit are shown in Fig. 150.38 for modes from 2 to 10 (mode 1 is just a shift of the center) with the value taken as the average of the fits to the two sides of the contour and the error bar defined by the minimum and maximum of

![Figure 150.35](image)

Lineout through image in Fig. 150.34(a) and through the convolved super-Gaussian ellipse fit to that image (fit) and through the unconvolved fit (inferred fit).

![Figure 150.36](image)

Single KBFRAMED image from Fig. 150.34(a) with \( r_{17} \) contours superposed.

![Figure 150.37](image)

Fractional-radial deviation of \( r_{17} \) contours from Fig. 150.36 as a function of the angle from the semi-major axis of the super-Gaussian–ellipse fit.

![Figure 150.38](image)

Legendre-mode spectrum of the fit to the measured \( r_{17} \) contour in Fig. 150.36.
those two fits. The Legendre fit to the hot-spot envelope at $r_{17}$ is, as expected, closer to the observed shape, although the average radius differs only slightly from the elliptical fit (26.7 μm for the observed and Legendre fit as opposed to 26.9 μm for the elliptical fit). In this particular image, modes 2 through 5 are significant although all are less than 0.1 (i.e., less than 10%), whereas modes 6 through 10 are less than −2%. Note that with an emission region of this radius, mode 10 is expected to be suppressed by the resolution of the instrument by approximately a factor of 2, i.e., the true limit for mode 10 is less than −4% for an observed limit of −2%.

The dominant modes of the hot-spot envelope are those expected from on-target illumination nonuniformities coming from beam-intensity imbalance, but this observation does not determine that they are the source of the perturbations. Also, it is important to note that the major axis of the ellipse is within 2° of the vertical (91.4° best fit), which is approximately parallel to the direction of the stalk that holds the cryogenic target in place in the OMEGA target chamber (KBFRAMED is located 10° below the equator of the OMEGA chamber and the stalk direction is downward in the images). The stalk and the glue spot that binds the stalk to the CD shell that surrounds the DT ice layer are known to be the largest mass perturbation at the surface of the target. The effect of a stalk is complex in nature but, simply put, it causes the hot spot to become elongated in the direction of the stalk. This example serves to illustrate the benefit of the increased resolution of the KBFRAMED instrument and the type of information that can be obtained from these images.

Conclusions
A novel 16-image KB microscope design that couples to a high-speed framing camera has been implemented on the OMEGA Laser System. This instrument, known as KBFRAMED, obtains framed images of x-ray emission from laser-generated plasmas with ~6-μm spatial resolution, ~30-ps time resolution over a region of ~400 μm in the energy band from 2 to 8 keV. It was specifically designed to measure the stagnation region (hot spot) of cryogenically cooled DT target implosions that have typical sizes of ~60-μm diameter and durations of ~100 ps. The spatial resolution and time sampling of KBFRAMED allow one to measure the time-varying size and shape of these hot spots.

ACKNOWLEDGMENT
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REFERENCES


Modeling Tritium Interactions with Metals

Introduction
Quantitatively modeling the interaction of tritium with metals is a vital step toward understanding the mechanisms of tritium migration through the metal. In turn, understanding these fundamental mechanisms is necessary to interpret experimental results and to make accurate predictions regarding tritium migration in metals during exposures to tritium gas and subsequent storage periods and/or cleaning procedures. While the literature contains several attempts at creating a quantitative model,1–4 each attempt lacks one fundamental aspect: tritium migration across the surface–metal lattice interface. Many reports in the literature show that tritium adsorbs onto the surface as tritiated water,1,5,6 while tritium absorbs into the metal lattice as atomic tritium.5 This difference in retention media results in the measured large differences in tritium concentrations in the surface and in the bulk metal.7 Including this physical condition into a quantitative model is necessary to accurately model the complete tritium–metal system.

In this article, we present a quantitative tritium migration model (QTRIMM) for modeling tritium migration in various metal substrates. This model includes the surface film of adsorbed water and relates the concentrations of tritium within the thin film of adsorbed water to the tritium concentrations within the substrate metal. Additionally, the tritium concentrations throughout the metal sample are output from the model calculation. Currently, this information is obtained experimentally only through acid etching or other destructive techniques. The model developed in this work provides two major advantages. First, inclusion of the thin film of adsorbed water in the model provides the first step toward a global model, capable of describing all experimental conditions. Second, because the tritium concentration profiles throughout the sample are calculated, the model has the potential to predict the location of tritium within a metal using only the loading and storage conditions. This avoids the necessity of destructive techniques to determine the tritium concentration profiles. Further, the model can be used to calculate the increase in gas-phase protium concentrations in a mixture of deuterium and tritium.

In this article the physical picture of a metal’s surface is presented first, along with the primary assumptions and the relevant equations. Following this, a detailed derivation of the model is included. This derivation is divided into two major parts: bulk migration and surface conditions. To model tritium migration in the bulk metal lattice, Fick’s second law of diffusion is solved numerically. This solution includes a condition for tritium diffusion through composite media. Next, the surface boundary conditions used to model tritium migration during an exposure to tritium gas and during a subsequent storage period are presented. Finally, a few predictions are made by using QTRIMM.

Surface–Metal Interface Condition
Tritium interacts with metal substrates by first adsorbing onto the metal surface and then permeating through the metal lattice. Under most conditions, tritium does not adsorb directly onto the metal’s surface. Instead, the tritons adsorb in the form of tritiated light water (HTO).1,5,6 After adsorption onto the surface, tritium can migrate from the water layers into the bulk metal lattice. The tritons occupy interstitial locations and defect sites within the metal’s crystal lattice and diffuse through the lattice by migrating between the various sites. This process is illustrated in Fig. 150.39. Here, the approximate thickness of both the water layers and the metal oxide layers are shown for reference. The remainder of this section will address the tritium migration mechanisms treated in QTRIMM and will discuss the major assumptions of the model.

The first step in the tritium permeation process is adsorption onto the metal surface. This process presumably occurs through isotope exchange between tritiated species in the gas phase and water molecules adsorbed on the metal surface.6 This exchange process is expected to occur rapidly, relative to tritium migration into the underlying metal lattice. Assumption of rapid equilibrium across the surface–gas interface is justified by comparing the flux of gaseous tritium to the surface and the diffusive flux of tritium into the metal lattice.8 Additionally, the average residence time of adsorbed species is of the order of femtoseconds at 25°C.
Rapid equilibrium across the surface–gas interface makes it possible to determine the surface concentration of tritium. At equilibrium, the surface concentration of tritium is related to the concentration of tritium within the gas phase.

\[ c_{\text{surf}}^{\text{eq}} = f \cdot c_{\text{gas}}^{\text{eq}}. \]  

In this formula, the mole fraction of tritium on the surface \( (c_{\text{surf}}^{\text{eq}}) \) is related to the mole fraction of tritium in the gas phase \( (c_{\text{gas}}^{\text{eq}}) \) by a constant factor \( (f) \). This factor accounts for various isotope exchange probabilities. In the limiting case of equal exchange probabilities, the scaling factor is unity. In reality, it is likely that the formation of double-isotope species, such as \( \text{T}_2\text{O} \), is not as probable as the formation of mixed isotope species of water, such as HTO or DTO (tritiated heavy water). This more-realistic scenario of nonequal reaction probabilities would reduce the scaling factor to a fractional value ranging between zero and unity. For simplicity, we take \( f = 1 \) in our calculations.

Using the above relation between the mole fractions in the two phases makes it possible to determine the absolute concentration of tritium in the adsorbed water layers. Assuming the density of these water layers does not change significantly with each successive layer, the concentration of tritium on the metal’s surface is given by

\[ c_{\text{eq}} = \frac{\rho_{\text{H}_2\text{O}} \cdot f \cdot c_{\text{surf}}^{\text{eq}} \cdot \delta_{\text{ML}}}{2 \text{ mol H} \cdot \text{mol H}_2\text{O}} \]  

Here, the concentration of tritium \( (c_{\text{surf}}^{\text{eq}}) \) is determined by the mole fraction of tritium in the gas phase \( (c_{\text{gas}}^{\text{eq}}) \), the isotope exchange scaling factor \( (f) \), the surface density of water \( (\rho_{\text{H}_2\text{O}}) \), and the thickness of a monolayer of water \( (\delta_{\text{ML}}) \). An additional factor of 2 is included to relate the number of moles of hydrogen to the moles of water on the surface. In the limit of \( f = 1 \), \( c_{\text{surf}}^{\text{eq}} \) is taken to be equal to a saturated surface solubility \( S_{\text{surf}} \). Under the assumptions of rapid equilibrium and a static gas phase, the equilibrium surface concentration in Eq. (2) remains constant throughout an exposure to tritium gas. In a situation where the fraction of tritium in the gas phase changes, the surface concentration will rapidly adjust to the new conditions.

After adsorption onto the surface, tritium can then permeate through the metal lattice by diffusing from site to site within the lattice. Permeation through the lattice begins with the triton crossing the surface–metal lattice interface. Migration across this interface is also assumed to be much faster than the rate of tritium diffusion in the underlying metal substrate. As such, the chemical potentials of tritium dissolved on either side of the interface are equal at the interface:

\[ \mu_{\text{surf}} (x_1) = \mu_{\text{metal}} (x_1), \]  

where \( \mu_{\text{surf}} \) and \( \mu_{\text{metal}} \) are the chemical potentials in the adsorbed water layers and the metal lattice, respectively, and \( x_1 \) represents the position of the interface. This equality leads to a relation between the concentrations of tritium within each region:

\[ \frac{c_{\text{surf}}^{\text{eq}} (x_1)}{c_{\text{metal}}^{\text{eq}} (x_1)} = \exp \left( -\frac{\mu_{\text{surf}} - \mu_{\text{metal}}}{RT} \right). \]  

This equation states that the ratio between the equilibrium concentrations in each region \( (c_{\text{eq}}) \) depends on the standard chemical potentials in each region \( (\mu') \), the gas constant \( (R) \), and the temperature \( (T) \). In a complete description of solute migration across an interface, the standard chemical potentials can vary from point to point as a function of depth. Varying standard potentials may result in gradual changes in the potential, culminating with equal standard potentials at the interface. Equal standard potentials result in equal concentrations at the interface.
interface and therefore in a continuous solute concentration profile across the interface region. Including spatially variable standard chemical potentials may not be necessary since some physical situations can be modeled well with the simplified approach of constant chemical potentials. Physical situations like this may be the result of well-defined boundaries between phases, similar to the interface between two solids. In these situations, the standard potential changes over a small distance are small and the equilibrium tritium concentrations at the interface are related by a constant factor.

The constant factor relating the tritium concentrations at the interface can be obtained by using Sievert’s law, where the ratio of the tritium concentrations given by Eq. (4) is equal to the ratio of the solubilities of tritium in each region \( S_i \), assuming the partial pressure of tritium at the interface is a constant:

\[
\frac{c_{\text{surf}}}{c_{\text{metal}}} = \frac{S_{\text{surf}}}{S_{\text{metal}}},
\]

This equation makes it possible to determine the tritium concentration in the metal lattice, given a tritium concentration on the surface. In order for the equality shown in Eq. (5) to hold, the metal lattice concentration must be multiplied by the isotope exchange scaling factor [Eq. (1)]:

\[
\frac{c_{\text{surf}}(f)}{f \cdot c_{\text{metal}}} = \frac{S_{\text{surf}}}{S_{\text{metal}}},
\]

For a constant temperature, the ratio of tritium solubilities in each region is a constant. Therefore, if the isotope exchange scaling factor is less than unity, the surface concentration will decrease, which means the concentration in the metal lattice must also decrease by the same factor.

This formalism for tritium migration across a boundary applies not only to transport across the adsorbed water–metal lattice interface but to any well-defined interface. One can extend Eqs. (3)–(6) to other composite media such as gold-plated metals or to metal substrates with artificially grown metal oxide layers, for example.

The final step toward describing tritium permeation is tritium diffusion through the metal lattice. For simplicity, we treat only lattice diffusion and ignore pathways such as grain boundaries and triple junctions and trap sites such as dislocations and vacancies. While these can be notable defects, the majority of the dissolved tritium is expected to migrate through the interstitial lattice sites in well-behaved metals with negligible contributions from defect sites.

**Modeling Tritium Migration in the Bulk Metal Lattice**

A model based on Fickian diffusion of tritium through the metal lattice has been developed. This model numerically solves for the tritium concentrations throughout a metal and includes a condition for solute diffusion through different, but interacting media, as outlined in the previous section. The numeric solution uses an intermediate time step and divides the solid into \( N \) cells. For simplicity, the model uses a one-dimensional solution. The quantity of solute in each cell can be determined by multiplying the cell thickness (\( \Delta x \)) with the surface area of the sample used in an experiment.

The concentration in each cell \( c \) is determined from a flux balance of tritium entering and leaving each cell. Because of this, the flux \( F \) is calculated at the intermediate positions, \( i \pm 1/2 \) as shown in Fig. 150.40. The rate of change in the solute concentration in each cell is determined by relating Fick’s First and Second Laws to yield

\[
\frac{\partial c}{\partial t} = -\frac{\partial F}{\partial x}.
\]

The finite-difference fluxes at the intermediate positions are

\[
F_{i+1/2} = -D_{i+1/2} \left( \frac{c_{i+1} - c_i}{x_{i+1} - x_i} \right)
\]

![Division of solid into equally spaced cells.](E26106JR.png)
where $D$ is the diffusivity. Using Eqs. (7)–(10), the rate of change of the concentration in cell “$i$” is written as

$$\frac{c_i^t - c_i}{\Delta t} = -\theta \left( \frac{F_{i+1} - F_i}{\Delta x_i} - \frac{(1-\theta)(F_{i+1} - F_{i-1})}{\Delta x_i} \right), \quad (10)$$

where the primes denote the next time step and $\theta$ is the degree of implicitness, which determines the degree of accuracy and stability; $\theta = 0.5$ provides the highest accuracy, while $\theta > 0.5$ enhances stability.

Equations (7)–(10) represent the basis for QTRIMM. Using these equations, a system of linear equations is obtained. Each equation in the set gives the concentration at the next time step for a particular location within the metal. Two additional equations are necessary to solve for the boundary conditions. These equations occur at indices of $i = 0$ and $i = N$ and will be discussed later. The total set of equations has one degree of freedom, allowing the system of equations to be solved exactly. Coupling Eq. (10) with a set of boundary conditions makes it possible to calculate the tritium concentrations in a homogeneous solid.

Tritium diffusion across an interface between two media requires that the flux equations [Eqs. (8) and (9)] be constrained. These fluxes must be modified to maintain a constant ratio of the tritium solubilities across the interface [Eq. (5)]. The other implicit condition is that the total quantity of tritium throughout the system must be conserved. Using these conditions, the equations for the modified boundary conditions can be derived.

To maintain mass balance throughout the system, the concentration at the interface must be given by the rate of solute entering and exiting the cell, i.e., we require flux balance. In the present diffusion model, the interface between two well-defined solvents is placed within one cell ($i = M$) in the discretized solid. The concentration at the interface position is then given by the average of the concentrations on the right ($c_M^R$) and left ($c_M^L$) sides of the cell because half of this cell is one solvent and half is the other:

$$c_M = \frac{(c_M^R + c_M^L)}{2}. \quad (11)$$

Combining Eqs. (5) and (11) yields

$$c_M^R = \left(\frac{2}{S_R+1}\right)c_M = \varphi_R * c_M. \quad (12)$$

Similarly,

$$c_M^L = \left(\frac{2 * S_R}{S_R+1}\right)c_M = \varphi_L * c_M. \quad (13)$$

The flux in and out of the interface cell must utilize $c_M^L$ and $c_M^R$:

$$F_{M+1/2} = -D_{M+1/2} * \left(\frac{c_M^{R+1} - c_M}{x_{M+1} - x_M}\right)$$

$$= -D_{M+1/2} * \left(\frac{c_M^{R+1} - \varphi_R * c_M}{x_{M+1} - x_M}\right), \quad (14)$$

$$F_{M-1/2} = -D_{M-1/2} * \left(\frac{c_M^L - c_M}{x_M - x_{M-1}}\right)$$

$$= -D_{M-1/2} * \left(\frac{\varphi_L * c_M - c_M^L}{x_M - x_{M-1}}\right). \quad (15)$$

Equations (14) and (15) can be inserted into Eq. (10) to determine the concentrations around the interface ($i = M-1, M, M+1$). These equations apply only around the interface position; the remaining equations are unchanged from the form derived for a homogeneous solid.

To solve for the concentration profiles in a metal sample, two equations for the boundary conditions are necessary. For the first boundary condition, we assume a symmetric solid. Under this assumption, the diffusion model must extend only to the center of the sample; the other side of the sample is a mirror image. To solve for the concentration at the centerline, we set the fluxes into or out of this cell to be equal but with opposite signs to reflect the opposite directions of flow. The boundary condition for the sample’s surface depends on the experimental conditions. Two general cases are outlined in the following section: the first case treats a storage condition where tritium is allowed to redistribute throughout the solid as well as to desorb from the surface; the second case treats a
condition where a metal is exposed to gaseous tritium and is loaded with tritium.

Boundary Conditions at the Metal’s Surface

We utilize slightly different surface boundary conditions, depending on whether the model calculates the concentration profiles during the loading phase or during a storage phase. During the loading phase, the surface concentration is assumed to be constant. During the storage phase, this concentration is allowed to vary because tritium is allowed to migrate into the metal lattice and to desorb from the surface at a constant rate. The derivations for these two conditions are shown below.

In all cases, multiple monolayers of adsorbed water develop on all metal surfaces that have been exposed to a humid atmosphere. These layers are distinct from the bulk metal and have a much higher solubility for tritium. To model tritium migrating from the bulk metal into the surface, the interface equations outlined in the previous section are used. Assuming that a rapid equilibrium develops across the surface–metal lattice interface, the concentrations on either side of the interface are related by the ratio of the solubilities for tritium in each region, as presented in Eq. (5).

The thickness of the metal substrate is much larger than the thickness of the adsorbed water layer as illustrated in Fig. 150.39. To account for this significant difference in thickness and to reduce the calculation time, two different cell sizes are used: one to calculate the concentrations in the bulk of the sample (Δx) and the other to calculate the surface concentration (δx). Given the small thickness of the water layer and the relatively slow rate of diffusion into the metal lattice, tritium concentrations likely equilibrate rapidly in the water layers relative to the bulk metal. Assuming this rapid equilibration and using the small water-layer thickness, the surface–metal interface is placed entirely in the first cell (i = 0) of the discretized solid as shown in Fig. 150.41. The thickness of the surface cell (i = 0) is determined by the surface concentration of adsorbed water (Q), the surface density of water (ρH₂O), and the thickness of a monolayer of water (δML):

\[ δx = \frac{Q}{ρ_{H₂O}} \cdot δ_{ML} \cdot 2. \]  

The factor of 2 in this equation scales up the thickness of this cell to account for the fact that half of this cell is adsorbed water and half is the bulk metal lattice (Fig. 150.41).

For the loading phase, the surface concentration is assumed to remain constant in time. To incorporate this fixed concentration into the solution matrix, we assume a linear relation between the concentrations in the first three cells:

\[ \frac{c'_1 - c_0 ϕ_R}{x'_1 - x_0} = \frac{c'_2 - c_0 ϕ_R}{x'_2 - x_0}. \]  

This equation follows the same form as Eq. (10) and can be inserted directly into the system of equations defined by Eq. (10).

During the storage phase, the tritium concentrations in the surface cell (i = 0) are allowed to vary by including two conditions for tritium transport: diffusion into the metal and desorption from the surface. Following the same formalism shown above, we define the diffusive flux into the metal lattice as

\[ F_{1/2} = -D \left( \frac{c_1 - c_0 ϕ_R}{x_1 - x_0} \right). \]  

The desorbing flux away from the surface is defined by

\[ F_{-1/2} = -v \cdot c_0 \cdot ϕ_L \left( \frac{δx}{2} \right). \]  

where \( ϕ \) is the desorption rate constant. The return flux from the atmosphere to the surface is ignored in the equation because the airborne concentration is assumed to be negligible. The
thickness of the surface cell, $\delta x$, is divided in half since half of this cell represents the adsorbed water layers illustrated in Fig. 150.41. The desorption rate constant $\nu$ determines the rate at which tritium desorbs from the surface. Sureau and McElroy measured this rate to be 0.91% per day for untreated stainless-steel surfaces. This is equivalent to $10^{-7}$ tritons desorbing per second, which is the value used in the following calculations. Equations (18) and (19) can be inserted into Eq. (10) to yield the surface boundary condition in the diffusion model.

**Model Predictions**

Using the model outlined in the previous sections, we can predict the migration of tritium through a metal sample during an exposure to tritium gas, a subsequent storage period, and a decontamination procedure. The results of a series of simulations are presented in Figs. 150.42–150.44, which show the consequences of exposing stainless steel to tritium gas and then storing the metal for a period of time.

Table 150.III lists the hydrogen diffusivity in stainless steel, copper, and aluminum; Table 150.IV lists the solubilities for these materials. Figure 150.42 shows the calculated tritium concentration profiles that develop within a stainless steel sample, exposed to 612 Torr of hydrogen gas containing 60% tritium for 24 h. Initially, the surface concentration was fixed at the value shown, while the remaining concentrations in the metal were set to zero. The profiles are plotted in increments of 1 h. For clarity, only the concentrations within the metal lattice are shown; the high surface concentrations are off-scale. Using the final profile, we calculate that tritium penetrates a mean distance of $\sim 10 \mu m$ into stainless steel. This compares favorably with the expected mean migration distance, which is found from the semi-infinite solution to the diffusion equation. The expected mean migration distance is

$$\langle x \rangle = \sqrt{4Dt} = 11.5 \mu m,$$

where $\langle x \rangle$ is the mean migration distance, $D$ in $m^2/s$ is the diffusivity of tritium in stainless steel, and the exposure time $t$ is 24 h.

### Table 150.III: Hydrogen diffusivity in select materials.

<table>
<thead>
<tr>
<th>Material</th>
<th>Stainless steel</th>
<th>Copper</th>
<th>Aluminum</th>
</tr>
</thead>
<tbody>
<tr>
<td>Frequency factor ($m^2/s$)</td>
<td>$7.2 \times 10^{-7}$</td>
<td>$7.9 \times 10^{-7}$</td>
<td>$1.46 \times 10^{-6}$</td>
</tr>
<tr>
<td>Activation energy (kJ/mol)</td>
<td>52.9</td>
<td>38.6</td>
<td>30.0</td>
</tr>
<tr>
<td>Diffusivity at 25°C ($m^2/s$)</td>
<td>$4.1 \times 10^{-16}$</td>
<td>$1.4 \times 10^{-13}$</td>
<td>$5.0 \times 10^{-12}$</td>
</tr>
</tbody>
</table>

### Table 150.IV: Hydrogen solubility in select materials.

<table>
<thead>
<tr>
<th>Material</th>
<th>Stainless steel</th>
<th>Copper</th>
<th>Aluminum</th>
</tr>
</thead>
<tbody>
<tr>
<td>Frequency factor ($mol/m^3 \cdot atm^{1/2}$)</td>
<td>342</td>
<td>1691</td>
<td>4416</td>
</tr>
<tr>
<td>Activation energy (kJ/mol)</td>
<td>13.0</td>
<td>39.3</td>
<td>28.5</td>
</tr>
<tr>
<td>Solubility at 25°C ($mol/m^3$)</td>
<td>1.8</td>
<td>0.22*</td>
<td>$4.4 \times 10^{-2}$</td>
</tr>
</tbody>
</table>

*Scaled by 1000.

In the second phase, the tritium profile evolution is tracked over a storage period of 50 days. During this time, tritium is allowed to redistribute throughout the metal, as well as to desorb from the surface. Tritium that has desorbed from the surface does not return to the sample in this example. The final concentration profile calculated for the loading procedure illustrated in Fig. 150.42 was used as the initial condition for the storage phase. The resulting concentration profiles in the metal lattice are shown in Fig. 150.43. Again, the high surface concentrations are not shown. Finally, the evolution of the concentration profiles are plotted in one-day increments. As expected the tritium concentrations in the near-surface region decrease while the tritium concentrations in the deeper regions increase with increasing storage time. This trend in the concentration profiles with increasing time produces profiles similar to those
During the storage periods, tritium could diffuse deeper into the sample as well as desorb from the surface. During loading, the surface concentration was fixed by the given exposure conditions and the tritium content in the sample bulk was set to zero. For reference, the 1000-DPM/dm$^2$ threshold is shown as a dashed line in each of the following two plots, where DPM is disintegrations per minute.

Figure 150.45 shows the activity on the metal surface immediately prior to the cleaning. The exposure and cleaning sequence indicates that tritium migrates back to the surface within each 24-h storage period. Further, regardless of the initial exposure conditions, the surface activity present after the first 24-h period indicates that the surface is above the 1000-DPM/dm$^2$ threshold. In all exposure cases considered, one additional decontamination was sufficient to bring the surface below the threshold. Additional decontaminations performed after the second day showed no significant depletion of the surface activity. A sufficient quantity of tritium is present within the metal lattice to replenish the surface with tritium to nearly the same level each following day.

During the storage periods, tritium could diffuse deeper into the sample as well as desorb from the surface. During loading, the surface concentration was fixed by the given exposure conditions and the tritium content in the sample bulk was set to zero. For reference, the 1000-DPM/dm$^2$ threshold is shown as a dashed line in each of the following two plots, where DPM is disintegrations per minute.

Figure 150.44 shows the dependence of relative distribution of tritium in stainless steel on storage time following an exposure to tritium. Included in this plot are the relative quantities of tritium that desorbed from the sample, as well as those remaining on the surface and in the bulk metal lattice. $P = 612$ Torr, $X_{\text{tritium}} = 0.6$, $t = 24$ h, $Q = 10^{-5}$ mol/m$^2$, $\nu = 1 \times 10^{-7}$/s.

These authors measured a local minimum in the concentration profile near the surface. They attributed this minimum to the chronic desorption of tritium from the surface, which is an effect predicted by the model.

Figure 150.44 shows the calculated relative distribution of tritium within the stainless-steel sample over the same storage period represented by the concentration profiles in Fig. 150.43. This figure shows the relative tritium inventories contained on the surface and in the bulk metal lattice along with the relative quantity that desorbed from the surface. For reference, the total mass of the system is also included. From these results, we can see that the surface contained $\sim 60\%$ of the total tritium inventory immediately after the 24-h exposure. This quantity rapidly decreases over the first ten days since tritium not only diffuses into the sample but also desorbs from the sample’s surface. After $\sim 30$ days of storage, the relative distribution does not change significantly because the concentration gradients in the sample illustrated in Fig. 150.41 are less steep.

In the following example, a stainless-steel sample was exposed to various atmospheres of tritium gas for differing periods of time and then stored for either zero or one day prior to decontamination. The cleaning protocol was repeated over a ten-day period. In these two examples, the surface activity was reduced to zero during decontamination and the sample was stored for one day before repeating the cleaning procedure.

In the following example, a stainless-steel sample was exposed to various atmospheres of tritium gas for differing periods of time and then stored for either zero or one day prior to decontamination. The cleaning protocol was repeated over a ten-day period. In these two examples, the surface activity was reduced to zero during decontamination and the sample was stored for one day before repeating the cleaning procedure.
for 24 h after exposure and before the first decontamination. In general, the results show that storing the samples for one day results in surfaces containing higher quantities of tritium after each decontamination cycle. Compared with the samples that were decontaminated immediately after exposure, most of the stored samples did not reach the target threshold, even after ten decay cycles. Storing the sample after exposure gave the tritium more time to diffuse into the bulk metal lattice. This higher reservoir of tritium in the metal in turn resupplied the surface with tritium, leading to the higher surface activities.

The surface activity of aluminum after exposure to tritium and a series of subsequent decontaminations was also simulated. The operating conditions were identical to those used for stainless steel. As in the stainless-steel cases, two main scenarios were considered. In both scenarios (Fig. 150.47), the metal was exposed to a range of tritium gas concentrations for a variety of durations. In the first set of cases evaluated, decontamination proceeded immediately after exposure. In the second case, decontamination was initiated one day after the tritium exposure. In both cases, once the decontamination-dwell sequence was started, the surface was decontaminated after a 24-h hiatus over ten days.

The decontamination simulations indicate that more decontamination cycles are required for aluminum that was stored for one day prior to starting the decontamination sequence. The surface activity for the stored samples does not drop below the threshold until a minimum of seven decontamination cycles have been performed, compared to a minimum of two cycles for the case when the decontamination starts immediately after exposure. As in the stainless-steel case, the one-day storage period permitted tritium to permeate deeper into the bulk metal lattice. This leads to an increased reservoir of tritium in the lattice, which subsequently migrates to the surface after each successive cleaning.

For the same exposure conditions, stainless steel requires more decontamination cycles than aluminum when the exposed samples are stored for one day before starting the decontamination sequence. This difference is attributed to the fact that the diffusion of tritium into steel is $2 \times 10^4$ slower compared to aluminum. The higher diffusivity of tritium into aluminum expedites tritium migration to the surface after each cleaning. As a result, a greater fraction of the tritium inventory is removed from the bulk with each successive decontamination and the surface activity remaining after each 24-h period decreases faster on aluminum than on stainless steel.

The tritium exposure and decontamination protocols described above were repeated using 3-cm-thick aluminum instead of a 0.3-cm-thick sample (Fig. 150.48).
Figure 150.47
Decrease in the surface activity on 0.3-cm-thick aluminum calculated for a range of exposure conditions for two cases. (a) Decontamination was initiated immediately after exposure. (b) Decontamination was initiated one day after the exposure. The surfaces were decontaminated once every 24 h over a ten-day period in both cases. $Q = 10^{-5}$ mol/m$^2$, $\nu = 1 \times 10^{-7}$/s.

Figure 150.48
Decrease in the surface activity on 3-cm-thick aluminum calculated for a range of exposure conditions for two cases. (a) Decontamination was initiated immediately after exposure. (b) Decontamination was initiated one day after the exposure. The surfaces were decontaminated once every 24 h over a ten-day period in both cases. $Q = 10^{-5}$ mol/m$^2$, $\nu = 1 \times 10^{-7}$/s.
The calculations show that increasing the thickness of the aluminum results in higher surface activities throughout the decontamination process. For the immediate decontamination scenario, the surface activity remains above the target threshold for the higher tritium concentrations and longer exposure condition and drops below the threshold for the shorter periods and/or lower concentrations. In the delayed decontamination scenario, the residual surface activity does not drop below the threshold for any of the exposure cases considered. Additionally, the decay in the residual surface activity in both scenarios does not follow the trend observed for the 0.3-cm case illustrated in Fig. 150.47. In the 0.3-cm case, the residual surface activity decreased linearly after the second surface cleaning. In the 3-cm case, the residual surface activity decreases at a slower, nonlinear rate. The increased metal volume provides a larger reservoir for tritium in the 3-cm case for identical exposure conditions. Increasing the thickness of the aluminum allows tritium to migrate deeper into the bulk during the dwell periods between surface decontaminations because of the higher diffusivity of tritium in aluminum.

In the following example, a stainless-steel sample was exposed to a tritium atmosphere containing $5 \, \text{Ci/m}^3$ until a steady state was reached. Afterward, the sample was subjected to the same decontamination cycles described above. The surface was decontaminated once a day for ten days where each decontamination was assumed to completely remove all surface-bound tritium. During the storage period between the surface cleanings, tritium repopulated the surface and could desorb from the surface. The results are shown in Figs. 150.49 and 150.50. Figure 150.49 shows the tritium concentration profiles in the bulk metal after each decontamination, while Fig. 150.50 shows the surface activity remaining prior to the next decontamination cycle.

Figure 150.49 shows that each successive decontamination further depletes the tritium content in the near-surface region. This depletion is in response to the removal of the surface activity and, consequently, the abrupt change in the tritium inventory at the interface on the metal side. The concentration gradient drives tritium to migrate out of the metal lattice and back into the surface layer. Storing the sample for one day between successive decontamination cycles provides the necessary time for tritium to migrate into the surface water layer. Increasing the storage time would increase the quantity of tritium migrating to the surface since tritium diffusion through the metal lattice limits the flux to the surface.

In addition to depleting the near-surface tritium with each decontamination, the simulation shows that the surface concentration cannot reach the 1000-DPM/dm$^2$ threshold within ten decontamination cycles. Figure 150.50 shows the evolution...
of surface activities with sequential decontaminations. While the decontamination cycles reduce the residual surface activity to lower values, progress to puncture the 1000-DPM/dm² threshold is slow. The difficulty in reducing surface contamination is because of the vast reservoir of tritium present within the metal lattice.

QTRIMM can also be used to calculate the temporal evolution of tritium pulses applied to various metals. In the example under consideration, one side of a metal wall is exposed to a deuterium–tritium (DT) gas mixture at room temperature. Following the exposure, the DT gas is evacuated and the wall remains under vacuum for a specified period of time. This cycle of exposure followed by vacuum outgassing is repeated several times. The tritium concentration profiles and tritium distributions are calculated for each cycle. In these calculations, the downstream boundary condition is modified to include a high-solubility surface from which tritium can desorb. This condition is identical to that used for tritium desorption from a surface in Eq. (19) except that the flux is positive since the flow is in the opposite direction. The permeation calculations compare the performance of three metals: aluminum, copper, and stainless steel using the literature survey averages of the solubilities and diffusivities for the three metals listed in Tables 150.III and 150.IV. The solubility of tritium in copper was increased by a factor of 1000 relative to the literature survey average to account for the observed increase in hydrogen solubility in copper at low temperatures. The parameters listed in Table 150.V are used for this example for the three metals. We cannot assume that the quantity of water absorbed on the metal surfaces is the same for the three metals because these quantities are metal dependent. Additionally, the desorption rates from the upstream and downstream surfaces may be different because of the different environments; however, these parameters are not expected to differ significantly from metal to metal.

The response of the metals to the tritium pulses are shown in Figs. 150.51–150.53. In each figure (a) shows the calculated

![Figure 150.51](E26146JR)

(a) Activity distribution and (b) tritium concentration profiles calculated for a stainless-steel wall exposed to repeated cycles of DT gas followed by vacuum desorption.

<table>
<thead>
<tr>
<th>Table 150.V: Parameters used to calculate tritium permeation through the selected metals.</th>
</tr>
</thead>
<tbody>
<tr>
<td>Mole fraction of tritium in the gas</td>
</tr>
<tr>
<td>DT pressure</td>
</tr>
<tr>
<td>Temperature</td>
</tr>
<tr>
<td>Quantity of adsorbed water on upstream (DT gas/vacuum) wall</td>
</tr>
<tr>
<td>Quantity of adsorbed water on downstream (lab air) wall</td>
</tr>
<tr>
<td>Desorption rate constant (both surfaces)</td>
</tr>
<tr>
<td>Exposure time</td>
</tr>
<tr>
<td>“Storage” time</td>
</tr>
<tr>
<td>Wall thickness</td>
</tr>
<tr>
<td>Number of exposure/vacuum cycles</td>
</tr>
</tbody>
</table>
Figure 150.52
(a) Activity distribution and (b) tritium concentration profiles calculated for a copper wall exposed to repeated cycles of DT gas, followed by vacuum desorption.

Figure 150.53
(a) Activity distribution and (b) tritium concentration profiles calculated for an aluminum wall exposed to repeated cycles of DT gas, followed by vacuum desorption.
Modeling Tritium Interactions with Metals

The simulations show several notable differences among the three metals. The first difference is evident in the calculated concentration profiles within each metal. The profiles indicate that tritium permeates through aluminum much faster than the other two metals, such that steady-state permeation is achieved within the first DT exposure. Additionally, while the subsequent vacuum portions of each cycle serve to decrease the concentrations throughout the aluminum wall, steady-state permeation is reclaimed during the following DT exposure. For copper, a negligible quantity of tritium permeates through the wall and steady-state permeation is not achieved during the five cycles. Finally, permeation through stainless steel is the slowest, with no tritium reaching the downstream surface. This trend in tritium permeation rates follows the same trend in the tritium diffusivity through each metal: tritium diffusivity is slowest in stainless steel and quickest in aluminum. As a consequence, the calculated quantity of tritium desorbing from the downstream side of the wall is largest for aluminum and smallest for stainless steel.

Conclusions

QTRIMM, outlined in this article, allows one to calculate the tritium concentration profiles within a metal sample. This model represents a novel approach to assessing the migration of tritium into, out of, and within a metal. It accounts for high concentrations of tritium on metal surfaces. The model predicts the evolution of the tritium concentration profiles that develop during an exposure to tritium gas, during subsequent storage periods, and during iterative decontamination cycles. This article illustrates the application of QTRIMM to show the tritium concentrations within stainless steel that was exposed to tritium gas and then stored for 50 days. Additionally, QTRIMM was used to predict the changes in surface activity as a result of decontamination cycles. The model demonstrates two well-known phenomena: (1) tritium can “reappear” on a decontaminated surface and (2) the longer one waits to clean a contaminated surface, the harder it is to decontaminate the metal. Finally, QTRIMM was used to predict the quantities of tritium that can permeate through aluminum, copper, and stainless steel. The calculations show that the greatest quantity of tritium permeated through aluminum compared to copper or stainless steel of the similar thicknesses. As a result of the quicker diffusion of tritium through aluminum, the model predicts that aluminum will contain the largest quantity of tritium after several exposures to DT gas compared to the other two metals but it is expected to be decontaminated more quickly than the other two metals.

ACKNOWLEDGMENT

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REFERENCES

Three-Dimensional Modeling of Neutron-Based Diagnostics to Infer Plasma Conditions in Cryogenic Inertial Confinement Fusion Implosions

Introduction

In direct-drive cryogenic inertial confinement fusion (ICF) implosions, a target comprising a shell of cryogenic deuterium–tritium (DT) fuel enclosing a vapor region is irradiated using multiple nearly identical laser beams. As the kinetic energy of the imploding shell is converted to the thermal energy of the hot spot, the shell undergoes deceleration and conditions relevant for achieving fusion reactions are obtained. Conditions achieved in the compressed core, at the time of peak fusion neutron production, from a typical cryogenic direct-drive implosion on the OMEGA Laser System are shown in Fig. 150.54. The profile was obtained from a spherically symmetric simulation using the code LILAC; nonuniformity was ignored in this calculation. Ideally, a hot core is surrounded by a high-density shell, although multidimensional simulations indicate that while this is largely true, the neutron-producing region is typically not centered and the high-density shell can be significantly perturbed. Perturbations are typically quantified in terms of spherical harmonics (from 3-D simulations) or Legendre modes (in 2-D experimental images or 2-D simulations).

The goal of cryogenic implosions on OMEGA is to understand the physics of directly driven ICF implosions. This requires observations that could potentially shed light on failure mechanisms including the role of long-wavelength asymmetries on target performance. Three-dimensional (3-D) direct-drive simulations that include the effect of long-wavelength asymmetries performed with the arbitrary Lagrangian–Eulerian (ALE) code HYDRA indicate that significant long-wavelength asymmetries corresponding to spherical harmonics up to 4 should be present during the time of neutron production. These asymmetries can be seeded by power imbalance among the 60 beams of OMEGA; by beam mistiming, beam mispointing, or nonuniformities on the inside of the cryogenic layer at the ice–vapor interface; or by the initial error in the placement of the target relative to the center of the target chamber, etc. Long-wavelength asymmetries can compromise performance by reducing the clean volume over which neutrons are produced. In Fig. 150.54, the clean volume has a radius of ~20 μm. The Rayleigh–Taylor growth of the nonuniformities during the deceleration phase of the imploding capsule results in growth at approximately this radius; spikes of the high-density shell penetrate the hot spot, whereas bubbles of lower-density material distort the high-density shell. These large bubbles in the high-density shell may permit heat and fuel to escape, thereby decreasing fusion yields. These asymmetries result in angular variations in areal density, defined as

$$\rho R = \int_0^R \rho(r)dr,$$

the radial integral of the areal density, where ρ(r) is the density along the radius of the target, and R is the outer radius of the target. Another manifestation of these asymmetries is large-scale fluid flow resulting in residual kinetic energy, i.e., kinetic energy that has not been converted to hot-spot energy. Diagnosing these asymmetries is important for identifying a potential source of performance degradation. Until recently, x-ray images from a single view have been used to infer the existence of asymmetries in OMEGA cryogenic implosions.

Figure 150.54
A typical density and temperature profile at peak neutron production in an OMEGA cryogenic implosion.
although quantitative measures are outstanding. Multiple views of neutron-based diagnostics have also been available, although a tool to interpret the results has been unavailable. In this work, we show that with additional views of the neutron-based diagnostics, asymmetries can potentially be diagnosed on OMEGA. Results from a newly developed postprocessor \textit{IRIS3D} (see Appendix, p. 110) for 3-D hydrodynamic codes are described here. The role of background in interpreting the various regions of the neutron spectra for the inference of areal density and its asymmetries is emphasized; the use of multiple detectors to infer a map of asymmetries is studied; and, finally, the detection of asymmetries through the effect of the residual fluid flow on neutron spectra is studied. This work is the first such analysis; future work will include detailed post-processing of 3-D simulations to compare quantities derived from neutron spectra with experiment.

The generalized Lawson criterion for ICF implosions\textsuperscript{8} provides a measure of target performance and can be written for OMEGA scale as

\begin{equation}
\chi_{1-D} \approx \langle \rho R \rangle \left[ \frac{(T) \text{ (keV)}}{4.4} \right]^{0.8},
\end{equation}

where \( \langle \rho R \rangle \) is the neutron-weighted areal density of the compressed target and \( (T) \) is the neutron-weighted ion temperature in the hot spot. A value of \( \chi_{1-D} \approx 1 \) indicates marginal ignition, where the ratio of the output fusion neutron energy to the input laser energy is \( \sim 1 \). This form of the generalized Lawson criterion is derived from a power law fit to a series of spherically symmetric simulations and is written in terms of quantities—neutron-weighted areal density and ion temperature—that are, in principle, observable. As discussed later, areal density is inferred from the elastically scattered neutron spectrum,\textsuperscript{9–11} and ion temperature is inferred from the width of the neutron spectrum of the D–T fusion neutrons.\textsuperscript{12} Asymmetrically driven implosion experiments limit, however, the ability to directly compare the experimentally inferred quantities with spherically symmetric simulations. Areal densities can vary around the imploding target and the inferred value depends on the viewing direction. Similarly, fluid flow in the hot spot can change the width of the neutron spectrum, resulting in direction-dependent influences on the apparent ion temperature. Therefore, multiple measurements are required to constrain the values attained in implosion experiments. In addition, background neutron-producing reactions in the target can introduce ambiguities in the interpretation of neutron spectra. Therefore, it is important to understand the role of backgrounds and perform studies on different patterns of asymmetry to be able to interpret observations. Simultaneously, comparing results from 3-D hydrodynamic simulations to observations is necessary to identify the adequacy of the modeling and infer the role of nonuniformity seeds in experiments. This requires a tool to post-process 3-D simulations and compare observables with experiment.

In this article, a Monte Carlo neutron-tracking code \textit{IRIS3D} (see Appendix, p. 110) is used to model neutron transport from three primary fusion reactions in a DT capsule: the 14.1-MeV D–T fusion neutrons, the 2.45-MeV D–D fusion neutrons, and the T–T three-body reaction, which results in a continuum of neutrons. Additionally, three secondary interactions of the primary D–D and D–T neutrons including the elastic scattering off the deuterium and tritium ions (used to diagnose areal density) and the neutron-induced deuteron breakup reaction (which provides a background to the elastic scattering reactions) are also modeled. Neutron-induced triton breakup is not included in this work because its cross section is \( \sim 4 \times \) smaller than the deuteron breakup reaction.\textsuperscript{13}

In the following sections, the basic well-known relationship between \( \rho R \) and elastic scattering is described,\textsuperscript{9–11,14} the effect of neutrons from other fusion processes in the compressed target on the inference of areal density, i.e., the role of backgrounds, is discussed; and the kinematics of the elastic scattering reaction to infer the areal density in different parts of the target is exploited. The inference of areal-density asymmetries with multiple detectors is discussed along with inferring ion temperatures from neutron spectral widths. It is shown that the effect of fluid flow is an increase in the width of the neutron spectrum, leading to an increase in the inferred ion temperature. Therefore, significantly different values of ion temperature inferred from different directions around a compressing core should indicate the presence of asymmetries. It is also shown that while the absolute values of the inferred temperatures might be inaccurate, the relative values still potentially track the underlying asymmetry. Finally, conclusions are presented.

**Areal Density and Elastic Scattering**

The number of neutrons that scatter elastically, \( y_n' \), as a number of primary neutrons, \( y_n \), that move along a path \( s \) is given by

\begin{equation}
y_n' = \int_s y_n n \sigma ds,
\end{equation}

where \( n \) is the number density of deuterons or tritons (particles per unit volume) and \( \sigma \) is the cross section of the scattering
interaction. For the D–T primary neutrons at 14.1 MeV, the cross section for scattering off deuterons is \( \sigma = \sigma_d \approx 640 \text{ mb} \), and for scattering off the background tritons is \( \sigma = \sigma_t \approx 930 \text{ mb} \). If the substitution

\[
n = \frac{\rho}{\overline{m}} f
\]

is used, where \( \overline{m} \) represents the average mass of particles in the material and \( f \) represents the fraction of particles that are of the type of interest (for DT fuel with equal proportions of D and T, \( M = 2.5 \text{ amu} \), and \( f = 1/2 \)), Eq. (3) can be rewritten as

\[
y_n' = \int y_n \frac{\rho}{\overline{m}} f \sigma ds.
\]

The primary neutron yield \( y_n \) is not a constant along the path length since some fraction undergoes various reactions in the compressed target, including the process of elastic scattering; however, \( y_n \) typically changes by less than 10% along a path. For the purpose of illustrating the dependencies, it is assumed to be constant. The code used here takes into account the reduction of \( y_n \) along a path. Therefore, since \( \sigma \) and \( f \) are constants along the path length of the neutrons, the ratio of the scattered to primary neutrons can be written as (also called the down-scatter ratio or DSR)

\[
\frac{y_n'}{y_n} \approx \sigma \int_{s} \rho(s) ds.
\]

Including scattering off both the deuterons and tritons, one can write an expression for the total DSR as

\[
\text{DSR} \equiv \frac{y_n'}{y_n} = \frac{f_d \sigma_d + f_t \sigma_t}{\overline{m}} \rho L.
\]

where

\[
\rho L = \int_{0}^{L} \rho ds
\]

is the areal density along a total path length \( L \). The DSR is the observed quantity in the experiment used to infer areal density. If the trajectories of the neutrons were purely radial, \( \rho L \) would correspond to \( \rho R \), where \( R \) is the radius of the plasma. Since fusion reactions occur throughout the hot core, neutron trajectories are not radial, even for spherically symmetric implosions. Therefore, comparisons between the calculated and the observed DSR are required to identify if the simulated areal density has been achieved in implosion experiments. In addition, \( \rho L \) is dominantly sampled when the neutron-production rate is the highest in the experiment. Consequently, the areal density inferred in the experiment corresponds to a neutron-weighted value. In the remainder of this article, \( \rho L \) and \( \rho R \) will be used interchangeably. It should be kept in mind, however, that the inference of an areal density from a neutron spectrum results in a value for \( \rho L \), whereas lineouts from the center of a simulated profile would provide a measure for \( \rho R \).

To test and illustrate the physics associated with neutron interaction and transport, an ice-block profile that approximates typical profiles in a cryogenic implosion was used (Fig. 150.55). This ice-block model is characterized by regions of constant density and temperature. Neutron spectra (Fig. 150.56) using a Monte Carlo approach were calculated for the profile. Calculated cross sections for elastic scattering, which previously have shown excellent agreement with measurements, were used in this calculation. This Monte Carlo code, which post-processes spherically symmetric simulations, has also been compared previously with experiment and shown to be in excellent agreement when detailed capsule simulations are post-processed to obtain spectra for a limited class of implosions. Of note in the neutron spectrum are the DT primary peak at 14.1 MeV and the deuteron and triton backscattered edges at 1.5 and 3.5 MeV, respectively.

The DSR is influenced by the continuous portion of the spectra as shown in Fig. 150.56. Parts of the spectra are measured in OMEGA experiments by two methods: (1) the magnetic recoil

![Figure 150.55](TC13276JR)

A spherically symmetric “ice-block” test profile. A 50–50 DT fuel ratio was used.
spectrometer (MRS)\(^9\) is used to measure the down-scattered neutron spectrum between 10 and 12 MeV. (2) Neutron time of flight (nTOF) is used to measure the ratio between 3.5 and 4.5 MeV, although energies up to 6 MeV are measured currently and extension to higher energy ranges is possible.

Kinematically, DSR inferred from a specific energy range and from a particular direction samples \(\rho L\) from only a specific portion of the target. If a neutron (or any particle) of mass \(m\) and with an initial energy \(E\) scatters elastically from a particle with mass \(M\) and emerges with an energy \(E'\), its angle of scattering \(\theta\) is given by

\[
\cos \theta = \frac{x(A+1)-(A-1)}{2\sqrt{x}},
\]

where \(A = \frac{M}{m}\) and \(x = \frac{E'}{E}\).

As a result, neutrons that scatter close to their original energies are deflected by small angles, whereas neutrons that emerge at low energies are nearly backscattered. When viewed from a single direction, the energy ranges of the neutron spectrum can be mapped to different parts of the target. The open squares in Fig. 150.57 show the locations of the MRS and nTOF detectors in the OMEGA target chamber. The contours represent 1-MeV energy ranges over the regions in the detector’s view under the assumption of a point neutron source from the center for scattering from deuterons [Figs. 150.57(a) and 150.57(c)] and tritons [Figs. 150.57(b) and 150.57(d)]. An extended source would blur the regions and extend the region somewhat but not significantly enough to change the conclusions. The 10- to 12-MeV range viewed by this detector has the advantage of focusing in on a narrower region of the target; however, a large number of detectors viewing this range would be necessary to get full, 4\(\pi\) sampling coverage as a result. Since multiple detectors for the higher energy range can be expensive, multiple views at lower energies (i.e., through time-of-flight measurements), which can map broader regions of the target, are considered here.

For an nTOF detector using a range of 1 to 6 MeV, a much larger portion of the target can be sampled [Figs. 150.57(c) and 150.57(d)]. Note that only the higher three sub-rings are visible for triton backscatter [Fig. 150.57(d)] since the backscattered peaks for triton scattering occur around 3.5 MeV. Breaking up this observed spectrum into smaller sub-ranges makes it possible for a detector to sample \(\rho L\) in smaller, narrower regions of the target. In addition, by using multiple detectors, not only \(\rho L\) but also asymmetry in \(\rho L\) can be mapped to different regions of the target. As discussed in the next section, however, backgrounds can significantly influence the spectrum at these lower energies and must be accounted for carefully.
Backgrounds for Areal-Density Inference

The spectrum shown in Fig. 150.56 includes the primary and the elastically scattered neutrons. Additional effects such as thermal and Doppler broadening and multiple scattering are discussed in this section.

Multiple scattering is taken into account in the particle tracking code by recursively scattering neutrons off the background fuel ions until the effect is no longer numerically significant. Figure 150.58 demonstrates the effect that multiple scattering can have on down-scattered spectra for two different areal densities. Because DSR is approximately proportional to $\rho L$, as we have seen, multiple scattering levels should be proportional to some power of $\rho L$. Consequently, the effect of multiple scattering should become more significant as $\rho L$ increases. This is shown in Fig. 150.58, where for two ice-block profiles with different areal densities, the neutron spectrum is calculated with and without multiple scattering. For the higher areal density, characteristic of cryogenic implosions at the National Ignition Facility (NIF), significant differences are observed in the neutron spectra at lower energies, indicating that multiple scattering is important under such conditions. For OMEGA-scale implosions, areal densities are typically around 0.15 to 0.25 g/cm$^2$. In this situation, the effect of multiple scattering is negligible and can be ignored in the calculations.

All dominant components of the neutron spectrum from a cryogenic DT capsule are shown in Fig. 150.59. Again, the ice-block spectrum in Fig. 150.55 is used to calculate the spectrum. Two additional primary neutron–generating reactions are likely to occur in a DT-filled capsule: $d(d,n)^3$He and $t(t,2n)^4$He. DD primary neutrons, generated around 2.45 MeV, are not a significant background to the down-scattered neutrons since they are clearly recognizable in the spectrum. They can interact, however, with the cold fuel, for example, by elastic scattering, causing them to contribute to the background below 2.45 MeV—e.g., around the deuteron backscattered peak.

![Figure 150.58](image_url)

Spectra generated by IRIS3D considering only DT primaries and deuteron and triton scattering, with and without multiple scattering: (a) for the ice-block profile shown in Fig. 150.56, which has a $\rho R$ of 0.32 g/cm$^2$; and (b) for the same profile, but with a higher shell density of 500 g/cm$^3$, resulting in a $\rho R$ of 1.52 g/cm$^2$.
Primary neutrons from the T–T reaction, on the other hand, do produce neutrons that act as background to the down-scattered signal because the T–T reaction produces three products, so the energies of the outgoing neutrons can be anywhere from 0 up to almost 10 MeV. Unlike the down-scattered spectrum, however, T–T reactions are independent of $pL$. Since IRIS3D also assumes that the reaction is isotropic (detailed cross-section information is unavailable), this background is assumed to be independent of the viewing angle.

Neutrons produced from the deuteron breakup reaction, $d(n,2n)p$, are yet another source of background to the down-scattered spectrum. Similar to the elastic-scattering reaction, this reaction is proportional to $pL$. In IRIS3D, this reaction is calculated only when the incoming neutron is a DT primary neutron because of the unavailability of cross sections for other energies. With an end point of 11.8 MeV, this interaction acts as a background mostly in the lower-energy portions of the spectrum. Note that when areal-density asymmetries are present, this reaction can result in ambiguities in the interpretation of areal densities from the down-scattered neutron spectrum. The directionality of neutrons from this reaction is shown in Fig. 150.60. Neutrons are launched radially from the center of the target and the locations of the product neutrons are plotted in neutrons per steradian. Figure 150.60 shows that the large majority of deuteron-breakup neutrons emerge in the forward direction, but because there is no 1-to-1 correspondence between energy and scattering angle for an individual deuteron-breakup neutron (because there are three product particles), even these forward-emerging neutrons can have arbitrarily low outgoing energies. Therefore, an exact interpretation of the down-scattered neutron spectrum is challenging since it requires knowledge of the areal-density asymmetry to subtract the background.

The cumulative neutron spectrum (including all backgrounds) is shown in Fig. 150.61. The figure indicates that at lower energies, backgrounds can significantly influence the spectrum and therefore the inferred areal density. This is also summarized in Table 150.VI, which indicates that the $pL$ inferred from the calculated spectra in the 3.5- to 4.5-MeV range including the background is considerably larger than the assumed areal density, whereas the background has a marginal effect on the inferred areal density in the 10- to 12-MeV range.

![Figure 150.60](TC13280JR)

Figure 150.60
Hammer projection representing angular distribution of deuteron-breakup neutrons. Each point on the surface of a projection represents a direction in the target chamber.

![Figure 150.61](TC13281JR)

Figure 150.61
Spectra generated by IRIS3D for the profile shown in Fig. 150.56 with and without background effects. The areal density inferred from each spectrum in the 3.5- to 4.5-MeV and 10- to 12-MeV ranges is shown in Table 150.VI.

<table>
<thead>
<tr>
<th>Inferred $pL$ (g/cm²)</th>
<th>3.5 to 4.5 MeV</th>
<th>10 to 12 MeV</th>
</tr>
</thead>
<tbody>
<tr>
<td>DT scattering only</td>
<td>0.337</td>
<td>0.338</td>
</tr>
<tr>
<td>All effects and interactions</td>
<td>0.559</td>
<td>0.340</td>
</tr>
</tbody>
</table>

Table 150.VI: The $pL$ values inferred from the spectra shown in Fig. 150.61 over the specified energy ranges. These values should be compared to a target $pK = 0.32$ g/cm².
Detecting Areal Density and Asymmetries

Background can also have an effect on the inference of $\rho R$ asymmetries. Figure 150.62 shows how the detection of areal density around an asymmetric target can be affected by the inclusion of background for two different asymmetry patterns. Asymmetric profiles are obtained by perturbing the ice-block profiles with Legendre modes. To illustrate the effect, an $\ell = 1$ pattern with a peak-to-valley amplitude of 0.2 g/cm$^2$ is imposed on the profile [Fig. 150.62(a)]. In what follows, it is assumed that infinite coverage around the target chamber is available. When the 3.5- to 4.5-MeV range is used to infer areal density, the opposite phase is measured for the mode [Fig. 150.62(b)]. This is because the neutrons being observed are backscattered and originate from the opposite side of the target from the detector. The minimum and maximum inferred areal densities when the background contributions are not calculated are 0.27 and 0.41 g/cm$^2$, respectively, so the absolute values are not inferred accurately, the overall asymmetry pattern is still encoded in the neutron spectra if the pattern is dominated by $\ell = 1$.

Somewhat different results are obtained when an $\ell = 2$ mode is imposed on the profile [Fig. 150.62(d)]. The $\ell = 2$ pattern is also apparent whether or not background contributions are considered. Additionally, the phase of the underlying areal-density pattern is reproduced by the inferred areal density [Fig. 150.62(b)]. Since the asymmetry being applied is now an even mode, sampling the rear of the target actually results in the same areal density being sampled as would be seen at the front of the target, when viewed from any particular direction. The maximum-observable contrast is $\approx 0.080$ g/cm$^2$; when all background contributions are considered, this number drops to only $\approx 0.075$ g/cm$^2$. The inclusion of background does not significantly hinder asymmetry detection for this mode, relative to the $\ell = 1$ perturbation. This is again because of the opposite parities of the $\ell = 1$ and $\ell = 2$ modes. Recall from Backgrounds for Areal-Density Inference (p. 104) that (1) the deuteron-breakup background is $\rho R$ dependent; (2) its neutrons emerge mostly in the forward direction (see Fig. 150.60); and (3) it is most significant in the backscattered (low-energy) portion of the neutron spectra, including the 3.5- to 4.5-MeV range. Therefore, when a detector infers $\rho L$ using this range, it samples $\rho L$ simultaneously from the opposite side of the target via elastically backscattered neutrons and the front side of the target via deuteron-breakup neutrons. In odd-mode asymmetries, these two locations will always have opposite areal densities, hindering the detection of those asymmetries, while for even-

![Figure 150.62](image)

Figure 150.62
(a) Hammer plot of the line integral of $\rho R$ from the center of the target for a profile with an imposed $\ell = 1$ mode. As in Fig. 150.57, each point represents a location in the target chamber. The profile used is the same as that shown in Fig. 150.56, but with the spherical-harmonic perturbation applied to the value of the density in each cell based on its location. (b) Hammer plot of observed $\rho L$ for the profile shown in (a). At each point is plotted the $\rho L$ inferred from a detector at that location looking at the 3.5- to 4.5-MeV range, as calculated by IRIS3D, considering only DT primaries and single deuteron and triton scattering. (c) Same as (b), but with all effects and interaction calculated; [(d)–(f)] same as (a)–(c), but an $\ell = 2$ mode imposed on the ice-block profile.
mode asymmetries, they will always have equal areal densities, enhancing the detection of the asymmetries. Also note that the contrast is reduced relative to the 0.2-g/cm² amplitude of the imposed mode. This is caused by the shorter wavelength of the asymmetry and the finite solid angle of the detector.

The decreasing contrast with increasing mode number is demonstrated in Fig. 160.63. Even with infinite coverage and background ignored, the contrast for an imposed mode with \( \ell = 4 \) falls to 8% of the imposed value for the backscattered neutrons (3.5 to 4.5 MeV) and to 17% for the forward-scattered neutrons (10 to 12 MeV)—extremely low values that likely cannot be detected, indicating that the limit in diagnosing asymmetries is likely longer than the wavelengths corresponding to \( \ell = 4 \).

Background subtraction is the subject of active investigation for low-energy neutron spectra and is not discussed further in this work. For the remainder of this article, it is assumed that background can be accounted for by fitting an extended spectrum, potentially up to 10 MeV. The discussion that follows considers theoretical limits on inferred areal densities with a more realistic scenario: a finite number of detectors.

An areal-density map can be reconstructed as follows: For each detector and each energy sub-range, a \( \rho L \) can be inferred based on the spectral height in that sub-range as measured by that detector, as is described in Areal Density and Elastic Scattering (p. 101) [see Eq. (6)]. This inferred \( \rho L \) can be projected in a ring on a sphere enclosing the target as shown in Fig. 150.57. As a way to combine deuteron and triton elastic scattering, an average value of \( A = 2.5 \) can be used when calculating projection angles via Eq. (7). After this is done for each detector and energy sub-range, rings that overlap can be averaged together in the regions of overlap, resulting in a reconstructed areal-density map in \( 4\pi \). Note that this projection assumes that neutrons are produced as a central source that introduces an error in the maps. This error is potentially quantifiable with neutron images that indicate the location of the primary source in the compressed core. Neutron imaging is not possible on OMEGA, however, because of limited neutron statistics; although, as shown below, if the perturbation were of a significantly large wavelength, the inference of areal-density variations would remain robust.

A proof-of-principle test of the reconstruction technique for \( \rho R \) was performed by post-processing a 3-D HYDRAcryogenic implosion simulation, including realistic nonuniformity seeds from laser beam imbalances and the offset of the target in the target chamber. Primary D–T neutrons and elastically scattered neutrons were tracked for 20 time slices around peak neutron production. The density profile at peak neutron production shows a dominant \( \ell = 1 \) Legendre mode in the compressed core, primarily because the target was not placed at the target chamber center (Fig. 150.64). The lineout of the areal density taken from the location of the peak neutron production is also shown in Fig. 150.64(b). With infinite coverage, the inferred areal densities from the neutron spectrum are shown in Fig. 150.65(a). Since the dominant mode is given by \( \ell = 1 \), the phase is reversed compared to Fig. 150.64(b). Four detectors were arranged in an ad hoc tetrahedral fashion around the target in IRIS3D, and the code was used to simulate their observed spectra from 1 to 6 MeV. The locations of these detectors are shown in Fig. 150.65(b). Neutron spectra are used to reconstruct an areal-density map using the procedure described previously. The shape of the reconstructed map compares well to the lineout map seen in Fig. 150.64(b), although the contrast

---

**Figure 150.63**  
(a) Same as Fig. 150.61(a), with a Legendre model \( \ell = 4 \) imposed on the profile;  
(b) same as Fig. 150.61(b), but for the profile shown in part (a) of this figure;  
(c) same as (b), but using the 10- to 12-MeV energy range.
is significantly lower, at about 0.07 g/cm$^2$, compared to about 0.26 g/cm$^2$ in the lineout or $\sim$27%. Figure 150.65(a) suggests, however, that even with arbitrarily many detectors viewing the low-energy portion of the spectrum, the maximum-observable contrast is $\sim$0.10 g/cm$^2$. This suggests that if one can account for background, a small number of detectors have the potential to reconstruct accurate areal-density maps of low-mode asymmetries around a target. It should be noted, however, that the success of any application of this method depends greatly on where the detectors happen to be placed relative to whatever asymmetries are present. If, as complementary experiments in room-temperature plastic shell implosions suggest, there is a systematic $\lambda = 1$ mode (Ref. 16), the locations for the new detectors can be optimally prescribed to detect the mode for these cryogenic implosions.

**Ion-Temperature Inference**

The width around the primary peak in the neutron spectrum, for D–T and D–D fusion neutrons, depends on target ion temperatures and fluid flow. The variance $\sigma_n^2$ in energy for a primary spectrum$^{17}$ is

$$\sigma_n^2 = \frac{2m_nT}{m_n + m_\alpha} + 2m_nE_0\sigma_v^2,$$

(8)

where $m_n$ is the mass of a neutron, $E_0$ is the mean energy of the primary neutron (14.1 MeV for D–T primaries and 2.45 MeV for D–D primaries), $m_\alpha$ is the mass of the non-neutron product ($^4$He for D–T and $^3$He for D–D), and $\sigma_v^2$ is the variance in the component of the fluid velocity along the direction of the detector; i.e.,

$$\sigma_v^2 = \text{Var}(\vec{v} \cdot \vec{d}).$$

(9)

The inferred temperature from the neutron spectrum is given by$^{18}$

$$T_{\text{fit}} = \frac{\Delta E_{\text{fit}}^2}{E_0} \frac{m_n + m_\alpha}{16m_n \log 2},$$

(10)

where $\Delta E_{\text{fit}}$ is the full width at half maximum (FWHM) of the spectrum. Therefore, the ion temperature can be inferred by measuring the primary neutron spectra. Note that effects such as fluid velocity or background contributions to the neutron spectrum, which widen D–T or D–D primary spectra, will increase the apparent D–T or D–D ion temperatures as seen from that direction (Fig. 150.66 and Table 150.VII). The neutron spectrum around the D–D fusion neutron peak [Fig. 150.66(a)]
can be broadened because of these effects while only marginally influencing the width of the D–T fusion neutron. The two contributions to the spectra are low-energy backgrounds from other neutron processes in the target and the lowering of the height of the D–D neutron peak because of its higher scattering cross sections, which cause a significant increase in apparent ion temperature (Table 150.VII). The apparent temperature calculated from just the D–D primary spectra is 2.97 keV compared to the thermal value of 3 keV (see Fig. 150.56), while that calculated from the spectra including all contributions without any corrections to the background is 3.89 keV—a 31% increase.

These effects are relatively insignificant when inferred from the D–T neutron spectrum [Fig. 150.66(b)]. This is because (1) the scattering cross sections of 14.1-MeV neutrons are smaller compared to those of 2.45-MeV neutrons and (2) there is a lack of backgrounds above 12 MeV, including all the neutron interactions. As Table 150.VII indicates, the increase is less than 1%.

The second term in Eq. (8), corresponding to fluid flow, can significantly affect neutron spectral widths. To isolate the effect of fluid flow, the same hydrocode profile from Fig. 150.64 was used and only D–T and D–D primaries were tracked. Figure 150.67 shows that the profile contains a jet of fluid flow directed from the high-density region of the \( \ell = 1 \) mode to the low-density side.

The inferred ion temperatures from D–T and D–D fusion neutron spectra (Fig. 150.68) correspond closely to this flow. Observed temperatures for both D–T and D–D neutrons are highest at the positions corresponding to the extremes in density shown in Fig. 150.64(b) and lowest in a ring at about 90° to these two points. For any direction in the ring around the \( \ell = 1 \) mode, the jet seen in Fig. 150.67 points orthogonally to that direction, so it has a very small effect on the value of \( \sigma_v^2 \) as seen from that direction as calculated in Eq. (9). For the directions that point along the \( \ell = 1 \) mode, however, there is a large variation in \( \vec{\nu} \cdot \vec{d} \) [Eq. (9)], leading to a large \( \sigma_v^2 \) [Eq. (8)].

Table 150.VII: Ion temperature values inferred from the neutron spectra shown in Fig. 150.66 from D–D and D–T reactions. These values should be compared to a thermal temperature of 3 keV.

<table>
<thead>
<tr>
<th>Inferred ion temperature (keV)</th>
<th>D–D reaction</th>
<th>D–T reaction</th>
</tr>
</thead>
<tbody>
<tr>
<td>D–T/D–D primaries only</td>
<td>2.97</td>
<td>2.9</td>
</tr>
<tr>
<td>All neutron interactions</td>
<td>3.89</td>
<td>3.0</td>
</tr>
</tbody>
</table>

Figure 150.66
(a) Spectra, near the D–D primary peak, generated by IRIS3D for the profile shown in Fig. 150.56 with and without other interactions included. Note that fluid flow is not accounted for in either case. (b) Same as (a), but for the D–T primary peak. The D–T and D–D temperatures inferred from each spectrum are shown in Table 150.VII.

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a result, the apparent temperature is much higher than it would be with no fluid flow in these directions. These results indicate that with multiple views of the primary neutron spectra, the relative ion temperature values inferred can provide an indication of low-order asymmetries in the hot spot.

Conclusions

Using a newly developed particle tracking code \textit{IRIS3D} (see Appendix, below), neutron-based spectra are used to study signatures of asymmetry in OMEGA-scale cryogenic implosions. These include areal-density variations and neutron spectral-width variations around the compressed target. Background subtraction from observed neutron spectra is extremely important for the lower-energy range (1 to 6 MeV) to infer true areal densities. It is found that because of the finite area of the target viewed by detectors, a limited number of detectors that use a neutron spectrum up to 6 MeV can provide information on the underlying asymmetric structure of the compressed shell. Although not presented in this work, it has been found that the shape of the neutron spectrum changes from the spherically symmetric shape. This information will be used to isolate the effect of asymmetry and backgrounds in future work. Background subtraction from the lower-energy D–D fusion neutron peak is critical for reliably inferring neutron width. It is also shown that measurements of neutron width can be correlated with an overall direction of fluid flow from both DD and DT ion temperatures, provided background correction has been performed for the DD temperature. Detailed simulations and comparisons with experiment for a range of implosion parameters will be presented elsewhere. We will continue to use \textit{IRIS3D} as a postprocessor for 3-D hydrodynamic codes and will pursue detailed comparisons with observations for OMEGA cryogenic implosions.

ACKNOWLEDGMENT

The authors thank Dr. Ken Anderson for providing the profiles from \textit{HYDRA} simulations used in this article. This material is based upon work supported by the Department of Energy National Nuclear Security Administration under Award Number DE-NA0001944, the University of Rochester, and the New York State Energy Research and Development Authority.

Appendix: Structure and Methods of \textit{IRIS3D}

\textit{IRIS3D} is a parallel Monte Carlo–based neutron-tracking code with variable particle weights. Neutrons generated by a variety of interactions are tracked through a spherical grid made up of hexahedral cells such as those shown in Fig. 150.69. Note that in Fig. 150.69 and throughout this appendix, \( r \), \( \theta \), and \( \phi \) refer to spherical coordinates with \( \theta \) denoting the polar angle and \( \phi \) denoting the azimuthal angle. Each cell is indexed as described in the figure and is characterized by a DT fuel density, an ion temperature, and a fluid vector velocity. To preserve the hexahedral structure of each cell, \textit{IRIS3D} imposes exclusion zones within a small distance of the origin and within a small angle of each pole that are not occupied by cells. Therefore, each vertex with an \( i \)-index of zero is not at \( r = \)}
0 but at a very small distance from the origin, and each vertex with a $j$-index of zero or the maximum index is not at $\theta = 0$ or $\pi$, respectively, but at 0 plus a small angle and $\pi$ minus a small angle, respectively. These exclusion zones are small enough that a very limited number of particles can pass through them in a given simulation. They have an insignificant effect on results, and particles that do pass through them are simply propagated across them by a reflection, where they continue on through the remainder of the grid as usual.

A trajectory in IRIS3D is generated in some cell, $(i, j, k)$, representing some number of neutrons $y_n$, all at some energy $E$, and moving in some direction $\vec{d} = (u, v, w)$, where $\vec{d}$ is a unit vector and $u, v,$ and $w$ are the direction cosines. The initial position of the trajectory, $P_0 = (x, y, z)$, is set so that the trajectory has entered an exclusion zone, as described above, the particle is appropriately reflected to continue along the grid. For each trajectory, this process is carried out until the particle exits the grid. When this occurs, the trajectory is binned. Note that the target is assumed to be very small compared to the target chamber; therefore, the target is assumed to be essentially point-like from the view of any detectors, so trajectories are binned only according to their direction $\vec{d}$.

IRIS3D currently models six neutron-emitting interactions. Spectra from these interactions are shown in Fig. 150.59: primary neutron–producing fusion reactions including $\text{d(d,n)}^4\text{He}$, $\text{d(d,n)}^4\text{He}$, and $\text{t(t,2n)}^4\text{He}$; and three secondary interactions including $\text{d(n,n')d''}$, $\text{t(n,n')t'}$, and $\text{d(n,2n)p}$. The secondary interactions involve the interactions of primaries as they traverse the grid.

The grid lines are at constant $\theta$ or $\phi$ values, so for each trajectory, its $(\theta, \phi)$ direction is calculated according to $\vec{d}$, and it is binned in whichever grid rectangle its $(\theta, \phi)$ direction lies. In addition, detectors at specific $(\theta, \phi)$ locations and with specific solid angles $\Omega$ are specified by the user (for example, these would correspond to existing detector locations on OMEGA). For each of these detectors, a trajectory is binned in them if and only if

$$\vec{d} \cdot \vec{t} \geq 1 - \frac{\Omega}{2\pi},$$

where $\vec{d}$ is a unit vector in the direction of the detector. Detectors collect both time-integrated and time-resolved neutron spectra.

In what follows, we discuss in detail how trajectories are launched. First, the number of fusion reactions in each cell is calculated. To do this, a reactivity $\langle \sigma v \rangle$ is calculated in each
cell based on cell ion temperatures using the analytic fit given in Ref. 19, and then the number of fusion reactions, $y_{\text{cell}}$, is calculated as

$$y_{\text{cell}} = \frac{n_i n_j}{1 + \delta_{ij}} \langle \sigma v \rangle_{ij} V \Delta t$$

$$= \frac{f_i f_j}{1 + \delta_{ij}} \left( \frac{\rho}{m} \right)^2 \langle \sigma v \rangle_{ij} V \Delta t,$$

where $n_i$ and $n_j$ are ion densities for the two different ions, $f_i$ and $f_j$ are the ion fractions for the two ions characterized during the target-fabrication process, $V$ is the cell volume, $\Delta t$ is the duration of the time slice being processed for neutron spectra, and $\delta_{ij}$ is the Kronecker delta ($= 1$ if $i = j$, i.e., for D–D fusion). If $N$ is the total number of trajectories that are to be launched across the entire grid for that time step (determined by requiring adequate statistics in the calculated spectra), then the number of trajectories launched per cell is $(y_{\text{cell}} / \sum y_{\text{cell}}) \cdot N$.

For each primary trajectory to be launched from a cell centroid, $d$ is chosen as a random unit vector (since the primary trajectory loses any directional character after the reaction), and $E$ sampled from a normal distribution with mean $\mu$ and standard deviation $\sigma$ determined by

$$\mu = E_0 + (\vec{v} \cdot \vec{d}) \sqrt{2m_nE_0}$$

and

$$\sigma = \sqrt{\frac{2m_n T_i E_0}{m_n + m_\alpha}},$$

where, for example, $E_0 = 14.1$ MeV is the mean neutron energy from the D–T fusion reaction; $\vec{v}$ the cell fluid velocity; $m_n$ and $m_\alpha$ are the masses of the neutron and alpha particles, respectively; and $T_i$ is the cell ion temperature. The second term in the expression for the mean of the energy distribution takes into account Doppler shifts in the energy of the neutron caused by the fluid velocity. The expression for the standard deviation takes into account the broadening of the neutron spectrum caused by the plasma temperature.

T–T reactions are treated similarly. Reactivity for temperatures below 1 keV is calculated using the analytic fit given in Ref. 20, while reactivity for ion temperatures above or equal to 1 keV is calculated by linear interpolation using the look-up table given in Ref. 21. The initial neutron energy $E_0$ is sampled from a distribution obtained by an $R$-matrix calculation, which agrees well with experimental measurements. Note that the T–T primary neutrons are not launched in pairs that obey the conservation laws but instead are launched one at a time, independently. Momentum and energy are conserved only with adequate statistics.

Secondary reactions are considered next. IRIS3D starts by estimating a value of $y'$, the total number of deuteron-scattered neutrons:

$$y' = y_{\text{total}} \frac{f_D \sigma_{14.1} \rho R}{m},$$

where $y_{\text{total}}$ is the total number of primary neutrons generated ($= \sum y_{\text{cell}}$), $\sigma_{14.1}$ is the cross section for deuteron elastic scattering for neutrons at 14.1 MeV, and $\rho R$ is a directionally averaged look-up table given by Ref. 23. The number $N_s$ of secondary trajectories launched from a given interaction location is a code input to obtain converged results; therefore the weight of each trajectory is $y_0 = y'/N_s$.

The number of deuteron (or triton)-scattering interactions, $y'$, is calculated as

$$y' = y \frac{f_D \sigma_{D} \rho \Delta s}{m},$$

where $y$ is the number of neutrons represented by the trajectory, $\sigma$ is the cross section for deuteron elastic scattering for neutrons with energy $E$, and $\Delta s$ is the path length. At this point, $y$ is replaced with $y' = y'$ since the primary trajectory loses any neutrons that scatter away. The weight of the trajectory is $y_0$, and a scattering angle $\theta$ is sampled by interpolating between a set of energy-dependent angular cross sections from Ref. 15. An azimuthal scattering angle $\phi$ is then picked randomly over the interval $(0, 2\pi)$, and the scattered neutron trajectory direction $\vec{d}'$ is set as

$$\vec{d}' = \vec{a} \sin \theta \cos \phi + \vec{b} \sin \theta \sin \phi + \vec{c} \cos \theta,$$

where $\vec{a}$ is the direction of the original trajectory and $\vec{a}$ and $\vec{b}$ are a set of unit vectors that are orthogonal to each other and to $\vec{d}$. Finally, $E'$, the energy of the scattered neutrons, is determined based on $\theta$ using Eq. (7). Once generated, the scattered trajectory is again handled as described previously.
The deuteron-breakup interaction is modeled somewhat differently. A look-up table for angular distributions of emerging particles is available only for 14.1-MeV incident neutrons, so the interaction is considered only in the case where the original trajectory represents D–T primary neutrons. A constant value of $\sigma = \sigma_{14.1} = 164.821 \text{ mb}$ is used for both Eqs. (16) and (17) (see Ref. 24). Two trajectories are launched for the product neutrons since two neutrons are produced in each break-up reaction. Since there is no deterministic relationship between the angle and energy of each emergent neutron, the scattering angle is sampled from a distribution calculated based on Ref. 25, and energy is sampled from a distribution based on that scattering angle. As was the case with the T–T primary reaction, the neutrons’ energies and scattering angles are both sampled independently, so momentum and energy are not necessarily conserved in any particular interaction but conservation improves with increasingly better statistics.

If a series of time slices from the same implosion are used instead, the process is repeated for each time step and the results are accumulated over time.

REFERENCES

9. P. G. Young, G. M. Hale, and M. G. Chadwick, ENDF/B-VII.1, IAEA Nuclear Data Services, 22 December 2011.
10. A. Deltuva, Institute of Theoretical Physics and Astronomy, Vilnius University, Vilinus, Lithuania, private communication (2016).
Flying Focus: Spatiotemporal Control of the Laser Focus

Introduction
The controlled coupling of a laser to a plasma has the potential to address grand scientific challenges including reaching the Schwinger limit, developing compact free-electron lasers, extending colliders to TeV energies, and generating novel light sources. Currently, many such applications have limited flexibility and poor control over the laser focal volume. In conventional near-diffraction–limited systems, both the minimum focal-spot size \( w_0 \approx f^2 \lambda \) and longitudinal focusing range \( Z_{R} \approx f^2 \lambda^2 \) are linked by the ratio of the focal length to twice the beam radius \( f = 2R \). As a result, these systems require large laser spots to extend their focusing range or waveguides to maintain small spots over long distances. At low energies, manipulation of the spatial phase overcame this limitation, but a long focal range introduced in this way does not possess dynamic properties. Pulse-front tilt was recently used to introduce a time-dependent rotation of the local wavefront in a scheme called “attosecond lighthouse,” but it lacked the long longitudinal focusing range.

“Flying focus” is an advanced focusing scheme, where a chromatic focusing system combined with chirped laser pulses enables a small-diameter laser focus to propagate nearly 100× its Rayleigh length while decoupling the speed at which the peak intensity propagates from its group velocity. This unprecedented spatiotemporal control over the laser’s focal volume allows the laser focus to co- or counter-propagate along its axis at any velocity. Experiments validating the concept measured subluminal \((-0.09c)\) to superluminal \((39c)\) focal-spot velocities, generating a nearly constant peak intensity over 4.5 mm. The flying focus allows simple, compact systems to exert novel control over laser–plasma interactions and presents opportunities to overcome current fundamental limitations in laser-plasma amplifiers, laser-wakefield accelerators, photon accelerators, and high-order frequency conversion.

Figure 151.1 shows a schematic of the configuration that generates a flying focus. A diffractive lens with a radially varying groove density \( G = r / (\lambda_0 f_0) \) is used to produce a chromatic focus, where \( f_0 \) is the focal length of the system at the central wavelength \( \lambda_0 \) and \( r \) is the distance from the optical axis. With this lens, the longest wavelength \( \lambda_r \) focuses a length \( L \approx f_0 (\Delta \lambda / \lambda_0) \) before the shortest wavelength \( \lambda_b = \lambda_r - \Delta \lambda \). By introducing a laser pulse with a temporally varying wavelength, the focus will move at a velocity given by \( v(z) = dz / dt \), where \( dz \) is the distance between two focused colors spectrally separated.
by \( \delta \lambda; \) \( d \tau = d x + d z/c \) is the time it takes for the two colors to reach their respective foci; \( d \tau \) is the time between the two colors (\( \delta \lambda \)) within the chirped laser pulse, and \( c \) is the speed of light. By changing the chirp of the laser beam, the time to reach focus for successive colors is varied to provide control of the focal velocity. In general, the velocity of the focus is given by

\[
\nu(\zeta) = c \left( 1 + (d\Lambda/d\tau)^{-1} (dz/d\Lambda)^{-1} \right)^{-1}, \tag{1}
\]

where \( dz/d\Lambda \approx -f_0/\lambda_0 \) is the longitudinal dispersion provided by the diffractive lens and \( \tau = t - z/c \). For a desired longitudinal focal-spot trajectory \( z(\tau) \), a laser chirp can be designed:

\[
\lambda(\tau)/\lambda_0 = \left( 1 - z(\tau)/f_0 \right)^{-1} \approx z_0(\tau)/f_0. \tag{2}
\]

For a trajectory with a constant velocity \( z(\tau) = v_0 \tau \), a linear chirp is required: \( \lambda(\tau) = (v_0 \lambda_0/f_0) \tau + \lambda_{r,b} \), where \( v_0 = L/T \), \( \lambda_{r,b} \) is the initial wavelength, \( T \) is the chirped-pulse duration, and \( \tau \approx T/2 \).

Figure 151.2 shows the velocity of the flying focus [Eq. (1)] for a linearly chirped laser beam (\( d\Lambda/d\tau = \Delta \lambda/T \)). When the wavelengths are arranged in time where the longest wavelength comes first (positive chirp), the focal spot propagates in the forward direction (i.e., away from the diffractive lens) at subliminal velocities. For a negatively chirped laser beam (i.e., when the shortest wavelength comes first), any focal-spot velocity is available. When the pulse duration of the laser is equal to the transit time of the light to propagate across the focal region (\( T = L/c \)), all of the colors focus simultaneously, generating a long line focus; from Eq. (1) this corresponds to an “infinite” focal velocity.

This article presents experiments that demonstrate the flying focus by measuring the temporal evolution of the focal-spot intensity at various longitudinal locations. From these measurements, the velocity of the focal spot was determined and compared with the theory. The following sections (1) describe the experimental setup where LLE’s Multi-Terahertz (MTW) laser\(^{26} \) was used to demonstrate the flying-focus concept; (2) present the main results where the laser pulse duration was varied to demonstrate unprecedented control of the focal volume; and (3) discuss the potential applications for the flying focus. In particular, we explore using the flying focus to accelerate an ionization wave at the group velocity of accelerating photons.

![Figure 151.2](E26377JR)

The measured (points, bottom axis) and calculated (curves, top axis) [Eq. (1): \( \nu/c = (1\pm cT/2L)^{-1} \)] focal-spot velocity plotted as a function of the pulse duration of the laser. The red (blue) symbols represent a positively (negatively) chirped laser pulse. For all but two of the data points, the error in the velocity measurements is smaller than the symbols (<2.5%). For the data point with a pulse duration of 14 ps (very close to the \( L/c \)), the error in the velocity measurement is large since the focal velocity is nearly 50\times the speed of light.

**Experimental Setup**

MTW is a Nd:glass optical parametric chirped-pulse–amplification (OPCPA) laser with a central wavelength of \( \lambda_0 = 1054 \) nm. The bandwidth (\( \Delta \lambda = 9.2 \) nm full width at 0.1\x maximum) was stretched to produce a 2.6-ns linear chirp, and a set of compressor gratings subsequently compressed the pulse to the desired chirped-pulse duration. Undercompression relative to the transform-limited pulse duration resulted in a positive linear chirp \([\lambda(\tau) = (\Delta \lambda/T) \tau + \lambda_r \] and overcompression resulted in a negative linear chirp \([\lambda(\tau) = -(\Delta \lambda/T) \tau + \lambda_b \]. A diffractive lens with a focal length of \( f_0 = 511 \) mm (at \( \lambda_0 \)) generated an \( \sim 15-\mu \)m-diam focus with a longitudinal separation of \( L \approx 4.5 \) mm between the extreme wavelengths. This focal region was nearly \( 100 \times \) the Rayleigh length \((Z_R = 0.05 \) mm\) of the \( f/7 \) system.

The velocity of the focus over the longitudinal separation was determined by measuring the radial intensity profile along
the laser beam’s axis as a function of time. The experiments used a parallel-path configuration (Fig. 151.3), where the collimated laser beam \((R = 3.5 \text{ cm})\) was split into two identical beams to form signal and reference paths that were then imaged onto a P510 Rochester optical streak system (ROSS) camera. Inside one of the parallel paths, the signal path was focused by the diffractive lens \((f_0 = 550 \text{ mm})\) and the reference path was focused by an achromatic lens with an \(f_1 = 400\)-mm focal length. Both legs used achromatic lenses \((f_{r,s} = 400 \text{ mm})\) to collimate the light that was then recombined with a slight angle to separate the images at the detector plane. The beams were focused to the detector with a final achromatic lens \((f_2 = 400 \text{ mm})\) that produced an image of the reference and signal focal regions. Modeling indicated that the optical system was \(\sim 3\times\) diffraction limited \((\sim 15 \mu \text{m})\) over the wavelength range of interest. The spatial resolution at the detector plane of the ROSS camera was \(\sim 50-\mu \text{m} \) full width at half maximum (FWHM). The reported pulse duration \((T)\) was determined using the reference pulse measured on the ROSS camera. The impulse response of the streak camera was measured to be 7-ps FWHM.

The diameter of the signal pulse as a function of longitudinal position \((z)\) along the longitudinal focal length was determined by moving the collection lens \((f_s)\) over successive positions spanning slightly beyond the range of extreme focal positions. At each \(z\) position, several images were recorded by the streak camera and averaged to increase the signal-to-noise ratio. The reference pulse was used to remove jitter between images. Each composite image generated a measurement of the time between the reference pulse and the signal pulse \((\tau)\).

Figure 151.4 shows the results for a negatively chirped laser pulse with a duration of \(T = 36.4 \pm 1 \text{ ps}\). The images indicate that the focal spot counter-propagated at a velocity of \(-0.77c\pm 2\%\). When measuring the focal spot at a position closest to the diffractive lens \((z = -1.5 \text{ mm})\), the diameter of the flying focus was measured to evolve in time from a large spot size to a best-focus spot size over the pulse duration (i.e., the laser spot does not come to focus until the end of the laser pulse). This is in contrast with the measurements that image a position 3.0 mm farther from the diffractive lens \((z = 1.5 \text{ mm})\). In this case, the focal-spot size was measured to start at its best focus.
and expand to a maximum diameter over the duration of the laser pulse (i.e., the laser spot starts at focus and expands until the end of the laser pulse).

The velocity of the focus \( v = \frac{\Delta z}{\Delta t} = c \left[ 1 + \left( \frac{\Delta \tau}{\Delta z} \right) c^{-1} \right] \) was determined by measuring the time of minimum foci (\( \tau \)) at each image plane (\( z \)). The slope of a best-fit line to the measured data [Fig. 151.4(b)] was used to determine \( m = \frac{c \Delta \tau}{\Delta z} \).

The error in the measurements shown in Fig. 151.2 is given by \( \delta v = v \delta m \), where \( \delta m \) is the uncertainty in each fit.

**Results**

Figure 151.5 shows measurements of the flying focus generated by both a negatively and a positively chirped laser pulse. The initial frame of the negatively chirped pulse shows the laser beam entering the focal region, but before it has reached focus.

The evolution of the flying focus intensity measured for a negative (left) and positive (right) chirped pulse, each with a duration of \( T \sim 60 \) ps. In each case, the laser is shown propagating into the measurement window (top left) at 0 ps. In the positively chirped case, the laser comes into focus at the left edge of the window (\( z \sim -2.5 \) mm), in contrast to the negatively chirped case, where the pulse is far from focus. At \( t = 25 \) ps (top middle), the negatively chirped case shows that the laser has reached focus at the back of the window (\( z \sim +2 \) mm). Over the next few frames, the focus propagates \( \sim -2 \) mm in \( \sim 20 \) ps, corresponding to \( -0.3c \), while over the same time, the positively chirped pulse moves forward slowly at \( \sim +0.2c \).
Over the next 20 ps, the laser reaches a focus at the far end of the system (z ~ 2 mm). This is in contrast with the positively chirped pulse, where the laser comes into focus initially at the front of the measurement window (z ~ -2 mm). Comparing the middle row for each data set shows that the focal spots are propagating in opposite directions. For the negatively chirped pulse, the peak intensity moved back toward the lens by Δz ~ 2 mm over the ~20 ps corresponding to a velocity of ~0.3c, while for the positively chirped pulse, the peak intensity moved forward by about the same distance in a comparable time corresponding to a velocity of about +0.2c. Figure 151.5 was constructed from temporal measurements of 30 longitudinal locations ranging from z = -3.75 mm to z = +3.75 mm. The measured images were sliced into temporal bins and recombined given their focal location and measured time (t).

The measured velocity of the focus as a function of the pulse duration of the laser compares well with the calculations using Eq. (1) (Fig. 151.2). The results show that when the laser pulse was negatively chirped with a duration of T = 34.4 ps, the focal spot counter-propagated at a velocity of v = -0.87c±2%. Reducing the pulse duration (T = 18.6 ps) resulted in a counter-propagating superluminal focus (v = -7.6c±20%). Extending the pulse duration to T = 232 ps slowed the focal spot propagating at vlc ≈ -0.09±1%. When the pulse duration was just less than the transit time of the light to propagate across the focal region, the focus was measured to propagate at nearly 50× the speed of light. A positive chirp provides access to a range of forward-propagating subluminal velocities. The focal-spot velocity for a positively chirped laser pulse with a duration of T = 65 ps was measured to propagate at v = 0.20c±1%.

Figure 151.6 shows snapshots of the longitudinal intensity profiles calculated for three different negative chirp cases. They illustrate propagating backward at the speed of light [Fig. 151.6(a)], propagating instantaneously across the focal volume [Fig. 151.6(b)], and propagating forward faster than the speed of light [Fig. 151.6(c)]. They were calculated by assuming Gaussian optics, I(z,t)/I_0 = [w_0/w(z,t)]^2, where w_0 ≈ 1/(2G_R) is the diffraction-limited spot size and

\[ w(z,t) ≈ w_0 \left( 1 + \frac{f_0^2}{4G_R^2} \left( \frac{z}{\lambda_0} + \frac{\lambda_0}{\Lambda(t)} \right) - 1 \right)^2 \]  

(3)

is the radius of the flying focus spot. The Rayleigh length for a diffractive lens is given by

\[ Z_R ≈ \frac{f_0^2\lambda_0}{4R^2} \approx 1/4G_R^2\lambda = 52 \mu m, \]

where G_R is the groove density at the radius of the laser beam (R). This is a reasonable approach to calculating the intensity profile provided that the pulse duration is much larger than the radial pulse front delay (T > T_RPFD = 5 ps).

The intensity of the flying focus across the longitudinal focal region is given by the spectral power, I(z,t) = P(λ) / π [w(z,t)]^2, which shows that the longitudinal intensity can be controlled by spectrally shaping the laser pulse. In the experiment, 1.6 nm of bandwidth was removed from the middle of a positively chirped spectrum, which demonstrated that the laser did not focus over the central region of the longitudinal focus. The measured laser
focus propagated subluminally \((v/c = 0.16\pm1\%)\) over the first ~2 mm and then did not focus again for ~26 ps, at which time the focus reappeared at \(z \approx 2.8\) mm and propagated to the end of the longitudinal focal region.

**Applications**

For more-exotic applications, the velocity of the focus can be varied by using a nonlinear chirp and/or a nonlinear chromatic optical system. From Eq. (1), it is evident that the focal velocity could be made to accelerate, decelerate, or oscillate across the longitudinal focal region depending on the design of the nonlinear chirp. An example that demonstrates the impact of the flying focus is a photon accelerator. A photon accelerator frequency upshifts light using rapidly changing density \(n_e\) generated by, for example, an ionization wave. Prior photon accelerator concepts have been limited by phase slip-page, where the upshifting laser beam accelerates out of the density gradient. A flying focus using a nonlinear chirp could mitigate this by making the velocity of the ionization wave follow the group velocity of the upshifting beam:

\[
\frac{dz}{dt} = v_g(t) = c\sqrt{1 - \frac{\omega_p^2}{\omega(t)^2}},
\]

where \(\omega_p^2 = n_e e^2 / m_e \epsilon_0\) and \(n_e\) is the maximum electron plasma density. In this case, the photons will be frequency upshifted from an initial frequency \(\omega_0\):

\[
\frac{\Delta \omega(t)}{\omega(t)} = \frac{\omega(t) - \omega_0'}{\omega_0'} = \sqrt{1 + \left(\frac{\omega_p}{\omega_0'}\right)^2 \frac{z(t)}{Z_R}} - 1,
\]

where \(z(t)\) is the trajectory of the ionization wave (i.e., the trajectory of the flying focus) and \(Z_R\) is an approximate width of the ionization wave that was assumed to be equal to the Rayleigh length of the flying focus.

Figure 151.7(a) shows the results from Eq. (5) where photons with an initial group velocity of \(v_g = 0.7c\) were accelerated to \(v_g = 0.99c\) over 4.5 mm (from \(\lambda_0 = 1054\) nm to \(\lambda' \approx 160\) nm at \(n_e = 5 \times 10^{20} \text{ cm}^{-3}\)). In a standard photon accelerator design where the ionization wave propagates at a constant velocity given by the initial group velocity of the seed photons, the accelerated photons would be limited to \(v_g = 0.9c\) (~550 nm). In this case, the accelerated photons overtake the ionization wave within the first 0.3 mm. The maximum photon energy in a photon accelerator driven by a flying focus is limited by the accelerator length, which is given by the total bandwidth in the laser \([L = f_D(\Delta \lambda/\lambda)]\).

Figure 151.7(b) shows the corresponding nonlinear chirp that is required to follow the accelerating trajectory. There are two solutions that both require a negative chirp. The solutions depend on whether the pulse duration of the flying focus is greater than or less than the time it takes for light to transverse the accelerator \((T = L/v_g \approx 15\) ps\). When the pulse duration is

---

**Figure 151.7**

(a) The velocity of the accelerating photons (left axis, dashed curve) and their wavelength (right axis, solid curve) are plotted as functions of accelerator length for a system where the ionization wave is produced by an accelerating flying focus. The electron density was assumed to rise from vacuum to \(n_e = 5 \times 10^{20} \text{ cm}^{-3}\) over the Rayleigh length of the flying focus \((Z_R = 0.05\) mm\)). (b) The nonlinear chirp is required for the flying focus to accelerate in phase with the frequency-shifted photons toward the diffractive lens (bottom axis) and away from the diffractive lens (top axis).
longer than the $L/v_g$, the flying focus will counter-propagate with respect to the flying focus beam; when the pulse duration is shorter than $L/v_g$, the flying focus will co-propagate. This nonlinear chirp accounts for the initial rapidly changing group velocity of the accelerating photons [Fig. 151.7(a)]. Extending the bandwidth to a typical value available in current ultrashort pulse lasers ($\Delta \lambda/\lambda_0 \approx 200 \text{ nm}/1000 \text{ nm}$) lengthens the accelerator to nearly $L \approx 10 \text{ cm}$, and the accelerated photons reach a final wavelength of 100 nm, assuming the same conditions for the ionization wave as above. The maximum wavelength shift could be significantly increased by using a density ramp to maintain a constant $\omega_p/\omega'(t)$ as the photons are accelerated.

**Summary**

The flying focus provides an avenue for novel control over laser–plasma interactions, removes the need for long-focal-length systems or guiding structures to maintain high intensities over long distances, and decouples the velocity of the focal spot from the group velocity of the light. In addition to photon accelerators, the spatiotemporal control of laser intensity achieved by the flying focus has the potential to change the way plasma devices are optimized and could be applied in many areas of physics. In a laser wakefield accelerator,\(^\text{28,29}\) the flying focus could eliminate dephasing by generating a focal spot that moves at a velocity that matches the accelerating electrons. This separation of the accelerator length from the plasma density will provide larger accelerating fields for a given accelerator length and could expand the options for optimizing laser-plasma accelerators. Furthermore, applying the flying focus to a laser-plasma amplifier will allow the ionizing pump laser intensity to propagate at $v = -c$ in order to generate a counter-propagating ionization wave just ahead of the amplifying seed pulse. This will enable one to control the plasma conditions observed by the seed and could be the enabling technology for an efficient laser-plasma amplifier (see the next article, **Raman Amplification with a Flying Focus**).

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**REFERENCES**

Raman Amplification with a Flying Focus

Introduction
Continuing to push the boundary of laser intensity using existing technology is increasingly challenged by the need for large, efficient, and damage-resistant gratings. Here, we propose a new laser amplifier scheme utilizing stimulated Raman scattering in plasma in conjunction with a “flying focus”—a chromatic focusing system combined with a chirped pump beam that provides spatiotemporal control over the pump’s focal spot. Localized high intensity is made to propagate at \( v = -c \) just ahead of the injected counter-propagating seed pulse. By setting the intensity in the interaction region to be just above the ionization threshold, an ionization wave is produced that travels at a fixed distance ahead of the seed. Simulations show that this will make it possible to optimize the plasma temperature and mitigate many of the issues that are known to have impacted previous Raman amplification experiments, in particular the growth of precursors.

Plasma-based laser amplifiers utilizing either stimulated Raman scattering (SRS) or strongly coupled stimulated Brillouin scattering have long been of interest. Lacking a damage threshold, compact plasma-based systems could produce unfocused intensities \( I \approx 10^{17} \text{ W/cm}^2 \)—more than six orders of magnitude larger than conventional systems. Typically, a moderate-intensity pump pulse with a duration of at least \( 2L/c \) propagates across a plasma of length \( L \). When the pump’s leading edge reaches the end of the plasma, an initially weak seed pulse is injected in a counter-propagating geometry. Tuned to satisfy the Manley–Rowe frequency- and wave-number–matching conditions, the beat wave between the two beams drives a plasma wave that mediates energy transfer from the pump to the seed.

An alternate scheme has been proposed to mitigate precursor growth in which the seed ionizes the plasma coincident with its amplification by the pump. However, this introduces additional constraints: the pump intensity must be below the threshold for ionization, limiting the Raman growth rate; conversely, the initial seed intensity must be high enough to photoionize the plasma, limiting the degree to which it can be further amplified; and the ionization itself damps the growing seed pulse. To our knowledge, this scheme has yet to be tested because of the added complexity.

A Raman amplifier with a flying focus retains the advantages of seed ionization while eliminating its downsides. A chirped pump is focused by a diffractive lens that introduces chromatic aberration in order to produce a longitudinally distributed focal spot. The temporal dispersion provided by the chirp, combined with the spatial dispersion provided by the lens, provides spatiotemporal control over the propagation of intensity isosurfaces. In the example shown in Fig. 151.9,
the pump has a negative linear chirp and a pulse duration that is equal to \(T = 2L/c\), where \(L\) is both the length of the focal region spanned by its bandwidth and the length of the amplifier interaction region. In this case, the desired pump intensity first appears where the pump exits the interaction region and subsequently propagates backward at \(v \approx -c\) at a constant value over a length that can be many times the Rayleigh length. More details regarding the flying focus optical system, along with additional applications, are contained in the companion article *Flying Focus: Spatiotemporal Control of the Laser*, p. 115.\(^1\)

To demonstrate the benefits of this concept, the coupled three-wave equations describing SRS in plasma were solved numerically (see, e.g., Refs. 13 and 26 and references therein). Such models have previously been benchmarked against particle-in-cell simulations and found to be in good agreement when plasma-wave amplitudes were kept below the wave-breaking limit and kinetic effects could be ignored [at low \(k_3\lambda_{De} \leq 0.3\), where \(k_3\) is the electron plasma wave’s (EPW’s) wave number and \(\lambda_{De}\) is the Debye length].\(^27\) The three-wave model is supplemented with a field ionization model to simulate the plasma ionization by the pump.\(^26\) The model is described in much greater detail in *Methods*, p. 125.

For all of the simulations, the initial density of hydrogen atoms was \(6 \times 10^{18} \text{cm}^{-3}\) and the interaction length was 4 mm, defining a pump duration of 26.7 ps. The pump wavelength was \(\lambda = 1 \ \mu\text{m}\) and the seed was upshifted by the EPW frequency. For the flying focus Raman amplification (FFRA) base case, the pump focusing system was \(f/5\) with the focus of each color located past the interaction region. To simulate focusing in this 1-D model, the pump enters from the left edge and its intensity increases as it propagates to the right in a manner that is consistent with the \(f\) number of the system. The blue leading edge of the pump converges to a spot diameter of 400 \(\mu\text{m}\) at the exit of the interaction region, where the intensity was set to be \(I_1 = 1.4 \times 10^{14} \text{W/cm}^2\).

In the simulations, the plasma mediating the energy transfer was formed by the pump beam ionizing the hydrogen gas within the interaction region. The ionization threshold of hydrogen is very close to the optimal pump intensity in systems designed for \(\lambda \approx 1-\mu\text{m}\) lasers. Since the pump first reaches this intensity at the right edge of the amplifier in the case of FFRA, plasma is initialized there and an ionization wave subsequently propagates backward with the intensity isosurface. The setup can therefore be tuned so that the plasma is formed just before the seed arrival at every point along the interaction region.

The peak of a 500-fs-duration (full width at half maximum) seed pulse with an initial intensity \(I = 1 \times 10^{11} \text{W/cm}^2\) was injected at the right edge just after the arrival of the pump’s leading edge (\(t = 14\) ps). Figure 151.10(a) shows three snapshots of the interaction as the injected seed travels from right to left across the interaction region for the FFRA case. The first frame shows that the gas is ionized only \((n_e/n_c > 0)\) close to the right edge, where the pump first reaches high intensity. The seed duration stretches as it grows in the linear regime. From the first to the second frame, it is clear that the ionization wave is propagating at an approximately fixed distance ahead of the seed. The nonlinear pump-depletion regime has been reached, with seed pulse compression and the formation of a secondary peak. This efficient amplification continues in the final frame. These results demonstrate the ideal behavior that is expected when the seed enters unperturbed plasma and competing instabilities are avoided.

Contrast Fig. 151.10(a) with the behavior observed in Fig. 151.10(b), which shows the results from a Raman amplifier without the flying focus. In this case, the intensity was set to \(I = 1.4 \times 10^{14} \text{W/cm}^2\) at the left edge of the amplifier and was assumed to be collimated as it propagated from left to right (consequently, the pump intensity seen by the seed pulse...
Figure 151.10
Results of three-wave model simulations. (a) With the flying focus, the pump first reaches high intensity at the right edge, where ionization is initialized. Constant intensity moves at $v = -c$ as different colors converge to different locations, so the ionization wave propagates at a nearly fixed distance ahead of the injected seed pulse. Ideal plasma amplifier behavior is observed. (b) When the pump is collimated within the interaction region and above threshold for ionization, the seed encounters higher temperatures along nearly its entire path, which reduces growth via increased Landau damping. (c) With a collimated beam as in Case 2 but holding $T_e$ fixed to be similar to Case 1, spontaneous stimulated Raman scattering (SRS) grows during the long time in which the pump propagates across the ionized plasma. Premature pump depletion degrades the resulting seed amplification. Flying focus Raman amplification (FFRA) Case 1 with noise initialized at the same level did not produce such precursors. EPW: electron plasma wave.

is nearly the same in both cases). The first frame shows that upon reaching the right edge, the plasma is ionized everywhere throughout the interaction region. While growth in the first frame is comparable, it slows rapidly compared to FFRA. Pump depletion and pulse compression fail to occur in this case.

The difference can be understood by looking at the electron temperature encountered by the peak of the seed pulse versus time [Fig. 151.11(a)]. In FFRA “Case 1,” after a brief initial growth period, $T_e$ levels off at $\approx 45$ eV because of the nearly constant duration of plasma heating by the pump prior to the seed’s arrival at each point along its path. With standard focusing (or a preformed plasma), the seed encountered plasma that was heated for a progressively longer duration as it propagated, producing a strong gradient in $T_e$ (Case 2). This model captures the fact that excessive heating can lead to debilitating levels of collisionless Landau damping, which acts to suppress the seed growth.\footnote{13,15} Figure 151.11(b) shows the sum of collisional and collisionless damping as a function of temperature. The former dominates at low temperatures and the latter at high temperatures; FFRA Case 1 is close to the temperature at which EPW damping is minimized.

Note that there could be additional impacts of elevated temperature that are not captured by this model. The thermal gradient seen by the seed pulse can lead to resonance detuning resulting from the Bohm–Gross frequency shift.\footnote{14} Detuning can also result from the kinetic nonlinear frequency shift that accompanies particle trapping.\footnote{16,19,25} Perhaps most importantly, the wave-breaking threshold is reduced in warm plasma,\footnote{17–19,28} which limits the plasma-wave amplitudes and thereby the energy transfer from pump to seed. This model, therefore, likely underestimates the adverse effects of high temperature and lack of temperature control with a conventional focusing and ionization scheme.

Given the uncertainties, a temperature of $\approx 45$ eV may not be optimal. A nice feature of the FFRA scheme, however,
is that the temperature can be easily tuned by adjusting the delay between the ionization wave and the injected seed pulse. Many parameters can influence this delay. Holding all else constant but injecting the seed 3 ps later, its peak encounters an electron temperature that is uniformly higher by ~20 eV [c.f., Fig. 151.11(a), Case 3]. Because of the higher temperature, it takes longer to reach pump depletion and the secondary peaks are suppressed. Both the interaction pump intensity relative to the ionization threshold of the gas and the pump’s f number are additional parameters for tuning the delay between ionization and seed injection.

To investigate nonthermal differences between FFRA and standard Raman amplifiers, a Case 4 was run, repeating Case 2 but with a fixed electron temperature ($T_e = 45$ eV). Although the seed encountered a similar electron temperature everywhere in Cases 1 and 4, the pump spent a longer time in ionized plasma prior to seed injection in Case 4 compared to FFRA Case 1. The debilitating effect of spontaneous SRS growing ahead of the seed is observed in Fig. 151.11(c). Although seed growth over the first half of the plasma proceeds in a similar fashion as Case 1, subsequent growth is suppressed because of premature pump depletion and interference with pre-existing EPW’s. Although noise was included in the same manner in FFRA Case 1, no spontaneous SRS growth was observed because of the limited distance over which it could grow ahead of the seed.

As with temperature, this model likely underestimates the negative impacts of spontaneous SRS. While the zeroth-order effect is competition for pump energy, there is some evidence that saturation of even low-level precursors can corrupt plasma conditions (e.g., with driven ion-acoustic waves or modified electron distribution functions) over relatively long time scales. In these situations, the seed does not encounter quiescent plasma and its growth is compromised. The controlled introduction of frequency detuning has been proposed to mitigate precursors without precluding the desired seed amplification (resulting from the larger resonance bandwidth of the latter in the nonlinear pump-depletion regime). Despite evidence that modern experiments have been adversely affected by too much frequency detuning, spontaneous SRS continues to be an issue and was recently observed to dominate the overall backscatter as the Raman growth rate was increased.

The use of a chirped pump beam—a feature of many previous experiments—is necessary for the flying focus but does introduce some frequency detuning for fixed plasma conditions that could degrade performance. It could be compensated for, however, by introducing a density gradient along the seed path in order to exactly satisfy the frequency-matching condition everywhere. While perfect resonance may result in undue levels of spontaneous SRS in a typical plasma amplifier, it would not degrade FFRA because of the alternative means by which FFRA suppresses precursor growth.

**Methods**

The basic three-wave equations are

\[
\begin{align*}
\partial_z - v_1 \partial_x + v_1 a_1 &= K a_2 a_3, \\
\partial_z - v_2 \partial_x + v_2 a_2 &= K a_2 a_3, \\
\partial_z - v_3 \partial_x + v_3 + i\delta\omega a_3 &= K a_4 a_2^* + S_3,
\end{align*}
\]

where the subscripts 1, 2, and 3 refer to the pump, the seed, and the EPW, respectively; $v_i$’s are group velocities; $\nu_i$’s are damping rates; $K = \omega (n_e / n_c)^{1/4} / 2$ is the wave-coupling parameter, where $n_e$ is the electron density and $n_c$ is the critical density; $a_{1,2} = 0.855 \times 10^3 \lambda_{1,2} (\mu m) / \sqrt{f_{1,2}} (W/cm^2)$ are normalized laser vector potentials, and $a_3 = \left| E_3 \right| / m_e c (\omega / \omega_{pe})$ is the normalized envelope of the EPW, with pump frequency $\omega$ and EPW frequency $\omega_{pe}$. Advection of the plasma wave can be neglected ($v_3 \approx 0$), and here detuning was also neglected ($\delta \omega = 0$) since it has been explored
extensively elsewhere.\textsuperscript{13,14,18,21,22} The pump and seed are damped collisionally, $\nu_{1,2} = \nu_{e} (\omega_{pe}^{2}/\omega^{2})$ with $\nu_{e} = 2.9 \times 10^{-6} \text{Zn} \text{cm}^{-3} \Lambda T_{e}^{-3/2}$; $\nu_{3} = \nu_{e} + \nu_{L}$ includes both collisional absorption and collisionless (Landau) damping, with $\nu_{L} \approx \sqrt{\pi/2} [\omega_{pe}^{2}/(k_{3}v_{e}^{3})] \exp[-\omega_{pe}^{2}/2(k_{3}v_{e}^{3})^{3/2}]$. $S_{3}$ is a noise term that is included to investigate spontaneous SRS growing from undriven plasma fluctuations. Following Ref. 13, $S_{3} = c_{f} \nu_{e} T_{e}$ is assumed to be proportional to the EPW damping rate and electron temperature, but a multiplier $c_{f}$ was added to test the sensitivity to the initial noise level. Experiments often find that plasma fluctuations are elevated over the expected thermal levels.

The three-wave model was supplemented with an ionization model to simulate the plasma ionization by the pump\textsuperscript{26}

$$\partial_{t} n_{e} = n_{n} w(a_{1}),$$

$$\partial_{t} n_{n} = -n_{n} w(a_{1}),$$  \hspace{1cm} (2)

where $n_{n}$ is the neutral gas density and $w(a_{1})$ is the ionization rate that depends on the local pump intensity. In the regime of interest, the Keldysh formula is valid.\textsuperscript{29,30} For $\gamma = \sqrt{2U_{1}/m_{e}c^{2}/a_{1} > 1}$, where $U_{1}$ is the ionization potential, the multiphoton ionization rate $w(a) \approx a_{n}N^{3/2}(2\gamma)^{-2N}$ is appropriate, where $N = 1 + \text{int}(U_{1}/\hbar\omega)$ is the number of photons required to overcome the ionization potential. For $\gamma < 1$, the tunneling formula is more accurate:

$$w(a) = 4\Omega_{0} \left(\frac{U_{1}}{U_{H}}\right)^{3/2} \frac{\sigma_{1}}{a_{1}} \exp\left[-\frac{2}{3}\left(\frac{U_{1}}{U_{H}}\right)^{3/2} \frac{\sigma_{1}}{a_{1}}\right].$$

with atomic frequency $\Omega_{0} = 4.1 \times 10^{16} \text{ s}^{-1}$, hydrogen ionization potential $U_{H} = 13.6 \text{ eV}$, and the hydrogenic electric-field normalized vector potential $a_{1} \approx 3.05 \times 10^{14}/a_{1}$. An exponential fit was used to fill in the region between the multiphoton and tunneling regimes. The molecular nature of hydrogen was approximated by using the molecular ionization potential $U_{1} = U_{H} = 15.4 \text{ eV}$ (Ref. 26). To conserve energy, an additional damping term on the pump was added to Eq. (1) by balancing the equation $n_{e}(m_{e}c^{2}/2) \partial_{t} v_{e} = -(U_{1} + e) \partial_{t} n_{e}$, where $e = m_{e}v_{osc}^{2}/2$ is the assumed birth energy and $v_{osc}$ is the oscillation velocity of electrons in the pump laser’s electric field. The electron temperature was initialized locally at the birth energy, but may subsequently evolve to balance collisional absorption of the pump and seed.

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**REFERENCES**


Full-Wave and Ray-Based Modeling of Cross-Beam Energy Transfer Between Laser Beams with Distributed Phase Plates and Polarization Smoothing

Introduction

In direct-drive inertial confinement fusion (ICF), a millimeter-scale spherical capsule is uniformly illuminated by symmetrically oriented laser beams.\(^1,2\) The capsules have an outer ablator layer and an inner fuel layer. The lasers ablate the outer layer of the capsule, which generates pressure to implode the fuel. The ICF program relies on radiation–hydrodynamics codes for designing capsules and tuning laser conditions to optimize the hydrodynamic efficiency of implosions.\(^3\) An essential component of these codes is a model for coupling laser energy to the capsule ablator.\(^4,5\)

Laser energy is coupled to the ICF capsule primarily through electron–ion collisional absorption of the laser beams propagating through a coronal plasma that forms around the irradiated capsule.\(^6\) Because the wavelength and period of the lasers are typically much smaller than the hydrodynamic spatial and temporal scales, respectively, the eikonal approximation can be used to solve the steady-state electromagnetic-field equations along ray trajectories using the instantaneous plasma conditions.\(^7\)

In addition to collisional absorption, nonlinear laser–plasma interaction (LPI) processes affect laser-energy deposition. The main nonlinear processes that are thought to be energetically important in ICF experiments are the three-wave processes: stimulated Raman scattering (SRS), stimulated Brillouin scattering (SBS), and two-plasmon decay (TPD).\(^8\) Stimulated Brillouin scattering is the parametric coupling between two electromagnetic waves and an ion-acoustic wave (IAW). When the seed and pump electromagnetic waves in the three-wave SBS process correspond to distinct laser beams, it is referred to as cross-beam energy transfer (CBET).\(^8,9\)

Ray-based laser-energy deposition models that have been adapted to include CBET predict that laser absorption is reduced by \(\sim 10\%\) to \(20\%\) (Ref. 4) and that laser-energy deposition uniformity is significantly modified in typical direct-drive ICF experiments.\(^3,10\) Hydrodynamic simulations that include CBET show significantly better agreement with measured scattered-light spectra and implosion trajectories, but ray-based CBET calculations must still be modified using \textit{ad hoc} multipliers and field limiters to give quantitative agreement with experimental observations.\(^11,12\) In addition to the eikonal approximation,\(^7\) ray-based CBET models that are used in radiation–hydrodynamics codes assume steady-state linear convective gains, pairwise coupling between rays, and local plane-wave laser beams. Direct-drive ICF experiments also employ polarization smoothing to improve drive-beam uniformity by splitting each laser beam into two beams with orthogonal polarization and a small angular divergence.\(^13,14\) This is accounted for in ray-based models by assuming random relative polarizations of interacting beams and spatially averaged incoherence between the two polarization components of each beam.\(^15\) A more-complete model of CBET is required to test the validity of these approximations.

This article compares wave- and ray-based CBET calculations in the presence of laser beam speckle and polarization smoothing using the full-wave LPI code \textit{LPSE}.\(^16\) The wave-based calculations suggest that laser beam speckle and polarization smoothing can lead to significantly more CBET than is predicted by ray-based calculations. To account for speckle effects in the ray-based model, a modification is presented that gives excellent agreement with wave-based calculations. Full-scale wave-based calculations in hydrodynamic profiles based on direct-drive experiments on the OMEGA laser suggest that beam speckle has a small (<1%) effect on direct-drive laser absorption.

This article (1) puts the current work in the context of previous work on the beam speckle’s effect on CBET; (2) describes the ray-based CBET model; (3) describes the equations solved by \textit{LPSE}; (4) compares the ray- and wave-based CBET models for a variety of laser and hydrodynamic configurations using both linearly polarized beams and beams with polarization smoothing; and (5) summarizes the conclusions.
Relation to Other Work

The theory of the speckle statistics of laser beams generated by distributed phase plates was developed using the formalism of Gaussian random fields. There has been considerable theoretical work on the impact of laser speckle on filamentation,21,22 deflection,23,24 SBS,25–27 and CBET in the paraxial approximation relevant to indirect-drive ICF.28 Most of the previous studies of laser beam speckle in the ICF context have focused on ponderomotive self-focusing and filamentation. A recent study looked at the interaction between ponderomotive self-focusing and CBET.29 The present study focuses on the effect of laser speckle on CBET in the absence of filamentation because this is the situation most relevant to direct-drive ICF experiments, where the single-beam laser intensities are typically well below the filamentation threshold and multibeam filamentation should not be important.30 The simulations were performed using a full-wave solver that does not make the paraxial approximation and solves the vector wave equation in 3-D, which is essential for studying direct-drive ICF where polarization smoothing is used and beams cross at arbitrary angles.

Ray-based modeling of laser-energy deposition and CBET is now routine in radiation–hydrodynamic simulations of ICF experiments.4,5 Ray-based CBET models typically inject the beams as plane waves with an intensity equal to the spatially averaged intensity of the speckled laser beams that are used in the experiments. A recent study used the ray-based paraxial complex geometric optics (PCGO) approach to calculate CBET between speckled beams and compared the results to a paraxial wave-based code.31 The PCGO approach gives a pseudo-speckle pattern that produces a statistical intensity distribution similar to a real speckle pattern over a limited range of speckle intensities. We present a ray-based model for calculating CBET between speckled beams that gives improved agreement with wave-based calculations by using ray tracing to directly solve the electromagnetic-field equations within the eikonal approximation. This produces a real speckle pattern that, in the absence of CBET, exactly matches the wave-based calculation in regions of space where the eikonal approximation is valid and shows excellent agreement with the wave-based CBET calculations up to gains ≥5.

We also present a study of CBET between polarization-smoothed beams. Ray-based models rely on the assumption of uncorrelated polarization and phase to calculate the interaction between laser beams that employ polarization smoothing, but the limits of this approximation have not been previously studied because most previous studies were based on wave-based codes that solve only the scalar wave equation for the electromagnetic fields.

Ray-Based CBET Modeling

Cross-beam energy transfer is calculated along ray trajectories by numerically integrating the steady-state homogeneous gain along ray trajectories using the local plasma conditions.32 The differential change in the energy of ray \( i \) at the \( j \)th location along its path caused by an interaction with ray \( k \) at the \( l \)th location along its path is

\[
\frac{dW_{ij}}{ds} = - \frac{W_{ij}}{\alpha_{ij}} + W_{ij} \xi_{kl} \left( \frac{W_{kl}}{\varepsilon_0} \frac{dS_{kl}}{dS_{kl}} \right),
\]

where \( W_{ij} = \sqrt{\varepsilon_0} \left( dS_{ij} / dS_{il} \right) |E_{ij}|^2 \) is the ray energy, \( E_{ij} \) is the enveloped electric-field amplitude, \( \varepsilon_0 = 1 - n_e / n_c \) is the permittivity, \( n_e \) is the electron density, \( n_c = m_e \omega_0^2 / (4\pi e^2) \) is the critical density for light with frequency \( \omega_0 \), \( m_e \) is the electron mass, \( e \) is the electron charge, \( dS_{il} / dS_{ij} \) is the ratio of the initial to current cross-sectional area of ray \( i \) (proportional to the ray intensity),

\[
\xi_{ij}^{kl} = \frac{5.88 \times 10^{-2}}{T_e (1 + 3 T_i / Z T_e)^3} \left( \frac{n_e}{n_c} \right)^2 \alpha_s \nu_{IAW} P(\eta_{ij}^{kl}),
\]

\[
P(\eta) = \frac{\left( \nu_{IAW} / \omega_s \right)^2 \eta}{\left( \eta^2 - 1 \right)^{3/2} + \left( \nu_{IAW} / \omega_s \right)^2 \eta^2},
\]

\[
\eta_{ij}^{kl} = \frac{\alpha_{kl} - \alpha_{ij} - (k_{ij} - k_{kl}) \cdot u}{\alpha_s},
\]

\( \lambda_0 \) is the laser wavelength in vacuum (in microns), \( \omega_{ij} \) is frequency of the \( k \)th ray at the \( j \)th location along its path in the lab frame and \( k_{ij} \) is the corresponding wave vector, \( \omega_s \) is the acoustic frequency, \( \nu_{IAW} \) is the IAW energy-damping rate, \( T_e (T_i) \) is the electron (ion) temperature in keV, \( Z \) is the ionization state of the ions, and \( u \) is the plasma flow velocity. \( \alpha_{ij} \) is the local laser absorption length, which for electron–ion collisional absorption is equal to the group velocity over the...
energy-damping rate \( \alpha_{ij} = c \sqrt{\epsilon_0} n_i c / \nu_{ci} n_e \) evaluated at the local plasma conditions. The collisional damping rate is \( \nu_{ci} = \sqrt{2 \pi e^4 Z^2 n_i \Lambda_{ci}} / (3 \sqrt{m_e T_e^{(2)}}) \), where \( \Lambda_{ci} \) is the Coulomb logarithm. Additional corrections related to laser absorption and the temporal derivative of the background density profile are included in ray-based models used in radiation-hydrodynamics codes, but these corrections were not used in the calculations here.

Equation (1) can be solved in numerous ways that all essentially come down to defining a procedure for breaking rays up into discrete segments and a procedure for mapping the energy of nearby rays onto a given section of a ray’s path. The discretized version of Eq. (1) is

\[
W_{i,j+1} = W_{ij} \exp \left[ \frac{1}{\alpha_{ij}} + \sum_{kl} \frac{\xi_{ik}}{\epsilon_0} \frac{dS_{kl}}{dS_{ij}} W_{kl} \right],
\]

where \( s_{ij} \) is the length of the \( j \)th section of ray \( i \)’s path. Equation (2) is a nonlinear system of equations that can be solved using fixed-point iteration, but written in its current form it converges slowly because information about upstream changes along a ray propagate only one path step per iteration. The rate of convergence can be improved significantly by noting that

\[
\frac{dx}{dt} = k,
\]

\[
\frac{dk}{dt} = \frac{1}{2} \nabla e_0(x),
\]

where the wave-vector magnitude is normalized to the vacuum value \( (k = \sqrt{\epsilon_0}) \). The solution to Eq. (4) is single valued at every point in \( (x, k) \) phase space, but it does not necessarily have a single-valued projection onto \( x \) space. The divisions between regions of the phase space solution that have a single-valued projection onto \( x \) space occur at caustics, and when the solution is divided into distinct sections that individually have single-valued projections onto \( x \) space, the regions are referred to as “sheets.” In ray-based CBET calculations, a ray from each sheet interacts with every other sheet at a given point in \( x \) space. Accordingly, the sum in Eq. (3) is restricted to rays on distinct sheets. The time-enveloped electric field is reconstructed from the eikonal solution by summing over sheets:

\[
E(x) = \sum_{j \text{ sheets}} E_j(x) e^{i \phi_j(x)},
\]

where \( \phi_j \) is the phase corresponding to the field amplitude \( E_j \). Additional subtleties are involved with ray-based CBET calculations at caustics because the electromagnetic-field amplitude in the eikonal approximation is singular. This topic will be discussed in a future publication (none of the ray-based calculations presented in this article included caustics).

The ray-based CBET calculations presented in this article discretized the ray trajectories on a Cartesian grid. Figure 151.12 shows a representation of the ray indexing scheme for two interacting beams, where rays 1 and 2 correspond to one beam and rays 3 and 4 correspond to the other beam. As an example of the indexing for the interactions, the crossing of rays 2 and 4 in grid cell 3 corresponds to the ray energies \( W_{2j} \rightarrow W_{22} \) and \( W_{4j} \rightarrow W_{42} \) with the interaction coefficient \( \xi_{2j} \rightarrow \xi_{42} \). Note that this example uses the ray energies at the grid cell entrance to calculate the interaction, but the numerical scheme converges to the same result if the midpoint or endpoint is used instead. Discretization onto a grid has the advantage of simplicity when determining the ray interactions because this step is reduced to simply looking at the other rays crossing a given grid cell, which also makes this step straightforward for parallel computation. The disadvantage, relative to interpolation-based
approaches, is that the rays must be dense enough that at least one ray from each sheet passes through each grid cell in the region where CBET is occurring for the solution to be valid.

1. Speckle in Ray-Based CBET

The laser beams used in ICF experiments pass through distributed phase plates (DPP’s) that produce a speckle pattern consisting of many local minima and maxima at the focal plane. The boundary condition for the electric field of a DPP beam injected at $z = 0$ can be modeled using

$$E(x, y, z = 0, t) = \sum_k |E(k, t)| e^{i(k \cdot x + \phi_k)},$$  \hspace{1cm} (5)

where the sum is over beamlets generated by the DPP with distinct wave vectors $k$ and random phases ($\phi_k$) (Ref. 18). The eikonal solution for the electromagnetic fields with this boundary condition cannot be calculated from a single sheet of rays because the wavefront is not locally a plane wave, but it is a superposition of many plane waves, so the field is obtained by taking the coherent sum of the ray-trace solution for each term in Eq. (5). For a given boundary condition, this gives the same solution for the fields as the wave-based calculation in the absence of CBET (in regions of space where the eikonal approximation is valid). The superposition solution used to calculate the fields cannot be applied to CBET, however, because the energy transfer is not linear in the electric field.\(^{37}\)

To approximately include speckle effects in the ray-based CBET calculation, the local intensity variations were calculated using the superposition solution for the unperturbed electromagnetic fields and interpolated onto the ray trajectories used in the plane-wave calculation. The rest of the CBET calculation is identical to the plane-wave case with the intensity variations along the ray trajectories appearing as an additional term in the exponent in Eq. (3).

The primary increase in computational cost associated with the speckle model is that the speckles must be resolved on the CBET grid, which typically requires several-times-better spatial resolution for convergence than is needed for plane-wave beams. The transverse correlation length of a speckled beam is given by the product of the laser wavelength of the $f$ number and the focusing optic ($\lambda_f$) (Ref. 38), which is $\sim 2 \mu m$ for the typical laser configurations used in ICF experiments ($f/6.7$ lenses and 0.351-$\mu m$ light).\(^{13}\)

We expect that a plane-wave approximation will be sufficient over some range of interaction configurations, and that the proposed model for including speckle effects in ray-based CBET calculations will extend this range, but a more-complete model is required to test the limits of these approximations.

**LPSE**

LPSE solves the time-enveloped Maxwell’s equations coupled to the low-frequency plasma response in the fluid approximation. The plasma response is linearized around an inhomogeneous background density and flow velocity profile.$^{16}$

The time-enveloped wave equation for the electric field is

$$\frac{2i\omega_0}{c^2} \frac{\partial}{\partial t} E + \nabla^2 E - \nabla \cdot E + \frac{\omega_0^2}{c^2} \varepsilon(\omega_0; x, t) E = 0,$$

where $\varepsilon$ is the plasma dielectric function

$$\varepsilon(\omega_0; x, t) = 1 - \frac{\omega_{pe}^2(x, t)}{\omega_0(\omega_0 + i\nu_{ci})},$$

and $\nu_{ci}$ is the electron–ion collision frequency. The physical electric field is given by $\mathbf{E} = \Re \left[ E(x, t) \exp(-i\omega_0 t) \right]$. 

Figure 151.12

Indexing scheme for the ray paths and cross-beam energy transfer (CBET) grid for two interacting laser beams. The ray paths are divided according to their intersections with grid cell boundaries.
The equations for the low-frequency plasma response are

\[
\left[ \partial_t + U_0(x) \cdot \nabla \right] \left( \frac{\delta n}{n_0} \right) = -\mathcal{W},
\]

\[
\left[ \partial_t + U_0(x) \cdot \nabla + 2\nu_{IAW} \right] \mathcal{W} = -\nabla^2 \left[ \frac{\delta n}{c_s n_0} + \frac{2e^2}{4m_e m_i \omega_0^2} \left| E \right|^2 \right],
\]

where \(\mathcal{W} = \nabla \cdot \delta U\), \(U_0\) is the background flow profile, \(c_s\) is the sound speed, \(m_i\) is the ion mass, and \(\nu_{IAW}\) is a phenomenological operator used to reproduce Landau damping of IAW’s, which is implemented by applying a constant damping of \(\nu_{IAW}/2\) in \(k\) space (the factor of 1/2 appears because \(\nu_{IAW}\) is the energy damping rate, and the damping is applied to the wave amplitudes in LPSE). The density is \(n(x,t) = n_0(x) + \delta n(x,t)\) and the flow velocity is \(U(x,t) = U_0(x) + \delta U(x,t)\). LPSE uses a total-field/scattered-field approach, where the laser beams are injected inside of the total-field region and the scattered-field region acts as an absorbing boundary. Further details of the numerical algorithm and benchmarking can be found in Refs. 16 and 39.

**Comparison Between LPSE Numerical Solutions and the Ray-Based Model**

This section is divided into subsections corresponding to interactions between linearly polarized beams in a plasma with a constant density and linearly varying flow, interactions between linearly polarized beams in ICF-relevant plasma conditions, and interactions between beams with polarization smoothing.

1. Homogeneous Plasma

Figure 151.13 shows the magnitude of the steady-state electric field from a 2-D LPSE simulation of the interaction between two counter-propagating speckled beams in a plasma with a linear flow velocity profile given by \(v_{flow}(x) = -3c_s(0.005x + 1)\) and a constant density \(n_e/n_c = 0.01\). The other plasma parameters were \(T_e = 2\) keV, \(T_i = 1\) keV, \(Z = 3.1\), \(A = 5.3\) (ion mass in amu), and \(\nu_{IAW}/\omega_s = 0.2\). The grid size was \(80 \times 240 \mu m^2\) (3168 \times 9504 grid cells) and the simulations were run until a steady state was established (4 ps). The speckle patterns correspond to \(f/6.7\) lenses and were generated by launching 128 beamlets from each boundary with top-hat intensity distributions in wave-vector space [Eq. (5)] and fourth-order super-Gaussian distributions in physical space. The average initial intensity of the pump (seed) beam was \(2 \times 10^{15} W/cm^2\) (1 \(\times 10^{12} W/cm^2\)); both beams were polarized out of the plane (the average intensity is defined here as the peak intensity that a plane-wave beam would have for the same beam width and flux). The beams were launched 4 \(\mu m\) inside the simulation boundaries, and the outer 2 \(\mu m\) of the simulation grid were absorbing. The seed beam is not visible at its injection point (\(x = 36 \mu m\)) because of its low initial intensity, but it can be seen in the scattered region at the bottom of the image, where it exits with an average intensity of \(3.2 \times 10^{13} W/cm^2\) corresponding to a gain of 5.8 (the gain is defined as the log of the incident seed-beam energy over the outgoing seed-beam energy).

Figure 151.14 shows lineouts of the LPSE field magnitude from Fig. 151.13 for the (a) pump and (b) seed beams at their respective exit planes (blue) and the corresponding result from the ray-based calculation (red). The two methods give nearly identical results for the field of the pump beam because it lost only 32% of its initial energy. The seed beam was amplified by more than 300 \(\times\), and the two speckle patterns look completely different. Despite the significant differences in the structure of the calculated fields, the gain predicted by the ray-based model was only 0.8% lower than the wave-based calculation (the gain from the ray-based calculation using plane-wave beams was 38% lower). The ray-based calculation does not reproduce the detailed structure of the fields because the local variations in the direction of the wavefront were neglected when the magnitude of the fields was interpolated onto the ray trajectories from the plane-wave calculation. The interpolation procedure gives the correct statistical variations in the laser intensity, which is why the average CBET from the wave-based calculation was reproduced, but after energy is transferred, it is forced to fol-
poorly to many beam systems, the grid resolution requirement is typically lower by a factor of 10 to 100 per dimension.

Figure 151.15 shows the results of a number of calculations similar to the one depicted in Fig. 151.13, where the pump intensity and the angle between the pump and seed beam were varied. The gains from the speckled-beam calculations are compared to those obtained in calculations that were identical except the speckled beams were replaced with plane-wave beams. For the case of counter-propagating beams, there is a difference of more than two e foldings in the energy gain for the speckled beams relative to the plane-wave beams at the highest pump intensity. At lower intensities, the difference diminishes. When the angle between the beams is increased, the difference between the plane-wave and speckled cases is reduced, and for more than 30° between the beams, there is no significant difference. Figure 151.15(d) shows the beam geometry used for these comparisons, which was chosen so that the peak of the CBET resonance was always at Mach 1.

low the ray trajectories whereas it follows the local wavefront in the wave-based calculation.

The advantage of the ray-based calculation is a significant reduction in computational cost. The LPSE calculation depicted in Fig. 151.13 took one hour on 792 CPU cores, while the corresponding ray-based calculation took only a few minutes on a desktop computer. For a given grid resolution, the computational cost of the LPSE calculation is proportional to the number of grid cells \( O(N^2) \). The relevant computational requirement for the ray-based CBET calculation is the number of pairwise interactions between rays, which is \( O(N_B^2) \), where \( N_B \) is the number of laser beams. The grid resolution in LPSE is determined by the need to resolve the wavelength of light, while in the ray-based calculation, it is necessary to resolve spatial variations in the laser intensity and hydrodynamic conditions. Although the quadratic dependence on the number of laser beams causes the ray-based calculation to scale poorly to many beam systems, the grid resolution requirement is typically lower by a factor of 10 to 100 per dimension.

Figure 151.15 shows the results of a number of calculations similar to the one depicted in Fig. 151.13, where the pump intensity and the angle between the pump and seed beam were varied. The gains from the speckled-beam calculations are compared to those obtained in calculations that were identical except the speckled beams were replaced with plane-wave beams. For the case of counter-propagating beams, there is a difference of more than two e foldings in the energy gain for the speckled beams relative to the plane-wave beams at the highest pump intensity. At lower intensities, the difference diminishes. When the angle between the beams is increased, the difference between the plane-wave and speckled cases is reduced, and for more than 30° between the beams, there is no significant difference. Figure 151.15(d) shows the beam geometry used for these comparisons, which was chosen so that the peak of the CBET resonance was always at Mach 1.
Figure 151.15 also shows the corresponding gains calculated using the ray-based CBET model. The ray-based model shows excellent agreement with the full-wave calculation for the interaction between plane waves, which indicates that the assumptions made in the ray-based model that are not related to beam speckle are valid in this configuration. The ray-based model also shows very good agreement with the speckled-beam results, and the calculated gains in the ray-based model always fall within one standard deviation of the LPSE gain averaged over different speckle realizations. The fact that the ray-based speckle model is in such good agreement with the wave-based calculations even when the plane-wave assumption gives a very poor approximation suggests that this is an extremely useful modification for including speckle effects in ray-based laser-plasma simulations. An additional benefit of the ray-based speckle model is that it inherently gives a more-realistic laser-energy deposition profile than a plane-wave approximation.

The maximum intensity used in Fig. 151.15 was chosen to keep the filamentation control parameter\(^{38}\)

\[
\frac{\langle P \rangle}{P_c} = 0.04 f_n^2 n_c \left( \frac{I_{14}}{I_0} \right)^{1/2}
\]

below 1 (\(f\) is the \(f\) number and \(I_{14}\) is the laser intensity in units of \(10^{14}\) W/cm\(^2\)). At the highest average intensity \(\langle P/P_c \rangle = 0.22\). This is right at the limit where the highest-intensity speckles could potentially filament, but no filamentation was observed in the simulations depicted in Fig. 151.15. This was necessary for the comparison to the ray-based CBET model because the equations solved by LPSE implicitly include filamentation but the ray model does not.

The trends in Fig. 151.15 can be understood qualitatively by considering the average intensity over the CBET interaction region because the energy transfer is exponential in both the pump intensity and the interaction length. When the beams are counter-propagating, some of the speckles will see an increased pump intensity along the entire interaction length, giving exponentially larger gains, while the partially compensating reduction in the fraction of the beam profile that undergoes significant CBET is only a linear effect. When the angle between the beams is large, a given seed speckle will interact with many high- and low-intensity pump speckles, and the product of interaction length and pump intensity integrated over the interaction region averages out to the same value as for plane-wave beams. The same logic applies to the longitudinal extent of the speckles. In this example the speckles were longer than the interaction region, but as the length of the interaction region is increased relative to the length of the speckles, the plane-wave gains are eventually recovered regardless of the relative beam orientation.

2. Amplitude of Density Perturbations

In addition to modifying the CBET gain, most of the CBET between speckled beams happens in localized hot spots, which can lead to larger density perturbations than occur when plane-wave beams interact. Figure 151.16 shows the ratio of root-mean-square (rms) density perturbations for a speckled-beam simulation and the corresponding plane-wave simulation. The ratio is always larger than 1, indicating that the typical density perturbation is larger for the speckled-beam interaction than for the plane-wave interaction. Although LPSE does not include the relevant physics, nonlinear effects become important at large \(\delta n / n\) and cause CBET to saturate.\(^{15,38}\) Note that the amplitude of the density perturbations is insensitive to the relative beam angle, which shows that this is not simply a result of increased CBET.

![Figure 151.16](image_url)

Figure 151.16

Root-mean-square (rms) density perturbation from LPSE simulations for the interaction between speckled beams divided by the rms density perturbation for the interaction between plane-wave beams. The rms perturbation was taken from the region \(y = [-70,70]\) and \(x = [0,30]\). The region was offset toward larger \(x\) so that the energy gain of the seed beam did not have a significant impact on the amplitude of the density perturbations. The error bars correspond to the standard deviation for three different speckle realizations.

To test the qualitative impact of nonlinear saturation, the LPSE simulations depicted in Fig. 151.15 were repeated with a clamp on the amplitude of the density perturbations \(|\delta n / n| \leq 0.01\). This clamp was chosen to be more restrictive than the expected \(|\delta n / n| \sim 0.1\) threshold\(^{38}\) to exaggerate the impact of...
nonlinear saturation. Previous studies using a code similar to LPSE have shown that an even more restrictive clamp on $\delta n/n$ is required to obtain quantitative agreement with scattered-light measurements in indirect-drive ICF experiments.\textsuperscript{40} Figure 151.17 compares the LPSE gains plotted in Fig. 151.15 and the clamped simulations (only one speckle realization was simulated at each condition and compared to the corresponding simulation that used the same seed for random phase generation). At all relative beam angles and intensities, the perturbation-limited results are essentially indistinguishable for the plane-wave simulations, but in the speckled-beam simulations, the clamp reduced the amount of CBET significantly.

3. OMEGA Implosions

Although the subscale calculations presented in the previous section show that speckles can increase the CBET gain, the effect on laser absorption in direct-drive ICF experiments is expected to be small for two main reasons: (1) the single-beam intensities in ICF experiments are usually about an order of magnitude lower than the lowest intensity from Fig. 151.15; and (2) in a 3-D spherical implosion, the fraction of the solid angle that corresponds to nearly counter-propagating rays is small. The effect could be greater in indirect-drive ICF where many of the beams are nearly co-propagating although the single-beam intensities are still relatively low.\textsuperscript{15}

Figure 151.18 shows the steady-state magnitude of the electric-field envelope and the corresponding ion-density perturbations for plane-wave and speckled-beam 2-D LPSE simulations of two $f/6.7$ s-polarized beams interacting in a realistic direct-drive ICF plasma profile from the 1-D radiation–hydrodynamics code LILAC.\textsuperscript{41} The beams were injected at normal incidence 22 $\mu$m inside the minimum $x$ and $y$ boundaries. The critical density is denoted by white dashed circles. The electromagnetic-field (density-perturbation) grid was $40,000 \times 40,000$ (80,000 $\times$ 80,000) cells. The small-scale rings in the fields are aliasing artifacts associated with the wavelength-scale perturbations on the high-resolution grids.

![Figure 151.17](image1.png)

Gain as a function of pump intensity from LPSE simulations for speckled beams (red) and plane-wave beams (blue) with (circles) and without (crosses) a limiter of $|\delta n/n| \leq 0.01$ for beams separated by (a) 0°, (b) 15°, and (c) 30°.

![Figure 151.18](image2.png)

Amplitude of the electric-field envelope in a full-scale two-beam LPSE simulation (at 12 ps) using ICF-relevant plasma conditions with (a) plane-wave beams and (b) speckled beams at initial intensities of $2 \times 10^{14}$ W/cm$^2$. The corresponding density perturbations are shown in (c) and (d). The critical density is denoted by white dashed circles. The electromagnetic-field (density-perturbation) grid was $40,000 \times 40,000$ (80,000 $\times$ 80,000) cells. The small-scale rings in the fields are aliasing artifacts associated with the wavelength-scale perturbations on the high-resolution grids.
angle between the beams was set to 90° because it gives nearly maximal CBET.\textsuperscript{10} In Fig. 151.18(a), the field amplitude of the incident beams is visibly reduced after they cross through the caustic of the other beam. Because the outgoing portion of the beams gain energy from the incident beams, this interaction leads to a reduction in the total laser absorption. The density perturbations shown in Figs. 151.18(a) and 151.18(b) indicate where CBET occurs. The dominant interaction regions are in the beam caustics crossing through the middle of the other incoming beam and near to the critical density where the beams are reflected. The interaction at near-critical density is seeded by the reflected beam and is referred to as self-CBET. Self-CBET is typically not as energetically important as the interaction between distinct beams because it occurs at high densities and most of the energy that is transferred is still absorbed. The speckled-beam calculations are qualitatively similar to the plane-wave calculations in terms of where CBET occurs with the additional restriction that most of the CBET occurs in high-intensity speckles.

Two-beam simulations were performed at two different single-beam intensities, $1 \times 10^{14}$ W/cm\textsuperscript{2} and $2 \times 10^{14}$ W/cm\textsuperscript{2}. These intensities are higher than the single-beam intensities in ICF experiments, but there are many more beams interacting in that case. The intensities were chosen to give a non-negligible reduction in laser absorption resulting from CBET while staying below the filamentation threshold. In the absence of CBET, the percentage of the incident laser energy that was absorbed by the plasma in the simulations similar to those shown in Fig. 151.18 was 94% regardless of beam type/intensity. For laser intensities of $1 \times 10^{14}$ W/cm\textsuperscript{2}, the absorption was 90.5% (89.2%) for the plane-wave (speckled) beams, and for laser intensities of $2 \times 10^{14}$ W/cm\textsuperscript{2}, the absorption was 85.4% (82.8%) for the plane-wave (speckled) beams. In both cases the reduction in laser absorption was ~30% larger for the speckled-beam simulations (relative to the no-CBET simulations). However, this configuration should significantly overestimate the impact of speckles because of the high single-beam intensities.

To test the effect of using many lower-intensity beams, 2-D quarter-scale, 16-beam LPSE simulations were run using single-beam intensities of $4 \times 10^{14}$ W/cm\textsuperscript{2}. The beams were injected uniformly at 22.5° increments. In the quarter-scale, 16-beam configuration, the plane-wave simulation had 67% absorption and the speckled-beam simulation had 66.8% absorption (98% absorption without CBET). Despite the significant reduction in absorption caused by CBET, the plane-wave and speckled-beam simulations had nearly the same total laser absorption, which suggests that the use of many lower-intensity beams smooths out single-beam speckle effects but does not diminish the impact of CBET.\textsuperscript{30} Accordingly, a plane-wave approximation is expected to be sufficient for calculations of CBET between linearly polarized beams in many-beam direct-drive ICF applications.

4. Polarization Smoothing

Polarization smoothing is typically accounted for in ray-based CBET models by multiplying the gain coefficient calculated for parallel-polarized beams ($\xi_{ij}$ in Eq. (3)) by a factor of $(1+\cos^2 \theta)/4$, where $\theta$ is the angle between the interacting rays. This factor is obtained by assuming random relative polarizations of the interacting beams and spatially averaged incoherence between the two polarization components of each beam.\textsuperscript{15} This approximation can be tested by comparing the ray model to LPSE simulations, which have full polarization effects.

Figure 151.19 shows the results of a series of comparisons between the LPSE and the ray-based model using polarization smoothing and two different $f$ numbers (6.7 and 3) at pump-beam intensities of $5 \times 10^{14}$ W/cm\textsuperscript{2}. The plasma conditions were the same as discussed in Homogeneous Plasma (p. 132) (homogeneous density and constant flow velocity gradient), and polarization smoothing was obtained by splitting the energy of the incident beams evenly between $s$ and $p$ polarizations. The two polarizations had statistically independent speckle realizations. The gain calculated by LPSE was higher than the gain in the ray-based model for all but the orthogonal $f/3$ beams. The gains are higher for the $f/6.7$ because their longitudinal correlation length ($2\pi f^2\lambda_0 \approx 100 \mu$m) is longer than the interaction volume,$^{38}$ so the assumption of incoherence between the polarization components when averaged over the

![Figure 151.19](image-url)
interaction region is violated. The speckle length for the 2/3 beams is expected to be in better agreement with the ray-based calculation because of the reduced speckle length (~20 μm).

There is a significant difference between the wave-based and ray-based results shown in Fig. 151.19, even though the intensity was only $5 \times 10^{14}$ W/cm², and the gains were much lower than for the corresponding linearly polarized calculations. The discrepancy between LPSE and the ray-based model in the case where the beams are orthogonal is readily explained by noting that the in-plane components of the beams’ polarizations do not interact, and the out-of-plane components have the same interaction as two linearly polarized beams with half of the energy. If the average intensity amplification for the linearly polarized case is written $II_0 = e^G$, then this should give $I / I_0 = (1 + e^{2G})/2$, whereas using the factor of $(1 + \cos^2 \theta)/4$ gives $II_0 = e^{G/4}$.

In the case of counter-propagating beams, the difference can be understood by using a simple square-wave model, where a speckled beam is treated as having a beam spot profile with twice the average intensity over half of the beam profile and zero intensity over the rest. If the speckles are random, the interaction between two linearly polarized counter-propagating beams will have the same statistical properties as the interaction between the two beam profiles depicted in Fig. 151.20(a). The interaction can be broken up into four situations that occur with equal probability, only one of which results in any CBET of the interacting beams were not random. Figure 151.21 shows the results of calculations where the initial polarization of the two beams was chosen randomly (and independently) before applying phase plates and polarization smoothing. For the 2/3 beams, the average gain over realizations still differs from the ray-based model because of the large speckle length. The calculations using 2/3 beams are in good agreement with the ray-based model, which suggests that the gain multiplier used to correct for polarization smoothing in the ray-based models is accurate when the underlying assumptions are satisfied.

The same model can be used for polarization-smoothed beams, but each of the beams must also be split into two orthogonal polarizations. There are now four equally probable combinations of amplitude and polarization within each beam and 16 possible types of interactions between beams that are depicted in Fig. 151.20(b). Adding up the various contributions to the amplification gives $(I / I_0)_{PS} = 1/2 + (3e^G + e^{2G})/8$. This should be compared to the amplification that is used in the ray-based model, $(II_0)_{PS(rays)} = e^{G/2}$, which to leading order gives $(III_0)_{PS} = (II_0)_{PS(rays)} \approx G/8$. The lowest-order correction is linear in $G$, consistent with the deviation between the wave- and ray-based calculations occurring at lower gains for polarization-smoothed beams than for linearly polarized beams.

The comparisons shown in Fig. 151.19 are useful for illustrating why CBET between polarization-smoothed beams might not agree with ray-based calculations, but they do not represent a fair comparison because the relative polarizations of the interacting beams were not random. Figure 151.21 shows the results of calculations where the initial polarization of the two beams was chosen randomly (and independently) before applying phase plates and polarization smoothing. For the 2/3 beams, the average gain over realizations still differs from the ray-based model because of the large speckle length. The calculations using 2/3 beams are in good agreement with the ray-based model, which suggests that the gain multiplier used to correct for polarization smoothing in the ray-based models is accurate when the underlying assumptions are satisfied.

Figure 151.20
Beam profiles with the same statistical properties as a square-wave speckle model for (a) linearly polarized beams and (b) polarization-smoothed beams split evenly between $s$ and $p$ polarization.
The discrepancy at larger f numbers could be significant in ICF experiments because the gradient scale lengths are typically comparable to the speckle length of the drive beams. In Fig. 151.18(c), the majority of the CBET is occurring in the “wings” in the lower left part of the image, where the caustic of the outgoing beam gains energy as it crosses the center of the incoming beam. The spatial extent of this transfer region is \( \leq 100 \) \( \mu \text{m} \).

**Summary**

The impact of beam speckle and polarization smoothing on CBET was studied using the full-wave LPI code \textit{LPSE}. The results were compared to ray-based calculations using a code that is based on the ray models used to simulate ICF experiments. The ray-based model tends to underpredict the amount of CBET when the assumption of spatially averaged incoherence over the length of the interaction region is violated. A ray-based speckle model was presented that gives excellent agreement with the wave-based calculations over a broad range of gains and relative beam angles.

At all relative angles, the CBET interaction between speckled beams generates larger rms density perturbations than the corresponding plane-wave interaction. These enhanced density perturbations could lead to the earlier onset of nonlinear saturation of CBET between speckled beams. The single-beam intensities in ICF experiments are not sufficient for this effect to be significant, but it could play a role in many-beam interactions.

For linearly polarized beams, the large gain (\( \geq 1 \)) and small relative beam angle that were required to see a significant difference between the plane-wave and speckled-beam calculations suggest that a plane-wave approximation should not result in a significant error in laser-absorption calculations for direct-drive ICF. This conclusion is supported by \textit{LPSE} simulations in hydrodynamic profiles relevant to direct-drive ICF that showed a modest reduction in laser absorption for two-beam interactions and almost no reduction in laser absorption for a 16-beam interaction. For polarization-smoothed beams, there is a significant difference between the wave- and ray-based results at modest gains and over a broader range of relative beam angles, which could have an impact on CBET calculations for ICF.

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**REFERENCES**


Direct-drive inertial confinement fusion (ICF) is one of two laser-based techniques being pursued for achieving controlled nuclear fusion at the 1.8-MJ National Ignition Facility (NIF). In direct-drive hot-spot ignition designs, laser ablation of a spherical shell drives the implosion and compression of a cryogenic deuterium–tritium (DT) fuel layer, into which a fusion burn wave propagates after first being initiated in a central, low-density hot spot. To achieve ignition, the fuel must be compressed to an areal density greater than 0.3 g/cm², which can be achieved by keeping the pressure close to the Fermi-degenerate pressure. Preheat of the DT fuel by suprathermal electrons generated by laser–plasma instabilities (LPI’s) increases this pressure, degrades compression, and inhibits ignition. Consequently, control of LPI suprathermal or “hot”-electron production is critical for a successful implosion.

Stimulated Raman scattering (SRS) and two-plasmon decay (TPD) are two instabilities that are capable of generating hot electrons since they both excite electrostatic waves in the plasma that provide accelerating fields. SRS entails the decay of a laser light wave into an electron plasma wave and a scattered-light wave at densities at or below one quarter of the critical density of the laser, while TPD is the decay of a laser light wave into two electrostatic plasma waves (plasmons) near the quarter-critical density. Previous studies of SRS and TPD have examined single-beam thresholds, quantified suprathermal electron production, explored collective multibeam processes, and investigated the spatial properties of TPD. The interaction conditions in a previously unexplored regime. Until the experiments described herein, carried out on a MJ-scale facility, it was not possible to simultaneously achieve the density scale length, laser intensity, electron temperature, and transverse plasma dimensions that are characteristic of ignition-scale direct-drive implosions.

This article presents the first exploration of the LPI origins, scaling, and possible mitigation of hot electrons under direct-drive ignition-relevant conditions. These new observations indicate the dominance of SRS over TPD, a result not previously anticipated, with significant implications for direct-drive–ignition designs.

Planar targets were irradiated from one side with 351-nm laser light using a subset of the NIF’s 192 beams, with 1-D smoothing by spectral dispersion at 90 GHz. These beams are arranged into cones that share a common angle with respect to the polar axis. There are four such cones in each hemisphere: the “inner” cones have angles of 23.5° and 30° (32 beams in each hemisphere), while the “outer” cones have angles of 44.5° and 50° (64 beams in each hemisphere). All targets described here were irradiated using beams in the southern hemisphere. The targets were thick CH (or Si) disks with a 4.4-mm diameter and a 0.5-mm thickness, oriented toward a polar angle between 0° and 30°. Planar targets were chosen because they are the only way, currently, to achieve direct-drive ignition-relevant plasma conditions, while using a reduced laser energy (~200 kJ).
on the NIF. The use of planar targets also reduces the level of cross-beam energy transfer\(^{25}\) relative to spherical targets.

Time-resolved SRS diagnostics, with or without spectral resolution (400 to 750 nm), were located at polar angles of 23.5°, 30°, and 50° (Ref. 26), as shown in Fig. 151.22. The targets were irradiated with laser pulses of ≤8-ns duration at vacuum overlapped intensities of \(≤3 \times 10^{15} \text{ W/cm}^2\). The plasma evolution was simulated using the 2-D radiation–hydrodynamics code DRACO\(^{27}\) for comparison with experimental observations. The DRACO predictions for the density scale lengths and electron temperatures, in the vicinity of the quarter-critical density \(n_c = n_e/4\) [where \(n_c\) is the electron density and \(n_e\) is the critical density for the laser wavelength \(\lambda_0\) (µm)], with \(n_c \approx 1.1 \times 10^{21} \lambda_0^{-2} \text{cm}^{-3}\)] were \(L_n \sim 500 \text{ to } 700 \mu\text{m}\) and \(T_e \sim 3 \text{ to } 5 \text{ keV}\), respectively. DRACO simulations calculate that the laser intensity is attenuated by \(\sim 50\%\) on reaching the quarter-critical surface as a result of collisional absorption.

A time-resolved scattered-light spectrum obtained from NIF shot N160420-003 is shown in Fig. 151.22(a). It displays a narrow, intense feature at a wavelength slightly above 702 nm (2\(\lambda_0\)). A local (i.e., near \(n_c/4\)) electron temperature measurement can be obtained from this feature from the relation \(T_e, \text{keV} = \Delta \lambda / 3.09\) (Ref. 28), where \(\Delta \lambda\) is the shift of the spectral peak from 2\(\lambda_0\) after corrections for Doppler and Dewandre shifts\(^{29}\) have been applied.\(^{30}\) The electron temperature inferred from this technique is \(T_e = 4.5 \pm 0.2 \text{ keV}\). The DRACO calculations predict a consistent temperature (4.5 keV), giving confidence in the numerical modeling of the corona and indicating that ignition-relevant temperatures have been achieved. As a result of refraction effects, this spectral feature is emitted only perpendicularly to the density gradient (i.e., along the target normal),\(^{19}\) and its observation required that the target be tilted to face the diagnostic [Fig. 151.22(d)]. For this reason, it is not seen in Figs. 151.22(b) or 151.22(c).

Importantly, this feature demonstrates significant differences relative to the near-2\(\lambda_0\) spectrum obtained at smaller scales on OMEGA. A typical half-harmonic spectrum from a spherical implosion experiment (shot 80802) on OMEGA is shown in the inset of Fig. 151.22(a), sharing the same wavelength scale as the NIF spectrum. The characteristic half-harmonic features that are red- and blue-shifted with respect to 2\(\lambda_0\) seen

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**Figure 151.22**

Time-resolved scattered-light spectra at collection angles of (a) 0°, (b) 23.5°, and (c) 50° relative to the target normal. These images were obtained in two CH target experiments. The image in (a) corresponds to an experiment (d) with the target oriented toward a streaked spectrometer and (e) irradiated by a ramp-flat pulse at a peak quarter-critical laser intensity of \(1.3 \times 10^{15} \text{ W/cm}^2\). The images in (b) and (c) correspond to an experiment (f) with the target oriented toward the south pole of NIF and was (g) irradiated first by beams at incidence angles of 45° and 50°, followed by beams at 23° and 30°. The streaked spectrum from a spherical-geometry experiment on OMEGA [inset in (a)] is contrasted to the image in (a).

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in the OMEGA experiment are a definitive diagnostic of the presence of TPD.\textsuperscript{31} The doublet arises by processes such as inverse resonance absorption, inverse parametric decay, and self-Thomson scattering that convert the up- and down-shifted TPD daughter plasma waves into transverse (light) waves.\textsuperscript{28} The lack of a blue-shifted half-harmonic and the narrowness of the red-shifted feature seen in the NIF experiment is a strong indication that different physical processes are occurring at the quarter-critical surface. The sharp feature observed in the NIF experiment is a well-known signature of the absolute Raman instability that can occur at densities close to quarter critical.\textsuperscript{28} The OMEGA spectrum implies the absence of SRS around \(n_T/4\) and the presence of TPD, while the NIF spectrum implies the presence of SRS at and below \(n_T/4\). Although the presence of some TPD activity in the NIF experiment cannot be entirely ruled out on the basis of Fig. 151.22(a) since the conversion efficiencies of TPD waves to half-harmonic emission relative to absolute SRS are difficult to quantify, it seems most plausible that SRS, rather than TPD, is the dominant quarter-critical LPI mechanism in ignition-scale direct-drive experiments.

Simple considerations based on the absolute threshold intensities for SRS \((I_{\text{SRS,thr}} = 2377/T_{n,\mu m}^{1.3})\) and TPD \((I_{\text{TPD,thr}} = 233 T_{e,keV}^{-1}/L_{n,\mu m}^{1.6})\), for normally incident single plane-wave beams,\textsuperscript{9,10} further support this identification. In these expressions, \(I_{\text{thr}}\) is the threshold intensity in units of \(10^{14}\) W/cm\(^2\). As an illustrative case, Fig. 151.23 shows the ratio of the absolute TPD threshold to twice that for absolute SRS as a function of electron temperature and density scale length. This is intended to acknowledge the fact that while TPD has been observed to be a multibeam phenomenon, it may be the case that fewer beams contribute to SRS. The OMEGA experiment that produced the spectrum shown adjacent to Fig. 151.22(a) \(L_n \sim 500 \mu \text{m}, T_e \sim 2.5 \text{ keV, } I \sim 6 \times 10^{14} \text{ W/cm}^2\) is marginally unstable with respect to TPD and slightly less so to SRS if the total overlapped laser intensities are substituted into the expressions for the single-beam thresholds. In contrast, the NIF experiment at ignition-relevant conditions \(L_n \sim 525 \mu \text{m}, T_e \sim 4.5 \text{ keV, } I \sim 1.3 \times 10^{15} \text{ W/cm}^2\), which produced the spectrum shown in Fig. 151.22(a), is in the SRS-dominated regime: the threshold for SRS is exceeded by a factor of ~22, while the TPD threshold is exceeded by a factor of ~6. It is expected that this qualitative trend of SRS being increasingly prominent relative to TPD with increasing scale length and temperature\textsuperscript{32} applies also for more-complicated cases of multiple obliquely incident beams, although this is a subject of future work.

The broad spectral features seen in Figs. 151.22(a)–151.22(c) are characteristic of SRS occurring at densities below \(n_T/4\) (between 0.15 and 0.22 \(n_T\)). Figures 151.22(b) and 151.22(c) highlight SRS spectra obtained at two different angles of observation and two distinct irradiation conditions. The target normal was parallel to the NIF polar axis [Fig. 151.22(d)], and the target was irradiated symmetrically, first by the outer beams from \(t = 0\) to \(t = 4.5\) ns, followed by the inner beams from \(t = 4.5\) ns to \(t = 7.5\) ns [Fig. 151.22(g)]. The predicted quarter-critical plasma conditions during the outer (inner) beam drive were \(L_n \sim 500 \text{ (690) } \mu \text{m, } I \sim 1.6 (1.1) \times 10^{15} \text{ W/cm}^2\), and \(T_e \sim 4.7 (4.4) \text{ keV, respectively. Temporally resolved scattered-light spectra\textsuperscript{26} were obtained at 23.5° [Fig. 151.22(b)] and 50° [Fig. 151.22(c)]. SRS is observed by both diagnostics at early times during outer-beam irradiation and at later times when irradiated by the inner beams.

This observation is attributed to SRS sidescatter,\textsuperscript{33} for which newly developed theory and supporting simulations are described in a companion manuscript.\textsuperscript{34} In this process the SRS light waves propagate approximately tangentially to contours of constant electron density in the corona and see much greater gains relative to backscatter. The data shown in Figs. 151.22(b) and 151.22(c) are in agreement with the predictions of this theory and cannot be explained by narrow-angle backscatter simply caused by refraction, particularly for the SRS observed at 50°. Therefore, the propagation direction (and collection angle) of SRS light, after it has finished refracting and is in vacuum, is determined solely by its wavelength (i.e., the density where it was generated) and depends only weakly on the incidence angle of the beams that produced it. This is evident in Fig. 151.22(b), where SRS light at 23.5° is observed.
at \( \sim 650 \text{ nm} \) during both outer-beam and inner-beam irradiation. The SRS shifts to shorter wavelengths (\( \sim 620 \text{ nm} \)) when the observation angle is moved to 50°.

To determine the total amount of SRS generated in these experiments, absolutely calibrated photodiodes measured the SRS light collected in \( \sim 2 \times 10^{-3} \text{ sr} \) in the two full-aperture backscatter stations (FABS’s)\(^{26} \) at 50° and 30°. These measurements were then extrapolated to account for the total emission. This was accomplished using a ray-tracing code with parameters and geometry provided by DRACO simulations to obtain simulated SRS emission profiles that include refraction and absorption as functions of wavelength and angle of observation (transmission of SRS light from its origin ranges from 2% at 702 nm to \( \sim 50\% \) at 630 nm). These calculations assume \( 2\pi \) azimuthal symmetry around the target normal. With the above assumptions, it is estimated that between 2% and 6% of incident laser energy is converted to SRS light.

The inferred SRS light energy is compared to the energy in hot electrons, which is inferred from hard x-ray bremsstrahlung emission generated by the interaction of hot electrons with the target.\(^{35} \) This bremsstrahlung emission was detected using the NIF filter fluoroscencer (FFLEX) diagnostic.\(^{36} \) The FFLEX signals were analyzed by performing Monte Carlo electron–photon transport calculations with the EGSnrc code,\(^{37} \) using a single-temperature (\( T_{\text{hot}} \)) 3-D Maxwellian hot-electron distribution. These calculations relate the absolute intensity of hard x-ray emission to the total quantity of hot electrons that produce it. Figure 151.24 shows the corresponding fraction of laser energy converted to hot electrons (\( f_{\text{hot}} \)) as a function of laser intensity at the quarter-critical density as calculated by DRACO for a series of experiments that include both CH and Si targets. The hard x-ray data were integrated over the period of the experiment starting after 4.5 ns. For outer-beam irradiation, \( f_{\text{hot}} \) increased from 0.7±0.2% to 2.9±0.6% as the laser intensity increased from \( 5.9 \times 10^{14} \text{ W/cm}^2 \) to \( 14 \times 10^{14} \text{ W/cm}^2 \). For inner-beam irradiation of CH targets, \( f_{\text{hot}} \) increased from 1.2±0.2% to 2.6±0.5% for intensities of \( 6.2 \times 10^{14} \text{ W/cm}^2 \) to \( 11 \times 10^{14} \text{ W/cm}^2 \). The uncertainty in \( f_{\text{hot}} \) is based on the statistical uncertainty in the single-temperature fit to the hard x-ray spectra. For CH experiments, \( T_{\text{hot}} \) is inferred to be between 45 and 55 keV for the outer-beam drive and 62 keV for the inner-beam drive, independent of laser intensity, with an uncertainty of ±4 to 5 keV. The threshold intensity for the onset of measurable hot electrons in CH targets lies in the vicinity of \( 4 \times 10^{14} \text{ W/cm}^2 \).

The inferred energy and temperature of the hot electrons are consistent with simple arguments based on SRS being their source. By conserving wave action in the scattering process (i.e., the Manley–Rowe relations\(^{38} \)), it was determined that, for SRS wavelengths between 600 and 650 nm, the total energy in plasma waves is 70% to 85% of the total energy in SRS between 1.4% and 5% of the incident laser energy for the experiments shown in Fig. 151.22. It is quite plausible that kinetic mechanisms such as wave breaking or stochastic processes can convert the plasma-wave energy into hot electrons with an efficiency sufficient to account for the fraction that is observed (\( f_{\text{hot}} = 1\% \) to 3%). The characteristic temperature for SRS-generated electrons is often estimated by \( T_{\phi} = \frac{(1/2) m_e v_\phi^2}{\rho} \) (Ref. 8), where \( v_\phi \) is the phase velocity of the plasma wave. For our experiments, where SRS is observed from wavelengths of \( \sim 620 \text{ nm} \) to \( \sim 702 \text{ nm} \) (2\( \lambda_p \)), the corresponding hot-electron temperatures range from \( \sim 30 \) to \( \sim 85 \text{ keV} \) (\( T_{\phi} \sim m_e c^2/6 \) for \( n_z = n_e / 4 \)), which is consistent with the hot-electron temperatures that best fit the measured hard x-ray spectrum.

The combination of \( T_{\text{hot}} \) and \( f_{\text{hot}} \) inferred in these experiments is close to the level that can be permitted in direct-drive–ignition designs, typically considered to be \( f_{\text{hot}} \sim 0.5\% \) to 1% for \( T_{\text{hot}} \sim 50 \text{ keV} \) (Refs. 2 and 39). This estimate is based on an allowable coupling of \( \sim 0.1\% \) of laser energy to hot-electron preheat in the DT fuel and a near-2\( \pi \) angular divergence of hot electrons inferred in OMEGA spherical experiments.\(^{20} \) Based on these data, direct-drive–ignition designs using a CH ablator and quarter-critical laser intensities of \( \sim 5 \times 10^{14} \text{ W/cm}^2 \) may be acceptable, but for higher intensities, LPI mitigation is likely to be necessary. The discovery of a regime dominated by SRS, rather than by TPD as on OMEGA, necessitates a
re-evaluation of the angular divergence of hot electrons at direct-drive ignition-relevant conditions and may also require reconsideration of mitigation strategies.

One potential LPI mitigation strategy, originally proposed for TPD, uses strategically placed mid-Z layers in the ablator to locally shorten the density scale length, increase the electron temperature, enhance electron–ion collisional damping, and reduce Landau damping of ion-acoustic waves.\(^{40–45}\) This reduction in scale length and increase in temperature are predicted as well for planar Si experiments \((L_n \approx 690 \mu m \text{ in CH to } \sim 560 \mu m \text{ in Si}; T_c \approx 4.4 \text{ keV in CH to } \sim 5.2 \text{ keV in Si})\), for which hot-electron data are shown in Fig. 151.24. The use of Si ablators has a modest effect on hot-electron levels, although it does increase the hot-electron intensity threshold to around \(6 \times 10^{14} \text{ W/cm}^2\). The lack of hot electrons in this experiment also correlates with a minimal level of SRS observed in any of the spectrometers.

In summary, the first experiments to investigate LPI at direct-drive ignition-relevant coronal plasma conditions have revealed evidence of a regime dominated by SRS, with a significant contribution from tangential sidescatter. This result is in stark contrast to prior experiments on OMEGA at shorter scale lengths and lower temperatures, in which SRS was minimal and quarter-critical instabilities were identified as TPD. For the first time, intensity thresholds for LPI hot electrons have been evaluated at direct-drive–ignition scales, and the use of a Si ablator has been found to increase the threshold intensity slightly, from \(\sim 4 \times 10^{14} \text{ W/cm}^2\) to \(\sim 6 \times 10^{14} \text{ W/cm}^2\). These quarter-critical laser intensities present a viable design space for direct drive. As discussed, these results have implications for LPI hot-electron preheat mitigation in direct-drive–ignition designs, which traditionally have included strategies to mitigate TPD, but will have to consider SRS. In future experiments, it will be important to characterize the angular distribution of hot electrons, which strongly affects the tolerable level of hot-electron generation and may be different in this SRS-dominated regime than in TPD-dominated experiments on OMEGA.\(^{20}\) Optical Thomson scattering will ultimately be used on the NIF\(^{46,47}\) to directly probe and characterize plasma waves in the quarter-critical region, as has been done previously on OMEGA,\(^{16}\) in order to definitively assess the presence or absence of TPD.

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REFERENCES


30. The Doppler shift resulting from plasma flow affects the wavelength of both the incoming 351-nm light and the outgoing 3912 light, and constitutes an ~10% effect on the wavelength. The Dewandre shift results from plasma expansion and an increase in the plasma density between the quarter-critical region and the observer as 351 nm propagates outward, and is an ~2% effect on the wavelength.


Measurement of Cryogenic Target Position and Implosion Core Offsets on OMEGA

Introduction
Direct-drive inertial confinement fusion experiments are performed by uniformly illuminating spherical cryogenic deuterium–tritium (DT) fuel-bearing CH shell capsules with high-power laser beams. A goal of inertial confinement fusion (ICF) experiments is thermonuclear ignition and gain; in order for this to occur, the fuel must be symmetrically compressed to high areal densities, i.e., at least 0.3 g/cm², and the central hot-spot temperature must be at least 10 keV (Refs. 2–4). Target performance is degraded by imperfections in symmetric laser illumination and in the target itself. Perturbations of the intensity in the low Legendre modes (\(\ell \leq 6\)), which may include target offset, can distort the core at stagnation while higher-mode (\(\ell \geq 6\)) perturbations lead to Rayleigh–Taylor unstable growth, target breakup, and mixing of the materials in the shell and fuel. These perturbations reduce the peak temperature and areal density of the final fuel region; therefore, minimizing them is desired. Assessing the performance of the implosions requires one to simulate the implosion with one-dimensional (Ref. 6) and multidimensional hydrocode simulations, and the multidimensional simulations require accurate values of target offset from beam aiming to accurately simulate the implosions. This article describes our method of measuring initial target offset from the aim point of the beams and determining the core offset resulting from target offset from this aim point.

Measurements of Initial Target Offsets
Targets illuminated by the 60 beams of OMEGA at intensities ranging from \(\sim 10^{14}\) to \(\sim 10^{15}\) W/cm² emit x rays, easily imaged by pinhole cameras, in the range of 1 to 10 keV. A set of x-ray pinhole cameras (XRPC’s) is used on OMEGA to precisely align the laser beams to the target center. This is currently done to an accuracy of \(\sim 7\mu\)m rms (root mean square) using a set of fixed and retractable XRPC’s, all digitally recording the images with charge-injection–device (CID) cameras. This set includes five fixed XRPC’s, which are attached to the OMEGA target chamber, and six ten-inch-manipulator (TIM)–based XRPC’s, which are retractable through a vacuum gate valve. The fixed XRPC’s remain in use during both cryogenic target and non-cryogenic target implosions.

Since the targets used to precisely align the OMEGA beams by locating the x-ray spot emitted by each beam on pointing shots are positioned by visible light cameras, all other non-cryogenic targets are aligned to this same point. This position is referred to as target chamber center (TCC), although it is really the aim point of the beams determined through precision pointing using the CID-based XRPC’s. The precise locations of TCC in the XRPC images are determined by measuring x rays emitted by a precisely located non-cryogenic target; these TCC reference images are all from target shots taken on the same day as the cryogenic target shots. This effectively eliminates the possibility of changes in the XRPC’s contributing to the determined offsets. Positioning cryogenic targets is complicated by the need to view the cryogenic target through windows in the shroud that maintains the target at near the triple point (\(\sim 20\) K). These windows refract the light passing through them by an amount that must be determined by testing prior to the actual shots. Furthermore, vibration of the cryogenic target stalk while the shroud is in place and impulses transmitted to this stalk when the shroud is retracted (\(\sim 50\) ms before the shot) can misplace the cryogenic target. This can cause the cryogenic target to be offset from TCC at \(t_0\) (the beginning of the laser pulse).

The example of a non-cryogenic target XRPC image in Fig. 151.25(a) shows a 1.5 × 1.5-mm region at the target plane. The outer edge of the x-ray emission, which occurs at \(t_p\), is determined from the maximum gradient using a Sobel filter. This set of positions is then fit to a circle whose center position is then determined [overlaid circle and central cross in Fig. 151.25(a)]. An example cryogenic target x-ray image is shown in Fig. 151.25(b). The fusion neutrons created by the implosion of this target (\(\gamma_n = 3 \times 10^{13}\)) have generated copious amounts of noise including a gamma-ray–induced background, single-pixel upsets caused by neutron-scattering events that produce protons, and line and column noise caused by similar interactions with the readout structure of the CID camera. This noise can in large part be removed by first filtering the image using a single-pixel upset detection and replacement algorithm, next by removing the average line and column noise measured away from the image itself, and lastly by using a...
median filter to reduce additional noise. The result of performing this noise removal procedure is shown in Fig. 151.25(c), and the position of the center of the cryogenic target is determined in like fashion to the reference non-cryogenic target. The pixel-location differences of the two centers are then converted to microns, and the difference between the cryogenic target position and the reference target position is a measured projected offset at $t_0$.

Two methods are employed to determine the three-space offset $\mathbf{r}$ of cryogenic targets at $t_0$ from TCC. Both methods use the projected offsets of the cryogenic target centers at $t_0$ from the reference non-cryogenic targets whose centers are at TCC. The view vectors for each XRPC are related to the target chamber vector coordinates by the following formulas:

\[
\mathbf{q} = \frac{\mathbf{z} \times \mathbf{v}}{|\mathbf{z} \times \mathbf{v}|} \quad \text{(1)}
\]

\[
\mathbf{p} = \frac{\mathbf{v} \times \mathbf{q}}{|\mathbf{v} \times \mathbf{q}|} \quad \text{(2)}
\]

where $\mathbf{q}$ is the horizontal vector in an image whose view direction is $\mathbf{v}$ and the normalized cross product of $\mathbf{z}$ (straight up) and $\mathbf{v}$, while the vertical direction in the image plane $\mathbf{p}$ is given by the normalized cross product of $\mathbf{v}$ and $\mathbf{q}$ (see Fig. 151.26).

The XRPC’s provide multiple quasi-orthogonal views of the target x-ray emission, from which $\mathbf{r}$ can be determined. The first method uses the projected offsets from pairs of cameras to determine the three-space offsets. For an offset in space of $\mathbf{r}$, the projections of $\mathbf{r}$ in a pair of camera views are given by

\[
\mathbf{r} \cdot \mathbf{q}_1 = r_x q_{x_1} + r_y q_{y_1} + r_z q_{z_1}, \quad \text{(3)}
\]

\[
\mathbf{r} \cdot \mathbf{p}_1 = r_x p_{x_1} + r_y p_{y_1} + r_z p_{z_1}, \quad \text{(4)}
\]

\[
\mathbf{r} \cdot \mathbf{q}_2 = r_x q_{x_2} + r_y q_{y_2} + r_z q_{z_2}, \quad \text{(5)}
\]

\[
\mathbf{r} \cdot \mathbf{p}_2 = r_x p_{x_2} + r_y p_{y_2} + r_z p_{z_2}, \quad \text{(6)}
\]

where 1 and 2 refer to the first and second view, respectively. The results can be combined into two different matrices by choosing to solve for $\mathbf{r}$ using either Eqs. (3), (4), and (5) or Eqs. (3), (4), and (6). This is equivalent to using the vertical offset from either camera 1 or 2. These choices can be written in matrix form as follows:

![Figure 151.25](E26466JR)

**Figure 151.25**

X-ray images from the OMEGA H4 x-ray pinhole camera (XRPC) charge-injection device (CID). (a) Target chamber center (TCC) reference image on shot 85780, (b) unfiltered image from a cryogenic target on shot 85784, and (c) filter image of the same. Both the reference image (a) and the filtered cryogenic target shot image (c) have the best-fit positions indicated by a circle and a cross in the center.

![Figure 151.26](E26467JR)

**Figure 151.26**

Vector representation displaying view direction, solution direction, and unit vectors of image plane with respect to each other.
Inverting the matrices gives two possible solutions to the offset as follows:

\[
\begin{align*}
\bar{r}_{121} &= \mathbf{M}_{q_1q_2p_1}^{-1} \mathbf{r}_{q_1q_2p_1}, \\
\bar{r}_{122} &= \mathbf{M}_{q_1q_2p_2}^{-1} \mathbf{r}_{q_1q_2p_2}.
\end{align*}
\]

The average of these two solutions is the choice used in this method and is given by

\[ \mathbf{r} = \frac{(\bar{r}_{121} + \bar{r}_{122})}{2}. \]

To improve the accuracy of determining \( \mathbf{r} \), values are computed from as many quasi-orthogonal camera pairs as possible. The results are averaged and the standard deviations of the values are used as an estimate of the errors of these values.

The second method of determining \( \mathbf{r} \) uses a least-squares approach. For a given assumed offset of the target \( \mathbf{r} \), the values \( r_{q_i} \) and \( r_{p_i} \) that would be observed in the \( i \)th view are given by

\[
\begin{align*}
\hat{r}_{q_i} &= r_{q_i} + r_{q_y} q_y + r_{q_z} q_z, \\
\hat{r}_{p_i} &= r_{p_x} p_x + r_{p_y} p_y + r_{p_z} p_z.
\end{align*}
\]

The least-squares search is performed to minimize the quantity \( \chi^2 \) given by

\[
\chi^2 = \sum_i (\Delta r_{i,j})^2 w_i^2,
\]

where the values \( w_i \) are the weights given to the \( i \)th view and the quantities \( \Delta r_{i,j} \) are the perpendicular offsets in the \( i \)th view in turn given by

\[
\Delta r_{i,j} = \sqrt{(\hat{r}_{q_i} - \hat{r}_{q_i})^2 + (\hat{r}_{p_i} - \hat{r}_{p_i})^2}.
\]

where \( \Delta q_i \) and \( \Delta p_i \) are the horizontal and vertical offsets of the target in the \( i \)th view. The value of \( \mathbf{r} \) that minimizes \( \chi^2 \) is taken as the best value, while the error \( d \mathbf{r} \) is given by

\[
d\mathbf{r} = \left( \frac{\chi^2}{\sum w_i} \right)^{1/2} \]

and is equivalent to the error of the mean of the best-fit value.

When only two views are available, the first method of determining \( \mathbf{r} \) is the best method to use, whereas when more than two views are available, the second method gives the most unbiased result. Table 151.I shows the current set of fixed XRPC’s used in this position analysis. Typical errors when determining the position are ~3 to 5 \( \mu \)m.

<table>
<thead>
<tr>
<th>Camera</th>
<th>Position ( \theta (\degree) )</th>
<th>Position ( \phi (\degree) )</th>
<th>Magnification</th>
</tr>
</thead>
<tbody>
<tr>
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<td>54.00</td>
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<tr>
<td>h13</td>
<td>9.74</td>
<td>342.00</td>
<td>4.043</td>
</tr>
<tr>
<td>p2</td>
<td>68.43</td>
<td>54.00</td>
<td>3.992</td>
</tr>
</tbody>
</table>

Table 151.I: X-ray pinhole camera (XRPC) parameters.

*Before 17 March 2017

**After 17 March 2017

Measurement of Implosion Core Offsets

The implosion cores are imaged by the gated monochromatic x-ray imager (GMXI)\(^1\) operating in time-integrating mode with four CID cameras recording the four images formed by the Kirkpatrick–Baez (KB) microscope optical assembly. Two of these images (GMXI-c and GMXI-d, filtered by 50.8 and 76.2 \( \mu \)m of Al, respectively) had signal levels that did not exceed the capacity of the CID cameras for all cryogenic and non-cryogenic target experiments that determined reference core positions. As for determining the \( t_0 \) offset, the non-cryogenic reference target is assumed to be perfectly centered at TCC. The energy bands are approximately the same for these two images being ~5 to 8 keV and ~5.5 to 8 keV for images GMXI-c and GMXI-d, respectively. The GMXI cameras observe the implosion cores from the common spherical coordinates \( \theta = 96.02°, \phi = 54° \) with respect to the target plane (each being ~1° away).

In contrast to the \( t_0 \) images where the limb of the image is used to determine the center, core images are centrally peaked, so the centroid is a better measure of the core’s position in the CID image. Figure 151.27 shows example GMXI images, trimmed to \( 200 \times 200 \mu \)m. The reference images used in this
case are from a target experiment with a non-cryogenic target consisting of an 18-μm-thick CH shell filled with 3 atm of D₂ gas, imploded with the same pulse shape used on the subsequent cryogenic target shots (these are referred to as pulse-shape setup (PSS) shots). Figures 151.27(a) and 151.27(b) are from the reference non-cryogenic target implosion (OMEGA shot 81056, GMXI-c and GMXI-d images); Figs. 151.27(c) and 151.27(d) are from a cryogenic target implosion (OMEGA shot 81060, GMXI-c and GMXI-d images). The “x” symbols denote the centroids of the images; the diamond symbols on the GMXI-c and GMXI-d images of OMEGA shot 81060 show the points of maximum cross correlation between the cryogenic and PSS shots in the GMXI-c and GMXI-d images, respectively. The square symbols on shot 81060 images denote the averages between the centroid centers and the cross-correlation maximums. The amount of core offset is taken as the amount by which the image must be shifted to maximize the cross correlation. Figure 151.28 shows the core offsets determined from the GMXI-c and GMXI-d images for a large number of cryogenic target shots; their consistency is evident. The average offset of the GMXI-c and GMXI-d images is therefore taken as the offset and the difference is an estimate of the error of this offset.

The \( t_0 \) offsets are compared with the GMXI offsets by computing the projections of the \( t_0 \) offsets in the view of the GMXI in the horizontal and vertical directions \( q_{0,GMXI} \) and \( p_{0,GMXI} \), respectively, given by

\[
q_{0,GMXI} = r \cdot q_{GMXI},
\]

\[
p_{0,GMXI} = r \cdot p_{GMXI},
\]

where \( q_{GMXI} \) and \( p_{GMXI} \) are the horizontal and vertical vectors, respectively, of the GMXI view. Since there is no other digitally recorded view of the core, the three-space core offset cannot be determined but the GMXI core offset and the projection of the \( t_0 \) offset into the GMXI view can be compared.

**Results**

Figures 151.29 and 151.30 show the measured core offsets compared to the projected \( t_0 \) offsets for horizontal, vertical, and
radial directions, respectively. The offsets in Fig. 151.29 are for a quasi-uniformly distributed sample of cryogenic target shots that span offsets from near zero to >100 μm and whose offset directions were nearly perpendicular to the GMXI view direction. Figure 151.29(a) shows that the horizontal displacement of the core is in the same direction as the \( t_0 \) offset and nearly equal in magnitude; i.e., the core is forming at approximately the position of the offset target center. In large part the core offsets confirm the accuracy of the \( t_0 \) offsets. Figure 151.30 compares the horizontal, vertical, and radial offsets of the implosion cores and the projected \( t_0 \) offsets for all recent cryogenic target shots (since 2015). The horizontal components of the \( t_0 \) and core offsets [Fig. 151.30(a)] are approximately uniformly distributed about the origin and most are <20 μm. The few large horizontal offsets agree in direction and are nearly of the same magnitude. In contrast, the vertical \( t_0 \) offsets are biased toward positive offset (in this case, from the TCC reference), whereas the core vertical offsets (y axis) are more uniformly distributed between positive and negative values. The reason for the positive bias of the \( t_0 \) vertical target offset is not known, but it is suggestive that the cryogenic targets are systematically above TCC at \( t_0 \) with an average offset of ~5 μm.

A large offset is expected to have a very detrimental effect on the fusion neutron yield, and even small offsets are calculated to have an effect on the yield under ideal simulated conditions,\(^{13}\) so placing the target at TCC as accurately as possible is desired. But in real experimental conditions where many other factors may affect the implosion in addition to target offset, it may be difficult to assess the importance of target offset alone. To explore this dependence, the measured neutron yield divided by the calculated yield [yield-over-clean (YOC)] by the 1-D hydrocode \( \text{LILAC}^6 \) is plotted in Fig. 151.31 as a function of the measured \( t_0 \) offset for all recent cryogenic target shots (since 2015). Figure 151.31 shows that the YOC varies from ~0.2 to 0.7 for offsets less than ~15 μm and is smaller (~0.3 or less)
for offsets greater than ~20 μm. These results are consistent with requiring a small offset to get a large value of the YOC but that a smaller value may be obtained at a small initial target offset for other unrelated reasons.

![Graph](image)

Figure 151.31
The measured cryogenic target neutron yield divided by the calculated yield ratio [yield-over-clean (YOC)] plotted as a function of the offset of the target at $t_0$.

Conclusions
This work describes the method for determining the offsets of cryogenic targets relative to the aim point of the OMEGA laser beams ($t_0$ offset) and shows measurements of the implosion core offsets from well-centered targets as determined in one direction (that of the GMXI). The $t_0$ offsets projected in the direction of the GMXI agree in direction and are close in magnitude to that of the core offsets with considerable scatter at small $t_0$ offsets (<20 μm). The approximate dependence of the YOC on target offset is such that no large YOC’s are obtained when the $t_0$ offset is large (>20 μm). Knowing the accurate value of the $t_0$ offset is therefore critical in assessing the fusion performance of the implosion.

ACKNOWLEDGMENT
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REFERENCES
Subpercent-Scale Control of 3-D Modes 1, 2, and 3 of Targets Imploded in a Direct-Drive Configuration on OMEGA

In laser-driven implosion experiments, a laser illuminates a spherical target either directly (direct-drive configuration) or after conversion into x rays (indirect-drive configuration). This absorbed laser energy leads to the ablation and extreme acceleration of the outer surface of the target through the “rocket effect.” This method is widely used to study plasma physics including high-energy-density physics and inertial confinement fusion (ICF). In all cases, maintaining spherical symmetry throughout the implosion is critical to obtaining a 1-D behavior that maximizes the internal energy of the imploded plasma at final compression. In ICF experiments, a capsule filled with deuterium (D) and tritium (T) is used to create a self-sustained fusion burn that will ignite the fuel and produce a net energy gain. At the end of the implosion, the kinetic energy of the imploding capsule is converted into internal energy, triggering the fusion reaction during stagnation.

Several simulations and comparisons with experiments have shown that target low-mode nonuniformities lead to a severe reduction in the implosion performance because of increased residual kinetic energy during stagnation and uneven compression that result in reduced core pressure and truncated burn. This degradation was shown to be particularly significant for modes $\ell \leq 3$, where $\ell$ is the order of the modes of the spherical harmonic decomposition of the shell’s shape. Consequently, reducing low-mode nonuniformity has been identified as one of the most-critical steps in demonstrating ignition at the National Ignition Facility (NIF) or conditions that are hydrodynamically equivalent to ignition when scaled from 26-kJ implosions on OMEGA to megajoule energies on the NIF.

Over the last decade, many studies have shown significant low-mode asymmetries of the imploding shell. Modes $\ell = 1$ have been typically identified from properties of the final assembly including asymmetry in its areal density, variation of its ion temperature along different lines of sight, hot-spot motion, and asymmetric x-ray emission of a Ti layer embedded at the inner surface of the shell. Modes $\ell \geq 2$ have been measured from the hot-spot shape, standard or Compton radiography, x-ray absorption spectroscopy, and self-emission shadowgraphy.

Several studies have focused on the causes of the asymmetries and the development of methods to correct them. In indirect-drive ICF, the laser’s beam-energy balance was modified to exploit cross-beam energy transfer and improve the sphericity of the core emission. The improvement was limited, however, because the observable (i.e., the core shape) was restricted to modes $\ell \geq 2$ and too indirect to give accurate access to the 3-D structure of the shell. In direct-drive ICF, simulations have identified different potential effects that create nonuniformities including target offset, beam-power imbalance, beam pointing, and beam timing. Success has been limited, however, in reproducing the experimental observables obtained on OMEGA because of the difficulty in evaluating and modeling each effect.

This article reports the first experimental demonstration that the amplitude of modes $\ell = 1, 2,$ and 3 of targets imploded in direct-drive configuration on OMEGA measured at a convergence ratio of $\sim 3$ can be controlled within $\pm 0.25\%$ by adjusting the laser’s beam-energy balance, leading to a total radial error of $1\%$. Over three shots, the 3-D shape of the imploding target was tomographically recorded by measuring four lines of sight of the ablation front. The projected ablation-front contours were measured with framing cameras using the x-ray self-emission shadowgraphy technique. The projected ablation-front motions were obtained by comparing the positions of the contours on the framing cameras with the corresponding contour positions measured on a nonimploding solid-CH-ball shot. The amplitudes of the modes were determined within $\pm 0.15\%$ by decomposing the contours oriented perpendicular to the lines of sight and shifted by the measured motions over spherical harmonics. The variations of the normalized target mode amplitudes ($\Delta r^m_\ell$, where $m$ is the mode order) between shots were shown to change linearly (within $\pm 0.25\%$) with the variation of the normalized mode amplitudes of the laser’s beam-energy balance ($\Delta e^m_\ell$) with a low-mode coupling coefficient ($C_\ell = \Delta r^m_\ell / \Delta e^m_\ell$) of $C_1 = -0.663 \pm 0.05$, $C_2 = -0.38 \pm 0.05$, and $C_3 = -0.18 \pm 0.03$ for modes $\ell = 1, 2,$ and 3, respectively. The decrease of $C_\ell$ with increasing mode number was expected because of the phase plates used with each beam on OMEGA.
Lateral thermal transport and amplification by the Rayleigh–Taylor instability were not expected to be important because of the long spatial wavelength (\(\lambda = 2\pi R/\ell\), where \(R\) is the averaged shell radius) of the low modes. The \(C_\ell\) values enabled one to evaluate within \(\pm 0.05\%\) of the amplitudes of the residual target modes that appear when the laser’s beam energies are balanced and to determine the laser mode amplitudes that mitigate them within \(\pm 0.25\%\).

The experiments employed 60 ultraviolet (\(\lambda_0 = 351\) nm) laser beams on the OMEGA Laser System.\(^{30}\) The beams illuminated the target and were smoothed by polarization smoothing,\(^{31}\) shaped by spectral dispersion,\(^{32}\) and distributed phase plates (fourth-order super-Gaussian with 95% of the energy contained within the initial target diameter).\(^{33}\) A 2-ns-long square pulse irradiated 866±3-\(\mu\)m-diam capsules with an energy of 20.2±0.4 kJ, resulting in an intensity \(\approx 4.3 \times 10^{14} \text{ W/cm}^2\). The shells were made of 19.2±0.2-\(\mu\)m-thick glow-discharge polymer (CH with a density of 1.03 mg/cm\(^3\) and each mode amplitude \(\leq 50\) nm) and filled with 17±1.5 atm of deuterium. An additional reference shot was made on an 856-\(\mu\)m-diam solid CH ball. For each shot, the target was placed at target chamber center with a maximum radial error of 1.5 \(\mu\)m measured with two high-speed video cameras (1000 images per second) that were used to automatically position the target before the shot.

The first shot used a standard laser beam-energy balance with a standard deviation of 2.5%. On the second and third shots, the beam-energy balance was varied to change the amplitude of the evaluated laser modes by minimizing

\[ \sum_{b=1}^{60} \left[ \sum_{l=0}^{\infty} \sum_{m=-l}^{l} \sqrt{4\pi e^m \ell^m} Y^m_l(\theta_b, \phi_b) - \bar{E}_b \right]^2 \]

with a larger variation for modes \(m = 0\), where \(\bar{E}_b\) is the energy of the beams normalized to averaged beam energy in percent, \((\theta_b, \phi_b)\) are the coordinates of the OMEGA beam ports, and \(Y^m_l(\theta, \phi)\) are the tesseral spherical harmonics.\(^{34}\) On the third shot, one beam was reduced by 80% as a result of hardware malfunction, further amplifying the mode amplitudes. The beam energies were measured with integrated spheres within \(\delta \bar{E}_b = \pm 0.5\%\) that were absolutely calibrated within \(\pm 2\%\) with calorimeters. This resulted in the same relative error for all mode amplitudes of \(\delta (e^m_l) = \delta \bar{E}_b/\sqrt{N_b} = \pm 0.06\%\), where \(N_b = 60\) is the number of beams.

Four x-ray framing cameras, located at different lines of sight, used arrays of 16 pinholes to image the soft x rays emitted by the irradiated target on four strips of a microchannel plate (MCP).\(^{35}\) The cameras were set up to magnifications of \(M = 6\) (two cameras) and \(M = 4\) (two cameras) with pinhole sizes of 10 \(\mu\)m and 15 \(\mu\)m, respectively. Their point-spread functions (PSF’s) result in about 2-D Gaussian convolutions of the images with a full width at half maximum of \(d_{PSF} \approx 10 \mu\)m and \(d_{PSF} \approx 15 \mu\)m, respectively.\(^{36}\) Four short, high-voltage pulses were sent to each strip to activate the signal amplification by the MCP and obtain time-resolved images. For all imploding shells, the electrical pulses were timed to ~0.4 ns, ~1.2 ns, ~1.5 ns, and ~1.8 ns, whereas for the reference shots, they were synchronized to ~0.2 ns after the beginning of the laser pulse (defined as 1% of maximum intensity). The absolute timings between the laser pulse and the images were known to an accuracy of 20 ps, and the interstrip timings were determined within 5 ps (Refs. 26 and 36). Along each strip, the images were separated by ~50 ps. Three cameras had integration times of 40 ps; one had an integration time of 200 ps. On all cameras, 25.4-\(\mu\)m-thick Be filters were used to record the soft x rays above ~1 keV. For each camera, the same pinhole array was used on all shots to maintain the distance between images.

On each self-emission image, the inner edge contour of the intensity peak [Figs. 151.32(a) and 151.32(b)] corresponded to the projection of the ablation-front surface along the line of sight of the diagnostic.\(^{29,37}\) The recorded intensity was the strongest near the ablation front because the emitting plasma had the largest density (which maximized its emission), and the integration distance of the emitting plasma to the detector was the longest. Just inside the ablation front, the recorded intensity dropped by a factor of 2 over a few microns as the plasma became optically thick, absorbing its emission and the emission coming from the back of the target. The time integration and spatial convolution of the diagnostic induced an inward shift, constant on a given image, of the inner gradient up to 4 \(\mu\)m and 20 \(\mu\)m for integration times of 40 ps and 200 ps, respectively.

The angular variation of the projected ablation-front surface (\(\Delta R_\theta\)) was determined from the difference between the angularly resolved contour radius (\(R_\theta\)) and the averaged contour radius \((\bar{R}_\theta)\). To reduce the error, self-emission images were angularly averaged over \(\Delta \theta = 20^\circ\), which was larger than the radial convolution \([360/2\pi] d_{PSF}/(\bar{R}_\theta) < 5^\circ\] and smaller than the scale length of the modes studied here (\(\lambda > 120^\circ\)). An error in \(R_\theta\) of \(\Delta R_\theta = \pm 0.8 \mu\)m was determined on the reference shot by fitting \(R_\theta\) with a normal distribution and taking the number at the 90th percentile. This error was larger than the error in \(\bar{R}_\theta\) of \(\Delta R_\theta/\sqrt{N_{\Delta \theta}} = \pm 0.2 \mu\)m, where \(N_{\Delta \theta} = 360^\circ/\Delta \theta = 18\) is the number of independent measurements. This resulted in \(\delta(\Delta R_\theta) = \delta R_\theta\).
On each image, the location of the projected center of the ablation-front surface on the framing camera was determined by finding the center of the circle that minimizes the standard deviation of its radial difference with the contour. This resulted in an accuracy in the center position of 

\[ \delta R_{\text{center,1}} \approx \delta R_{\theta} / \sqrt{\sum_{m} \bar{I}_m} = \pm 0.2 \mu m. \]

The center location was corrected from the electrical-pulse (EP) propagation that introduced a displacement of the contour along the strip by 

\[ -(\Delta R)_{\text{EP}} = V M (\Delta t)_{\text{EP}}. \]

where \( V \) is the implosion velocity, \( (\Delta t)_{\text{EP}} = V_{\text{EP}} (\bar{R}_0) M \), and \( V_{\text{EP}} \) is the electrical-pulse velocity that was characterized off-line within \( \delta (V_{\text{EP}}) / V_{\text{EP}} \pm 10\% \). 

\( \bar{R}_0 \) was determined by fitting the evolution of \( (\bar{R}_0) \) linearly for the images of first strip [within \( \delta (V) / V = \pm 10\% \) and using a third-order polynomial for the other images [within \( \delta (V) / V = \pm 4\% \) (Ref. 26)]. The associated error of 

\[ \delta R_{\text{center,2}} \approx \left( \frac{(\Delta R)_{\text{EP}} / V_{\text{EP}}}{(\delta (V) / V)^2 + (\delta (V_{\text{EP}}) / V_{\text{EP}})^2} \right)^{0.5} \]

grew with time up to \( \pm 0.5 \mu m \). When the images were on the same strip, the error in \( V_{\text{EP}} \) did not affect the distance between images since it was approximately constant.

The shift between each contour center measured on imploding capsule shots and the corresponding contour center measured on the reference shot was used to determine the projected 

\[ \langle R_0 \rangle \approx \langle R_0 \rangle_{\text{MLT}} + \delta R, \]

where \( \langle R_0 \rangle \) is the intensity peak. (c) Angular variations of the projected ablation-front surface \( (\Delta R_0) \) for the images in (a). (d) Projected ablation-front surface motions \( (\Delta R_{\text{center}}) \) as a function of the averaged contour radius \( (\bar{R}_0) \) along \( x \) (orange circles) and \( y \) (blue squares) obtained by comparing the contour centers in (a) with the contour centers in (b).

Figure 151.32: Comparison of the self-emission images recorded on (a) the second imploding capsule shot and (b) the reference shot. The circles correspond to the inner edge contours of the intensity peak. (c) Angular variations of the projected ablation-front surface \( (\Delta R_0) \) for the images in (a). (d) Projected ablation-front surface motions \( (\Delta R_{\text{center}}) \) as a function of the averaged contour radius \( (\bar{R}_0) \) along \( x \) (orange circles) and \( y \) (blue squares) obtained by comparing the contour centers in (a) with the contour centers in (b).
motion of the ablation-front surface ($\Delta R_{\text{center}}$). On the reference shot, each contour center corresponded to the projection of the initial target position on the diagnostic. On a capsule implosion, the contour centers determined on the first strip were aligned with the corresponding contour centers measured on the reference shot. This made it possible to correct for differences in diagnostic pointing and initial target position. Longitudinal and transverse differences were accounted for by comparing the diagnostic relative magnifications and the image locations, respectively. The error in $\Delta R_{\text{center}}$ was given by

$$
\delta(\Delta R_{\text{center}}) \approx \sqrt{2} \left[ \delta\left(R_{\text{center},1}\right)^2 + \delta\left(R_{\text{center},2}\right)^2 \right]^{0.5} = \pm 0.8 \mu m.
$$

The best estimations of the angular variation of the projected ablation-front surface ($\Delta R_\theta$) and motion ($\Delta R_{\text{center}}$) at an average radius of 150 $\mu m$ were obtained by linearly fitting their evolution with $\langle R_\theta \rangle$ ranging from ~300 $\mu m$ to ~100 $\mu m$ [Figs. 151.32(c) and 151.32(d)]. These evolutions were expected to be linear since, over these radii, there was no significant change in the laser intensity, leading to an almost constant pressure applied to the target: $\Delta R/\langle R_\theta \rangle \approx -\Delta P/\langle P_\theta \rangle \approx -\beta \Delta I/\langle I_\theta \rangle$, where $R_0$ is the initial target radius; $\langle P_\theta \rangle \approx \langle I_\theta \rangle^{0.5}$; $\beta \approx 0.5$ (Ref. 15); and $\Delta R$, $\Delta P$, and $\Delta I$ are the angular variations and $\langle R_\theta \rangle$, $\langle P_\theta \rangle$, and $\langle I_\theta \rangle$ are the angularly averaged values of the radius, pressure, and laser intensity, respectively. Errors in $\langle \Delta R_\theta \rangle_{150}$ and $\langle \Delta R_{\text{center}} \rangle_{150}$ of $\delta(\Delta R_\theta)_{150}$ = $\pm 0.4 \mu m$ and $\delta(\Delta R_{\text{center}})_{150}$ = $\pm 0.6 \mu m$ at the 90th percentile of the error distribution were determined by comparing $\Delta R_\theta$ and $\Delta R_{\text{center}}$ with their linear fits.38

The four measured projected contours were oriented perpendicular to the lines of sight of the corresponding framing cameras to determine the 3-D shape of the ablation-front surface [Fig. 151.33(a)]. Because of the 3-D nonuniformities, the center and averaged radii of each contour were slightly different than the center and averaged radius of the 3-D object. To account for this, one contour was used as a reference and the other contours were shifted transversally and magnified to suppress their radial differences with the reference contour at the two crossing points (i.e., where the polar and azimuthal angles are the same).

The 3-D motion of the ablation-front surface was determined by finding the point at the minimum distance between the four lines defined by the lines of sight of the framing cameras shifted by the measured projected motions and by the displacements introduced during the contour alignment process [Fig. 151.33(b)]. The four projected contours provided two measurements each of the three coordinates of the 3-D center, so the five extra measurements reduced the error in the three coordinates.

The amplitudes of modes $\ell = 1, 2, 3$ of the ablation-front surface were obtained by decomposing the four oriented contours shifted by the measured 3-D displacement using spherical harmonics

$$
[R(\theta_c, \phi_c) = \sum_{\ell=0}^{3} \sum_{m=-\ell}^{\ell} \sqrt{4\pi r_m^\ell} Y_m^\ell(\theta_c, \phi_c),
$$

where $R(\theta_c, \phi_c)$ is the radius normalized to the averaged radius in percent ($r_0^\ell = 100\%$) and ($\theta_c, \phi_c$) are the coordinates of the four contours. The errors in the mode amplitudes were evaluated by simulating the previously determined error distributions of $\delta(\Delta R_{\text{center}})_{150}$ and $\delta(\Delta R_\theta)_{150}$ and fitting the errors by the normal distribution. Errors of $\delta(r_m^\ell) = \pm 0.15\%$, $\delta(r_{0}^\ell) = \pm 0.1\%$, and $\delta(r_{-1}^\ell) = \pm 0.1\%$ were obtained at the 90th percentile for modes $\ell = 1, 2, 3$, respectively.
Figures 151.34(a)–151.34(c) show that, for each mode $\ell$, the difference in the mode amplitudes of the ablation-front surface between shots $\Delta r^m_\ell$ varied linearly with the difference in the corresponding normalized mode amplitudes of the laser’s beam-energy balance $\Delta e^m_\ell$ with low-mode coupling coefficients of $C_1 = -0.66 \pm 0.05$, $C_2 = -0.38 \pm 0.04$, and $C_3 = -0.18 \pm 0.04$. The negative values were due to the fact that the more intense the laser, the more accelerated that part of the target. The fact that the factor was the same between different shots shows that the effects that create nonuniformities other than the beam-energy balance (such as target position, beam pointing, beam timing) were reproducible between shots.

Errors in $\Delta r^m_\ell$ of $0.2\%$, $\Delta r^m_2 = 0.5\%$ and $\Delta r^m_3 = 1.0\%$ at the 90th percentile were obtained by comparing the points with their linear fits. These comparisons were also used to determine the errors at the 90th percentile of $C_\ell$.

The decrease of $C_\ell$ with mode number [Fig. 151.34(d)] was caused by the phase plates that reduced the amplitude of the modes on target. The laser mode on target is given by

$$E(\theta,\phi) = \sum_{\ell=-\infty}^{\infty} \sum_{m=-\ell}^{\ell} \sqrt{4\pi} e^0_\ell Y^m_\ell(\theta,\phi),$$

where $e^0_\ell = a_\ell \sum_{h=-\infty}^{60} E_h Y^m_\ell(\theta_h,\phi_h) = a_\ell e^m_\ell$; $a_\ell$ are coefficients that describe the profile of each beam,

$$\tilde{E}_b(\theta,\phi) = E_b \sum_{\ell=-\infty}^{\infty} a_\ell (2\ell + 1) / 4\pi P_\ell(\cos \gamma)$$

normalized to have $e^0_0 = 100\%$, $P_\ell$ is the Legendre polynomials, and $\gamma$ is the angle between $(\theta, \phi)$ and $(\theta_h, \phi_h)$. The SG5 phase plates reduced the values of modes 1, 2, and 3 by factors of 0.79, 0.47, and 0.2, respectively, which result in a constant $C_\ell / a_\ell = 0.85 \pm 0.07$ that relate the laser modes on target to the target modes [Fig. 151.34(d)].

The values of $C_\ell / a_\ell = \Delta R / \langle \theta \rangle \times \langle \theta \rangle / \Delta I$ resulted in

$$\beta = -150 / (150 - R_0) \times C_\ell / a_\ell \approx 0.44 \pm 0.035,$$

which was close to the theoretical value of 0.5. This shows that the smoothing of the laser modes by the lateral heat transport and the amplification by the Rayleigh–Taylor instability were negligible for those modes, as expected.

These linear evolutions allowed us to determine the residual target mode amplitudes $[r^m_{\text{res}}]$ that remain when the laser
beam energies are balanced and the optimum laser-mode amplitudes that compensate them \( \left[ e_{\text{opt}}^m \right] \). Over the three measurements \( \left( \Delta r_{\text{res}}^m \right) \) is obtained by averaging \( \left( r_{\text{res}}^m \right) = r_{\text{opt}}^m - C_i e_{\text{opt}}^m \) with an associated error of \( \delta \left( \Delta r_{\text{res}}^m \right) = \delta \left( r_{\text{opt}}^m \right) / \sqrt{3} = \pm 0.05\% \); \( \left( e_{\text{opt}}^m \right) \) is given by \( e_{\text{opt}}^m = \left( r_{\text{res}}^m \right) / C_i \). Applying these corrected laser modes would lead to a spherical implosion with a maximum radial error

\[
\approx \left[ \sum_{\ell=0}^{3} \sum_{m=-\ell}^{\ell} \left[ \delta \left( \Delta r_{\text{res}}^m \right) \right]^2 \right]^{0.5} = \pm 1\%.
\]

In summary, tomographies of imploding shells were used to determine the laser energy balance that suppresses target modes \( \ell = 1, 2, \) and 3. This is essential in direct-drive implosion experiments including ICF, where 3-D simulations predict significant enhancement in fusion performance.8

ACKNOWLEDGMENT

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REFERENCES


The Effect of Tritium-Induced Damage on Plastic Targets from High-Density DT Permeation

Introduction

Direct-drive inertial fusion experiments on LLE’s OMEGA Laser System\(^1\) and indirect-drive experiments at Lawrence Livermore National Laboratory’s (LLNL’s) National Ignition Facility\(^2\) use glow-discharge polymer (GDP) as the capsule material that contains the cryogenic DT fusion fuel.\(^3\) Knowledge of the outside diameter of the capsule and the fuel layer’s thickness and uniformity are critical so that appropriate laser conditions can be set for the implosion experiment.

Cryogenic targets measured in the cryogenic target Characterization Stations have had greater than expected outside diameters (OD’s) (up to 13 \(\mu\)m) from thermal contraction after cooling. The expected diameters were calculated from the General Atomics’ (GA’s) National Institute of Standards and Technology (NIST)-traceable, room-temperature–measured OD value and the coefficient of thermal expansion of GDP; this contraction was not observed. As a secondary effect, mismeasurement of the OD can influence the reported fuel-layer thickness. To examine this effect, several experiments were performed including (1) an optical system calibration check; (2) a comparison of OD’s measured in the cryogenic system with a NIST-traceable value (864.1±0.5-\(\mu\)m-OD silicon ball measured at GA); (3) a parametric study of how system variables can affect the OD measurement; and (4) a comparison of an opaque sphere versus a transparent sphere.

Experimental Configuration

1. Optical System Description

Cryogenic targets are characterized\(^4\) using the non-telecentric,\(f/5,\) long-working-distance objective shown in Fig. 151.35. The target is illuminated with a pulsed 630-nm-wavelength light-emitting diode (LED) to minimize the effects of target vibration. This wavelength, along with the \(f/5\) optics, gives a diffraction-limited (Rayleigh criterion) resolution of 3.8 \(\mu\)m. A 1000 \(\times\) 1000-sq-pixel, 12-bit charge-coupled device (CCD) is used to record an image of the target. This gives a 1-\(\mu\)m pixel size given the 1-mm object-space field of view; the image is oversampled and there is no loss in resolution resulting from the pixel size.

2. Image Analysis

The optical system is calibrated (\(\mu\)m/pixel) with a “grid target” that consists of an array of opaque 10-\(\mu\)m-diam aluminum dots that are 20 \(\mu\)m apart on center to within 0.1 \(\mu\)m (see Fig. 151.36). The distortion of the image and centration of the optical axis of the imaging system are also measured and corrected, if necessary, using this grid. Periodic confirmation of calibration using the grid target is performed, especially after any changes are made to the optical path of the system, such as replacing windows or adjusting the optical axis.

To find the outside diameter, 360 radii of the target’s image are traced from the center of the capsule. The region where the intensity transition from dark to light is analyzed, the locations where the intensity begins to transition from the local minimum to the local maximum are determined, and the halfway point between them is deemed the perimeter of the target. [See Fig. 151.39 (p. 162) for an example of a radius versus angular position plot.] The target’s radius is then calculated using the...
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1. Image Illumination

We first studied the saturation of the camera surrounding the capsule's image and its effect on the diameter reported by the software analysis. A GA-measured (864.1±0.5-μm-outer-diam) silicon (Si) ball was measured using Characterization Station #3. Examples shown in Fig. 151.37 give the measured outside diameters and LED currents that produced the images.

Using the traveling stage, 30 dots traversed the eyepiece cross hairs in both the x (parallel to stalk) and y (perpendicular to stalk) directions. The measured distances of \( x = 599.8 \, \mu m \) and \( y = 599.2 \, \mu m \) give a dot pitch in the \( x \) direction of 19.99 \( \mu m \) and in the \( y \) direction of 19.97 \( \mu m \), which agrees with the manufacturer's quoted pitch of 20.0±0.1 \( \mu m \).

2. Analysis Software Confirmation

To examine the reliability of the analysis software, synthetic data were generated and analyzed with the program. The analysis software reproduced exactly the quantities used to produce the synthetic data.

### Possible Sources of Error in OD Measurement

Several parameters were varied to determine their effect on the measured outside diameter as summarized in Table 151.II.

<table>
<thead>
<tr>
<th>Parameter studied</th>
<th>Effect</th>
</tr>
</thead>
<tbody>
<tr>
<td>Illumination intensity</td>
<td>Effect if background is saturated</td>
</tr>
<tr>
<td>Illumination geometry (numerical aperture)</td>
<td>No effect</td>
</tr>
<tr>
<td>Focus shift</td>
<td>Effect only if image is conspicuously out of focus</td>
</tr>
<tr>
<td>Position of the capsule along the optical axis of the imaging system (image refocused)</td>
<td>No effect</td>
</tr>
<tr>
<td>Position of the capsule laterally in the field of view</td>
<td>No effect</td>
</tr>
<tr>
<td>Characterization station</td>
<td>No effect</td>
</tr>
<tr>
<td>Moving Cryostat Transfer Cart</td>
<td>No effect</td>
</tr>
<tr>
<td>Opaque versus transparent sphere</td>
<td>No effect</td>
</tr>
</tbody>
</table>

Calibration Verification

1. Optical Calibration Confirmation

A grid target (manufactured by Applied Image\(^5\)) identical to the one currently being used with the Characterization Stations was measured with a compound microscope that had been calibrated using a Nikon stage reticule. It was the same grid target that was used when evaluating the cryo target characterization technique off-line during the technique's development. The accuracy of the microscope was also confirmed by correctly measuring a NIST-traceable, standard 1-mm-diam ball.

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To examine the reliability of the analysis software, synthetic data were generated and analyzed with the program. The analysis software reproduced exactly the quantities used to produce the synthetic data.

### Table 151.II: Effects of parameters studied on the measured OD of the Si ball.

<table>
<thead>
<tr>
<th>Parameter studied</th>
<th>Effect</th>
</tr>
</thead>
<tbody>
<tr>
<td>Illumination intensity</td>
<td>Effect if background is saturated</td>
</tr>
<tr>
<td>Illumination geometry (numerical aperture)</td>
<td>No effect</td>
</tr>
<tr>
<td>Focus shift</td>
<td>Effect only if image is conspicuously out of focus</td>
</tr>
<tr>
<td>Position of the capsule along the optical axis of the imaging system (image refocused)</td>
<td>No effect</td>
</tr>
<tr>
<td>Position of the capsule laterally in the field of view</td>
<td>No effect</td>
</tr>
<tr>
<td>Characterization station</td>
<td>No effect</td>
</tr>
<tr>
<td>Moving Cryostat Transfer Cart</td>
<td>No effect</td>
</tr>
<tr>
<td>Opaque versus transparent sphere</td>
<td>No effect</td>
</tr>
</tbody>
</table>

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The effects of illumination were then tested with a poly α-methyl styrene (PAMS) capsule with a GA-reported outside diameter of 867.4 μm (wall thickness = 19.2 μm). Results for the x-axis view are shown in Fig. 151.38 along with the LED currents that produced the images. In Fig. 151.38(a), when properly illuminated so that the full dynamic range of the camera is utilized, the OD of the x axis is 869.2 μm and the OD of the y axis is 871.2 μm. Figure 151.38(b) is underilluminated and the OD is slightly overestimated: x-axis OD = 869.4 μm; y-axis OD = 871.6 μm. Figure 151.38(c) is clearly saturated and the OD is significantly underestimated: x-axis OD = 862.2 μm; y-axis OD = 859.8 μm. The program is still reporting a slightly larger OD even with proper illumination. Note that the x-axis OD is closer to the GA value than the y-axis OD with proper illumination. The source of this discrepancy is unclear since the Si ball’s OD measured the same in both axes. It may be an effect of the capsule’s transparency and the illumination nonuniformity present in the frame; this is evident in the offset central bright region inside the capsule’s image.

The error in Fig. 151.38(b) may be caused by noise in the image that is clearly visible in the capsule’s darker periphery; note that the error is small when compared with Fig. 151.38(a). The OD in Fig. 151.38(c) was underestimated because of “blooming.” At saturation, pixels lose their ability to accommodate additional charge. This additional charge will then spread into neighboring pixels, causing them to either report erroneous values or also saturate. This spread of charge to adjacent pixels is known as blooming.

To prevent saturation, the pixel values in the background of the image are displayed by the software in real time by analyzing the image. After it is confirmed that the image is not saturated, data are recorded and analyzed. Only images with the correct illumination are collected for analysis; other than the systematic error of an ~3-μm overestimate of the OD, image illumination is not the source of OD discrepancy.

2. Illumination Geometry

Another test to see if the OD measurement was sensitive to the illuminating ray bundle was performed by adjusting the illuminator’s aperture to control the distribution of the rays coming from the light source. Three aperture conditions were tested: 100%, 50%, and 10% open. The 10%-open condition produced distinct diffraction rings around the image of the Si ball, whereas the others only reduced the intensity of the image. The intensity was adjusted to give the same background intensity for each condition, and images were captured and evaluated. All gave the same OD value as previous experiments: the measurement overestimated the OD by about 3 μm.
3. Focus Shift

Since the imaging systems in the Characterization Stations are not telecentric, the apparent diameter changes with focus adjustment. When examining the target’s surface for debris, the focus is shifted by several hundred micrometers. Returning to the “best focus” after these adjustments is subjective and may be operator dependent. In addition, the target can move in and out of focus because of vibration around its best-focus position.

An experiment was performed to test the effect of moving the Si ball out of the focal plane while holding the objective’s focal plane fixed; these results are summarized in Fig. 151.39. The radial unwrapping of each image, shown below the image, indicates the degree of blurring in its perimeter. The line is the location determined by the analysis software to be the edge of the ball. Images of the silicon ball were obtained with it shifted both toward and away from the objective lens by up to 85 μm [Fig. 151.39(b)] in 17-μm steps. The focus control on the objective was not adjusted to compensate for the shift. The measured OD was reproduced within ±0.5 μm of its average of 867.4 μm for all of the images.

This test was repeated, but this time the focus control on the objective was adjusted to compensate for the shift. The remeasured OD was reproduced within ±0.5 μm of its average of 868.4 μm for all of the images; this time the average was 1 μm larger, most likely because of the lack of telecentricity of the objective lens.

To test the operator’s reproducibility to refocus the objective lens, the Si ball was centered in the layering sphere and the objective’s focus knob was turned to produce a noticeably out-of-focus image. An image was taken, the objective was refocused, and a second image was taken. This was repeated 15 times and the OD of the refocused images had an average of 867.4 μm with a standard deviation of ±0.2 μm. The analysis software is surprisingly robust in that it underreported the OD by up to 6 μm, even for the grossly out-of-focus images [Fig. 151.39(c)]. Figures 151.39(a) and 151.39(b) demonstrate that the OD can be reproduced exactly, even when slightly unfocused.

Images of the PAMS capsule were also obtained with it shifted both toward and away from the objective lens by up to

![Figure 151.39](image-url)

Several images taken during the focus scan: (a) Si ball at focal plane; (b) 85 μm away from focal plane; and (c) deliberately out of focus. The radial unwrapping of each image, shown below the image, indicates the degree of blurring in its perimeter. The line is the location determined by the analysis software to be the edge of the ball. Note that the blurring also overestimates the asymmetry of the capsule’s OD, as indicated by the increased undulation of the line.
45 μm. In this test, the focus control on the objective was not adjusted to compensate for the shift. The measured OD’s were reproduced within ±0.5 μm of their average of 869.2 μm for all of the images.

A lineout along the diameter would allow the operator to more objectively determine best focus in real time. However, the OD measurement is not sensitive to being slightly out of focus, certainly within the operator’s qualitative ability to choose the correct focus.

4. Lateral Shift from Center
Since a cryogenic target is often vibrating both in and out of focus and laterally in the image during data acquisition, the sensitivity of the OD measurement to a lateral shift in the field of view was tested. As the capsule was shifted toward and away from the lens along the x axis in the focus-scan test, images were recorded along the y axis to determine if the measured OD changed with lateral position in the field of view. No difference in OD was measured even at extremes in lateral shift.

5. Characterization Station
The Cryogenic Target Facility contains three identical Characterization Stations. The Si ball was imaged in the same Moving Cryostat Transport Cart (MCTC) in all three stations; each was adjusted to the same illumination and focus conditions. The results are shown in Table 151.III. There was no statistical difference in OD measurement among the three stations.

6. Moving Cryostat Transfer Cart
During cryogenic target experiments, the capsule is stored, layered, transported, and characterized in a cryostat contained in a MCTC. There is some slight variation in window thickness and alignment between the layering spheres in these carts. The data shown in Table 151.III were taken with the Si ball in MCTC #2. The ball was transferred into MCTC #7 and characterized at Characterization Station #3. Using the same illumination and focus conditions, no difference was observed between the measurements made in each cart.

7. Warm Versus Cold Layering Sphere
The Si ball was cooled to 19 K in MCTC #2 at Characterization Station #3 and remeasured. The OD shrunk by, at most, 0.4 μm, as expected from the small thermal expansion coefficient of silicon. There was no statistical difference between the room-temperature and cryogenic measurements.

Data Analysis
1. Contraction of Cold, Unfilled, and D₂-filled Capsules
Images of two cold, unfilled GDP capsules were taken at 90° rotation intervals along both the x- and y-axis views, and the average OD was determined from each. These data, summarized in Table 151.IV, indicated that GDP capsules do contract when cooled; the difference being that they were

<table>
<thead>
<tr>
<th>Characterization Station</th>
<th>x-axis OD (μm)</th>
<th>y-axis OD (μm)</th>
</tr>
</thead>
<tbody>
<tr>
<td>1</td>
<td>867.0</td>
<td>867.4</td>
</tr>
<tr>
<td>2</td>
<td>866.2</td>
<td>866.8</td>
</tr>
<tr>
<td>3</td>
<td>866.6</td>
<td>866.8</td>
</tr>
<tr>
<td>Mean±σ</td>
<td>866.6±0.3</td>
<td>867.0±0.3</td>
</tr>
</tbody>
</table>

<table>
<thead>
<tr>
<th>Capsule Type</th>
<th>Target Number</th>
<th>Outer Diameter (μm)</th>
<th>Average change (μm, corrected)</th>
<th>Percent change (corrected)</th>
</tr>
</thead>
<tbody>
<tr>
<td></td>
<td></td>
<td>Warm x axis</td>
<td>Warm y axis</td>
<td>Cold x axis</td>
</tr>
<tr>
<td>Unfilled GDP</td>
<td>CRYO-ME-4Q13-12</td>
<td>868.8</td>
<td>871.2</td>
<td>859.0</td>
</tr>
<tr>
<td></td>
<td>CRYO-ME-4Q13-8</td>
<td>877.7</td>
<td>876.6</td>
<td>867.2</td>
</tr>
<tr>
<td>D₂-filled GDP</td>
<td>CRYO-2123-19-04</td>
<td>871.6</td>
<td>861.7</td>
<td>12.9</td>
</tr>
<tr>
<td>Mean±σ</td>
<td></td>
<td></td>
<td></td>
<td>12.9±0.3</td>
</tr>
</tbody>
</table>
never exposed to high-pressure DT, unlike the GDP capsules imploded during cryogenic target experiments that showed no contraction.

In addition, a single data point was obtained for a D$_2$-filled GDP capsule that also exhibited contraction. Although, since D$_2$ cryogenic target experiments have not been performed on OMEGA for many years, the fact that it contracted the same as the unfilled GDP capsule indicates that it is not mechanical stress from pressurization that causes the cold, DT-filled capsules’ OD’s to remain close to their room-temperature value.

2. Lack of Contraction of Cold, DT-filled Capsules

Past cryo target data were extracted from the database to compare warm versus cold OD’s as a function of fill date, fuel-layer thickness, Characterization Station number, and MCTC number. These data represent 129 different capsules over a time period from 26 August 2014 to 8 December 2015; they are shown in increasing change in OD in Fig. 151.40. The warm OD’s (measured by GA) and cold OD’s (measured at the Characterization Stations) differ, on average, by 0.06±1.2 μm or 0.01±0.13%. Note that the 3-μm systematic error was not corrected in these data.

The possibility that a step change in measurement accuracy took place at some time in the recent past was explored; these data are shown in Fig. 151.41(a). The data on OD change are also plotted versus Characterization Station number [Fig. 151.41(b)] and MCTC number [Fig. 151.41(c)]. There is no clear trend in any of these data.

![Image](image1)

**Figure 151.40**
LLE’s cold OD subtracted from GA’s warm OD for 129 capsules, with the order shown in increasing difference: (a) the absolute change and (b) the percentage change. On average, they differ by 0.06±1.2 μm or 0.01±0.13%, respectively.

![Image](image2)

**Figure 151.41**
Change in outside diameter versus (a) measurement date, (b) Characterization Station number, and (c) MCTC number. There is no clear trend in any of these data; the data are scattered evenly around zero change. MCTC: Moving Cryostat Transfer Cart.
3. Contraction of a Nonpermeation-Filled Capsule

In an unrelated experiment, a single GDP capsule with a 30-μm-diam hole laser-drilled in its wall was included in a permeation fill along with capsules of similar dimensions. Although not the original purpose of that experiment, the data from it can be used to evaluate if mechanical stresses from pressure gradients across the capsule are responsible for the lack of contraction of GDP capsules at cryogenic temperatures. The hole allows the GDP layer to be exposed to the same beta-decay bombardment inside and outside the capsule's wall in addition to that from tritium in solution within the wall, but without the mechanical stresses of the external pressure that enables permeation. The cold diameter was 3 μm less than a typical capsule that was permeation filled in the same batch, but this difference is not statistically significant compared with the range of changes exhibited in the ensemble of capsules shown in Fig. 151.40. A 3-μm-OD change falls within 2.5× the standard deviation from the nearly zero average OD change; if the sample of capsules was normally distributed, 98.8% of the OD-change values would also lie within 2.5 standard deviations from the mean.

Conclusions

The outside diameters of a silicon ball and two GDP capsules were measured while varying the illumination intensity, illumination geometry, focus shift, position of the capsule along the optical axis of the imaging system, position of the capsule laterally in the field of view, the Characterization Station, and MCTC. The greatest effect on OD measurement was illumination intensity, i.e., saturation of the image around the perimeter of the capsule. In addition, if the peak brightness of the illumination does not coincide with the optical axis and capsule center, intensity variations around the perimeter can locally affect where the analysis software determines the capsule’s edge. Unsaturated images reproduced the OD measurement even under low illumination. Secondly, focus does have an effect on the OD, but errors are produced only if the image is noticeably out of focus.

A systematic overestimation of the OD was revealed during this study; overall, the Characterization Station–measured OD was greater by ~3 μm than that measured at GA. The capsule data acquired during this study corrected for this offset; however, the historic data collected from our database did not correct for this offset since the offset had existed for some unknown time and comparison of historical data must include it. The ~13-μm lack of observed contraction was not a result of measurement error—the systematic error can account for only 3 μm; the remaining effect is real.

The OD’s of three GDP capsules that had not been exposed to DT were measured at both room temperature and 19 K. After the data were corrected for the 3-μm systematic error, they all contracted by 13 μm, which is 1.5% of their warm OD, as expected. A database comparison of 129 DT-filled capsules revealed that they contracted by an average of 0.06±1.2 μm or 0.01±0.13%. A lack of the ~10-μm anticipated contraction and the overmeasurement of the OD by 2 to 3 μm can explain the up-to-13-μm, larger-than-expected OD’s reported by the measurement software.

Radiation damage to the polymer while exposed to beta-particle bombardment during DT permeation explains the lack of contraction. GDP capsules are a highly cross-linked polymer. The average beta-particle energy from tritium decay is 5.7 keV—strong enough to break multiple molecular bonds in the polymer that are a few eV each. Therefore, broken carbon–carbon bonds can readily bond with the ionized hydrogen dissolved in the wall of the capsule. We postulate that the capsules, therefore, swell during permeation to a degree that is nearly compensated for by the contraction during cooling.

Conversely, polystyrene exhibits a high resistance to radiation damage: the polystyrene capsules experience less damage during permeation and contract as expected when cooled, as shown in Table 151.V. Polymers containing aromatic molecules generally are much more resistant to radiation degradation than are aliphatic polymers; this is true whether or not the aromatic group is directly in the chain backbone. Consequently polystyrenes, with a pendant aromatic group, and polyimides, with an aromatic group directly in the polymer backbone, are relatively resistant to high doses of radiation (>4000 kGy) (Refs. 9 and 10).

The GDP capsules containing thicker layers were exposed to DT for a longer period and at a higher concentration during permeation, yet there is no strong correlation of OD change with layer thickness (see Fig. 151.42). There is a possible shift in the median in the data toward less shrinkage as the layer thickness increases, but it is not a convincing trend. This would imply that damage and swelling occur early in the process and conclude quickly.

A DT-gas sample retrieved from the permeation cell following GDP capsule permeation was sent to LLNL to be analyzed with their magnetic-sector mass spectrometer. Many of the constituents in the sample were light hydrocarbons as shown in Fig. 151.43. Since the DT delivery system is constructed of stainless-steel tubing joined by either welded or metal-sealed
Table 151.V: Change in OD after cooling from 293 K to 19 K for a sample of DT-filled polystyrene shells. (The cold OD was reduced by 3 μm from the actual measurement to correct for the systematic error revealed during calibration testing.) The average percent change is 1.11±0.12%, close to the 1.44% thermal contraction calculated from the coefficient of thermal expansion and the temperature change.12

<table>
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<th>Target Number</th>
<th>Outer Diameter (μm)</th>
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<th>Percent Change (corrected)</th>
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</table>

Figure 151.42
Change in outside diameter versus final layer thickness. There is no convincing trend in these data; however, a possible slanting median to the data may indicate less shrinkage with increasing DT exposure.

ACKNOWLEDGMENT
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REFERENCES
The effect of Tritium-induced damage on Plastic Targets from high-density DT permeation


5. APPLIED IMAGE Inc., Rochester, NY 14609.


In direct-drive inertial confinement fusion (ICF), laser beams irradiate a plastic shell containing a thick layer of frozen deuterium–tritium (DT) and ablatively drive an implosion. The ultimate goal of ICF is ignition and energy gain; the minimum shell kinetic energy required for ignition (defined as when the energy from DT fusion reactions exceed the laser energy incident on the target) is given by 

$$E_{\text{min}} \sim \alpha^{1.88} P_{\text{abl}}^{0.77} v_{\text{imp}}^{-5.89}$$

(Ref. 1), where the three parameters of the implosion—$$\alpha$$, $$v_{\text{imp}}$$, and $$P_{\text{abl}}$$ (adiabat (the ratio of the fuel pressure to the Fermi-degenerate pressure at peak implosion velocity), implosion velocity, and ablation pressure, respectively)—are determined primarily by the deposition of the laser energy into the coronal plasma of the target, the subsequent heat conduction to the ablation surface, and the resulting equation of state (EOS) of the shell material. Cross-beam energy transfer (CBET) has been identified in direct-drive experiments on the OMEGA and National Ignition Facility (NIF) lasers to significantly reduce absorption, ablation pressure, and implosion velocity.

The role of CBET in direct drive was identified in early research but only recently identified as the leading cause of decreased energy coupling. When attempts were made to match multiple calculated observables (shell morphology, trajectory, scattered-light spectra and distribution, and shock timing) with experiments, the critical role of CBET became apparent: lowering laser absorption by 20% to 30%. Good agreement with the multiple experimental observables was obtained when both the CBET and nonlocal electron transport models were included in 1-D LILAC and 2-D DRACO simulations. Historically, the role of CBET was masked by using a flux-limited electron transport model that matched laser absorption.

CBET laser–plasma interaction results from two-beam energy exchange via stimulated Brillouin scattering (SBS), which reduces absorbed light and consequently reduces ablation pressure and implosion velocity. The dominant CBET loss mechanism in direct drive occurs when rays counter-propagate (backscatter mode), thereby increasing scattered light, as illustrated in Fig. 152.1(a). For the ignition-relevant overlapped beam intensities of $$\sim 8 \times 10^{14} \text{ W/cm}^2$$ for these NIF experiments, CBET is calculated to reduce laser absorption by 22%, the average implosion speed by $$\sim 9\%$$, and the average ablation pressure by 35% via simulations of the experimental

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**First Observation of Cross-Beam Energy Transfer Mitigation for Direct-Drive Inertial Confinement Fusion Implosions Using Wavelength Detuning at the National Ignition Facility**

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**Figure 152.1**
(a) The effect of cross-beam energy transfer (CBET) in polar direct drive (PDD) predominantly affects the equatorial region; (b) successful CBET mitigation benefits the same region.
conditions. These drive-related results are consistent with other ongoing OMEGA\textsuperscript{7} and NIF-scale\textsuperscript{8} experiments. Reducing the target mass compensates for CBET losses, but the thinner shells become compromised as a result of instability growth.\textsuperscript{12} As shown by the above equation for $E_{\text{min}}$, efficient laser energy coupling and hydrodynamic stability are essential aspects of direct-drive ICF, making CBET mitigation vital. Mitigation strategies of the deleterious CBET effects invoke combinations of spatial, temporal, and wavelength domains. Wavelength detuning, the focus of this article, works by altering the resonance condition between interacting beams.\textsuperscript{2} Wavelength detuning was first examined for indirect drive\textsuperscript{13} and subsequently for direct drive but was prematurely dismissed as a viable option.\textsuperscript{14}

The first direct-drive experiments have been designed for the NIF to study the efficacy of wavelength-detuning CBET mitigation. The target designed for these wavelength-detuning shots on the NIF was adapted from existing 600-kJ designs,\textsuperscript{8} where the trajectories and the shape of the imploding shell and scattered light were well described by the CBET model in DRACO. The basic target design is shown as the inset in Fig. 152.2, where the laser beam power (shown in red) produces a peak overlapped intensity of $\approx 8 \times 10^{14}$ W/cm\textsuperscript{2} at the initial target radius.

The indirect-drive NIF beam geometry distributes 192 beam ports (grouped into 48 quads, shown as projected circles in Fig. 152.3(a)) toward the poles of the NIF target chamber, forming cones of quads with a common polar angle.\textsuperscript{15} Repointing higher-intensity beams from lower latitudes toward the equator partially compensates for the NIF port geometry and higher incident angles when illuminating direct-drive targets. In this modified configuration, referred to as polar direct drive (PDD),\textsuperscript{16,17} CBET predictably dominates in the equatorial region where most of the crossing-beam interactions occur,\textsuperscript{18,19} as shown in Fig. 152.1(b). As a result, PDD implosions tend to become oblate because CBET reduces the laser drive preferentially in the equatorial region. With this motivation, a basic wavelength-detuning strategy exploits the PDD configuration, where each hemisphere has a different wavelength or color. However, the nominal symmetric wavelength mapping [see Fig. 152.3(a)] developed for indirect-drive targets precludes achieving hemispheric wavelength detuning using typical PDD repointing configurations.\textsuperscript{17}

The NIF fiber front end\textsuperscript{15} supports three separate initial colors or wavelength shifts $\Delta \lambda_0 = \{\lambda_1, \lambda_2, \lambda_3\}$ detuned from a central wavelength $\lambda_0 \approx 351$ nm. Currently, the three-color $\{\lambda_1, \lambda_2, \lambda_3\}$ mapping onto the NIF indirect-drive ports is symmetric about the equator [see Fig. 152.3(a)]. To induce a wavelength difference about the equatorial region, a dramatic repointing (referred to as “cone swapping”) is required in either the northern or southern hemisphere [see Fig. 152.3(b) for the southern case]. For cone swapping, in one hemisphere the higher-latitude ports (“inner cones”; $\{\lambda_1, \lambda_2\}$) are repointed to the equator and the lower-latitude ports (“outer cones”; $\{\lambda_3\}$) are repointed to the mid- and high latitudes. For the wavelength-detuning experiments described here, only two different colors were specified such that $\lambda_1 = \lambda_2 \neq \lambda_3$, although future enhanced experiments with three colors are planned. The current NIF
configuration, while not optimal, is capable of achieving a modest wavelength-detuning level $\Delta \lambda_0 = \{+2.3, +2.3, -2.3\}$ Å UV, which is adequate for these proof-of-principle experiments. Cone swapping plus wavelength detuning induces the desired partial hemispheric wavelength difference between beams crossing the equatorial region.

The far-field spot envelope [induced from distributed phase plates (DPPs)\textsuperscript{20} and small-divergence smoothing] quad-mapping is given by the current indirect-drive configuration on the NIF: the inner cones ($\lambda_1, \lambda_2$; red/green projected circles in Fig. 152.3) use a wide elliptical spot shape not well suited for the equatorial region, while the outer cones ($\lambda_3$; blue projected circles in Fig. 152.3) use a narrow elliptical spot shape. The values of the beam energy and repointing were additionally adjusted in the cone-swapping hemisphere to compensate for the swapped spot shapes and the higher incident angles using established PDD design principles.\textsuperscript{17} The cone-swapping repointing scheme and the fixed DPP quad-mapping result in nonoptimal implosion symmetry [see Fig. 152.3(b)]. For this reason, fusion yield and areal density are not metrics for these experiments, which concentrate instead on observables directly related to laser energy absorption: implosion trajectory, shell morphology, and scattered light. Future reconfigurations (optimal DPP’s for PDD,\textsuperscript{17,20} flexible color mapping, and larger wavelength separation) can relieve these constraints, and simulations predict improved overall fusion performance.

In direct drive, many overlapping beams interact with each other in a complicated tangle of intensity, directions, and wavelengths, depending on the beam-port configuration surrounding the imploding target. In addition, each beam strongly refracts and chirps in the expanding, evolving plasma atmosphere during propagation and then scatters energy spectra in a wide spread of exiting paths. The DRACO CBET package (Adaawam)\textsuperscript{21} is an integral part of the 3-D ray-trace package (Mazinis\textsuperscript{22}),\textsuperscript{17} which models each beam as a set of adaptively chosen rays to minimize noise. An extension to the plane-wave CBET model\textsuperscript{2} adapts the steady-state fluid model to 3-D interacting rays in Adaawam by generalizing the wave-vector phase-matching condition. The CBET model\textsuperscript{2} includes relevant SBS physics and results in gain/loss for a probe ray interacting with the total pump angular spectrum. Adaawam calculates the CBET interaction self-consistently in conjunction with the hydrodynamic evolution of the ICF target (via a split-step technique) and captures the necessary coupled interaction of the dynamic electron density profile, temperature, and plasma-flow velocity that dictates the behavior of CBET, and vice versa, since CBET and the hydrodynamics are strongly coupled. Adaawam uses advanced iterative feedback control to stabilize the CBET tightly coupled many-beam interactions while maintaining energy conservation. This model has been compared to many observables across a range of implosions on OMEGA\textsuperscript{23} and the NIF.\textsuperscript{7} An experimentally determined CBET-gain multiplier of 1.5 (from unrelated OMEGA shots\textsuperscript{23}) that use the first-principles EOS tables was applied to all pre- and post-shot simulations without attempting to fit the NIF shots having similar intensity but different scale lengths and pulse shapes. The CBET gain multiplier of 1.5 that applies across laser systems indicates a predictive ability on the initial wavelength-detuning shot campaign at the tested $-8 \times 10^{14}$ W/cm$^2$ intensity range.

Maximal CBET occurs in the rapidly expanding coronal plasma where two interacting rays satisfy the ion-acoustic-wave–matching conditions\textsuperscript{2,13,14} that account for propagation direction, wavelength, and fluid flow; e.g., a CBET resonance occurs at the Mach-1 surface given a radial plasma flow for directly opposed radially propagating rays of equal wavelength. The instantaneous ray wavelength is given by its initial value and the temporal derivative of the electron density (an extension of the common Doppler shift\textsuperscript{24}), which dynamically alters the instantaneous refractive index in space, and thereby the wavelength, and is independent of ray direction. Consequently, the CBET resonance features are altered as the coronal plasma evolves, which directly maps onto a chirped scattered-light measurement that can be employed to help analyze the implosions and laser–plasma interaction physics. A future publication will address the complete set of measurements and modeling.

Wavelength detuning between crossing beams responds differently in indirect- versus direct-drive ICF implosions, depending on the dominant CBET mode. In indirect drive, the sign of small wavelength detuning ($<2$-Å UV) is used to control the direction of energy transfer between interacting beams by leveraging the CBET resonance for the forward-scatter mode.\textsuperscript{13} While this mode occurs in direct drive, it does not increase scattered-light loss because the energy exchanged is spatially shifted and deposited in slightly different regions; however, distortions at small wavelength separations can arise.\textsuperscript{19} In contrast, an outbound ray in the dominant backscatter mode in direct drive experiences CBET gain regardless of the wavelength-difference sign or magnitude (for nominal levels) because the ion-acoustic wave’s contribution dominates the CBET resonance function.\textsuperscript{19} Under atypical conditions, the outbound ray may experience a loss resonance but insignificantly impacts scattered light because the outbound rays typically transport little energy. The ensemble CBET exchange is best described as an interaction volume (a weighted volume that determines
the interaction strength, which depends on path length, intensity, wavelength, electron density, coronal temperature, fluid velocity, etc.) because any high-gain region is equally matched by loss and significant CBET occurs only when the ensemble interaction volume is large. For example, there might be high intensity near a turning point over insignificant path lengths that form an ineffective and small interaction volume with minimal resulting CBET.

The resonant CBET gain region of the outbound rays in the backscatter mode never disappears but rather shifts into a smaller interaction volume because the relative instantaneous wavelength difference changes the ion-acoustic-wave–matching conditions of the interacting rays. The resonance region bifurcates and shifts both farther out in the corona (where the outbound rays have lower intensity and experience higher expanding fluid velocity and lower electron density) and closer inside the corona (where the interaction becomes shielded by the refractive shadow-boundary surface and/or outbound rays that have negligible intensity) [19] [see Fig. 152.1(a)]. A sufficiently large wavelength separation (detuning) significantly reduces CBET exchange for direct drive by decreasing the interaction volume. In contrast, an insufficient wavelength separation can lead to deposition and shell distortion via the forward-scatter mode. [19] The efficacy of wavelength-detuning CBET mitigation diminishes as the plasma expands and the target implodes, which causes the CBET resonance regions to gradually drift into larger interaction volumes during the drive pulse. [19] Larger wavelength-detuning values delay the onset of diminished mitigation. Simulations predict that wavelength-detuning CBET mitigation is effective for both symmetric direct drive (OMEGA) and PDD since the same mechanisms occur in both configurations, although the positive impact is more pronounced for PDD. [19]

With this motivation, for the first time in direct-drive ICF, wavelength-detuning CBET mitigation was demonstrated and shown to improve energy coupling. The NIF PDD wavelength-detuning CBET mitigation campaign shots were performed in three pairs; each pair consisted of one implosion backlit with ~6.7-keV x rays produced from a planar Fe foil target energized by two quads of NIF beams with 45 kJ (see the blue curve in Fig. 152.2) of UV laser energy per beam with an equatorial view of the compressing shell and a second implosion for self-emission images of the compressing target from equatorial and polar views. Additional diagnostics measured both hard x rays produced by energetic electrons arising from the stimulated Raman and possible two-plasmon–decay instabilities. The inferred levels contain at most only a few percent of the incident energy and do not affect the analysis of the laser–target coupling and CBET. [7] The first pair of control shots (N160405 and N160406) with the same wavelength for all the beams (zero detuning) were performed to establish the baseline experimental observables. Next, two pairs of experiments with a detuning mapping of $\Delta \lambda_0 = [+2.3,+2.3,−2.3] \, \text{Å}$ were performed to evaluate the efficacy of wavelength-detuning CBET mitigation. The zero-detuning and first-detuning shots (N160821-001 and N160821-002) employed southern-hemisphere cone swapping, as illustrated in Fig. 152.3(b). The second-detuning shots (N170102 and N170103) employed northern-hemisphere cone swapping, primarily to observe the expected image inversion and to effectively image the self-emission from the antipodal pole. The repointing (accounting for mirror-image cone swapping) and pulse shapes were nominally identical for all shots where the only intended difference was the wavelength configuration.

The simulated and measured backlit gated x-ray radiographs are analyzed to show shell morphology evolution as well as in-flight shell trajectory, which are used to infer energy coupling. The gated images (gate time ~100 ps) shown in Fig. 152.4 compare the shell morphology for the three backlit shots. The experimental framing-camera images are a composite of several images close in time for this slowly moving target that were cross-correlated and adjusted for magnification to enhance the signal-to-noise ratio; the measurements used a 30-µm pinhole. The DRACO simulations post-processed with the x-ray imaging code Spect3D [25] with matching pinholes and gates. The first two rows are radiographs of matched post-shot simulations and experimental results for the baseline zero-detuning and wavelength-detuning shots with northern-hemisphere cone swapping. The last row shows radiographs for detuning shots with northern-hemisphere cone swapping. All the backlit radiograph data show remarkable agreement between simulation and experiment, especially the expected trend for the detuning shots. A mere ~2% to 3% additional laser energy is absorbed with detuning, but since this energy is localized to the equatorial coronal volume fraction (~25%), and the deposition is redistributed to increase hydrodynamic efficiency, the result is dramatic as observed with the gated x-ray radiographs.

Most notable was the design prediction and measurement of the equatorial mass accumulation near the equator with active wavelength detuning (bottom two rows in Fig. 152.4). As predicted, the mass accumulation flipped orientation when cone swapping was applied to the opposite hemisphere. The wavelength-detuning design attempted to minimize the $l = 2$ Legendre mode while accounting for the spot shapes, point-
ing, and energies in conjunction with the expected increased drive in the equatorial region caused by CBET mitigation. The equatorial mass accumulation is a common feature in PDD designs (and not directly related to CBET mitigation), which is caused by lateral mass flow toward the equator (from primarily oblique incidence) when sufficient equatorial drive is available (e.g., from CBET mitigation) and when using non-optimal spot shapes while achieving a small $\ell = 2$.

The shell trajectory is inferred from the simulated and experimental backlit radiographs by first extracting the outer steepest gradient surface or radii [see Figs. 152.4 and 152.5 (inset)]. The majority of the CBET gain occurs in the equatorial region [Fig. 152.1(b)] and consequently the region expected to benefit from wavelength detuning. Both the surface-area-weighted average of the whole extracted surface and a range restricted to the equatorial region (shown here) demonstrate the benefit. When the extracted shell surface is restricted to the equatorial region ($\pm 30^\circ$ region about the equator) and plotted as a function of time (see Fig. 152.5), the inferred implosion speed increases as a result of wavelength-detuning CBET mitigation. The equatorial shell speed increases 9% from 144 to 157 $\mu$m/ns based on simulation (experimentally a 16% increase from 133 to 154 $\mu$m/ns) because wavelength-detuning CBET mitigation deposits 3% additional energy within the small volume over the equator. The enhanced equatorial velocity is consistently observed when comparing the extracted outer shell contours taken from zero detuning and detuning shots in Fig. 152.5 (inset), where the entire surface-area-weighted average implosion speed increases experimentally by 13%.

In conclusion, the first direct-drive wavelength-detuning CBET mitigation experiments on the NIF with a modest wavelength difference between crossing beams confirmed improved

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**Figure 152.4**
Comparison of backlit radiographs from post-shot DRACO simulations and NIF experimental results near the end of the laser pulse at $t = 8.5$ ns. The dashed lines indicate the outer shell surface extracted from each image defined by the steepest gradient in the inward radial direction.
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Future experiments are planned to scope out the capabilities of wavelength-detuning CBET mitigation to further improve coupling and to address the asymmetry by proposing system changes to both OMEGA and the NIF: adding multiple wavelength sources to OMEGA, expanding the NIF’s wavelength-detuning range, using SMA-DPP’s, different wavelengths within NIF’s quads, and remapping the NIF fiber front end to obviate cone swamping.

ACKNOWLEDGMENT

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REFERENCES


Figure 152.5

Equatorial shell trajectories from post-processed simulated (solid lines) and experimental (symbols) backlight radiographs. The red lines/symbols represent the baseline zero-detuning experiment (N160405). The blue lines/symbols represent the average of the two detuning experiments (N160821 and N170102). The inset shows superimposed extracted surfaces from the experimental radiographs of Fig. 152.4, exemplifying the equatorial mitigation.
First Observation of Cross-Beam Energy Transfer Mitigation


21. Inspired by Anishinaabe word Adaawam, meaning “to borrow something from someone,” to describe the cross-beam energy transfer (CBET), just as energy is borrowed from one beam to another. Resource: J. D. Nichols and E. Nyholm, A Concise Dictionary of Minnesota Ojibwe (University of Minnesota, Minneapolis, 1995).

22. Inspired by Anishinaabe word Mazinisin, meaning “be imprinted, have a design,” to describe laser energy deposition just as the laser imprints on and changes target morphology. Resource: J. D. Nichols and E. Nyholm, A Concise Dictionary of Minnesota Ojibwe (University of Minnesota, Minneapolis, 1995).


Mitigation of Cross-Beam Energy Transfer in Ignition-Scale Polar-Direct-Drive Target Designs for the National Ignition Facility

In inertial confinement fusion (ICF), a high-powered laser-driven ablation process is used to implode a spherical shell composed largely of fuel [approximately equimolar deuterium (D) and tritium (T)], producing a central volume (the “hot spot”) of high density and ion temperature. In the “direct-drive” scheme, the ablation is accomplished by direct illumination of the target using a spherical distribution of short-wavelength (λ ≤ 351-nm) UV laser beams. The inertially confined fuel ions rapidly undergo fusion reactions, producing 3.5-MeV alpha particles, some of which are stopped in both the central region and the surrounding dense DT shell. Ignition occurs when these alpha particles deposit enough energy to launch a thermonuclear burn wave, consuming a fraction of the fuel (depending on the imploled fuel’s total areal density) before the high pressure generated by the burn wave causes the target to disassemble. Ignition is predicted to occur if the fuel’s internal energy exceeds a minimum value,\(^1\)

\[
E_{\text{min}} \approx (50.8 \text{ kJ}) \alpha_{\text{in}}^{1.88} \left( \frac{v_{\text{imp}}}{300 \text{ km/s}} \right)^{-5.89} \times \left( \frac{P}{100 \text{ Mbar}} \right)^{-0.77},
\]  

(1)

where \(\alpha_{\text{in}}\) is the ratio of the pressure to the Fermi-degenerate fuel pressure (the “adiabat”) in the dense DT shell, \(v_{\text{imp}}\) is the peak shell implosion speed, and \(P\) is the ablation pressure. The direct-drive approach is of interest because, for the same incident laser energy, it couples ~3 to 5× more energy into the implosion capsule than indirect drive, enabling more fuel mass to be imploled and lowering the threshold on hot-spot energy \(E_{\text{min}}\) as well as the pressure and, therefore, convergence. Note that the threshold energy \(E_{\text{min}}\) in Eq. (1) depends sensitively on the implosion speed, which in turn is directly related to the energy coupled to the ablating shell.

Ignition and total fusion yield are directly connected to the volume of the hot spot, the central region in which the temperatures and densities are sufficient to initiate fusion reactions. This volume is reduced by perturbations on the inner edge of the shell that are seeded by a number of sources, including laser-drive nonuniformities and target imperfections, and grow as a result of the Rayleigh–Taylor instability as the shell is decelerated by the pressure of the interior gas.\(^2\) As this volume is reduced, so is the energy coupled to the hot spot. Deviations from 1-D implosions can also result in incomplete stagnation, producing residual kinetic energy and reduced hot-spot pressure. In its current configuration, the laser beam ports at the National Ignition Facility (NIF)\(^3\) are preferentially distributed toward the poles of the target chamber, designed primarily for use with x-ray–driven targets enclosed in a hohlraum. Direct-drive implosions using this configuration [“polar direct drive” (PDD)\(^4,5\)] require beam repointing to compensate for the lack of equatorial beams and higher incident angles in the equatorial region.

In order to credibly design PDD targets, it is critical to incorporate the important physics in the simulations. Laser direct-drive experiments on OMEGA\(^6\) and the NIF\(^7–9\) have demonstrated that it is necessary to model both cross-beam energy transfer (CBET) and nonlocal electron heat transport. CBET is seeded stimulated Brillouin scattering (SBS) in which two beams interact by means of an intermediate ion-acoustic wave,\(^10\) increasing the scattered light, reducing the ablation pressure, and decreasing energy coupling and shell velocity (especially in the equatorial region for PDD targets). Nonlocal electrons in the corona, by contrast, increase the conversion efficiency of laser energy to shell kinetic energy by means of their larger mean free paths and more-effective transport. These effects have been observed to be important in the modeling of numerous implosion experiments at comparable laser intensities on OMEGA.\(^11,12\) While nonlocal electron transport can increase the hydrodynamic efficiency of the implosion, CBET scatters a sizable fraction (~20% to 30%) of the incident laser energy, reducing both \(P\) and \(v_{\text{imp}}\) and raising \(E_{\text{min}}\). The magnitude of these combined effects is illustrated in Fig. 152.6, where the ablation pressure and shell speed are shown as a function of wavelength-detuning separation, \(\Delta \lambda\). The \(\Delta \lambda = 0\) limit indicates the effects of unmitigated CBET, compared with the much higher drive pressure and shell speed that can be achieved when CBET is mitigated by means of wavelength detuning, as described below.
The ablation pressure and implosion speed are shown for the ignition design as functions of wavelength-detuning separation. Also shown (dashed line) is the approximate in-flight aspect ratio (IFAR) expected for a shell with mass reduced by the amount needed to recover the original 400-km/s implosion speed.

The two PDD designs presented in this article—the alpha-burning design and the lower-adiabat ignition design—address the twin constraints of sufficient shell kinetic energy and implosion uniformity. These are the first ignition-scale direct-drive designs of any dimensionality to include the effects of nonlocal heat transport and CBET. Previous ignition designs did not incorporate this important physics and modeled these processes in an approximate way by using an *ad hoc* flux limiter applied to the classical expression for heat conduction. In these designs, the loss of drive related to CBET is mitigated by detuning the laser-beam wavelengths relative to one another. As with previous PDD designs for the NIF, the drive asymmetries caused by the disposition of the beams are controlled through a combination of independent pulse shapes for different groups of beams, tailored laser beam spot shapes, and beam repointing. The alpha burner has a shell adiabat of almost 5 for greater hydrodynamic stability. In a simulation that models only drive perturbations caused by beam geometry, the alpha burner is predicted to generate bootstrap heating with a yield enhancement of 5× by means of alpha deposition, producing over 10^{17} fusion neutrons. The lower-adiabat ignition design (−3) achieves a gain close to 30.

PDD designs have seen vast improvement since the original concept was first proposed by Skupsky et al. While that earlier design was not capable of ignition, it included components that have been used in most other laser PDD designs: use of different pulse shapes for different laser beams, repointing of beams toward the equatorial region of the target, and equatorial spot shapes that concentrate energy toward the equator. Each of these components compensates for the reduction in equatorial ablative drive because of the greater angles of incidence and resulting energy deposition at lower densities for the laser light driving that region. Marozas et al. presented the first igniting PDD design. Their design improved on the earlier design by using an automated tuning process for the pulse shape designs. Reference 5 also presented a general process for tuning PDD designs and demonstrated the importance of the time dependence of the relative beam-group energies. Their design also made use of spot shapes apertured by a high-order super-Gaussian envelope, reducing the amount of energy flowing around the target [spot-masking apodization (SMA)]. The first design to use a shell of DT ice with a CH ablator rather than a foam/DT shell was that of Collins et al. It also applied the previously introduced beam conditioning through smoothing by spectral dispersion (SSD), employing multiple-frequency modulators applied selectively prior to the main “drive” portion of the laser pulse. The more-recent designs of Lafon et al. use ablator composed of mid-Z elements to reduce perturbation growth resulting from laser imprint by increasing the size of the conduction zones between the laser absorption and ablation regions. PDD designs have also been developed for the shock-ignition scheme, in which a high-intensity laser spike at the end of the drive pulse drives a strong shock, thereby initiating ignition. These designs do not include any mechanism for mitigating CBET, nor do they model nonlocal electron heat transport; the work presented in this article includes the first such designs. Finally, an intermediate-energy PDD detuning design for CBET mitigation has been fielded for the first time on the NIF, demonstrating the effectiveness of this approach and exploring the physical mechanism of wavelength detuning. Observables such as the shapes and trajectories of the in-flight shell inferred through radiographs are well modeled with the CBET model described below. These validated models are used in the designs presented in this work.

The consistent result of each of these investigations is that the low equatorial drive can be successfully compensated for in a number of ways. The design of Collins et al. is the basis of the designs described here, which employs equatorial pulses with 50%-higher power (within the NIF laser performance envelope) than the polar beams, repoints beams toward the equator, and uses SMA to offset the loss of equatorial drive caused by PDD.

For direct-drive targets of sufficient density scale length and laser intensity, the SBS process responsible for CBET is dynamically important. As mentioned above, this process entails the parametric coupling of incident light with an ion-acoustic wave and a backscattered electromagnetic wave. The efficiency...
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of energy transfer is determined by a resonance function of the parameter \( \eta = \left( \omega_{\text{pump}} - \omega_{\text{probe}} \right) / k_a \cdot v / \left( c_a k_a \right) \), where \( \omega_{\text{pump}} \) and \( \omega_{\text{probe}} \) are the ray frequencies of the beams losing and gaining energy, respectively; \( c_a \) is the outflow sound speed; \( k_a = k_{\text{pump}} - k_{\text{probe}} \) is the ion-acoustic wave vector; and \( v \) is the outflow plasma velocity. Energy transfer from the incoming ray to the outgoing ray occurs as \( \eta > 0 \) is satisfied in backscatter mode under normal circumstances.\(^{15}\) CBET is well known in indirect-drive ICF experiments on the NIF, where it has been used to transfer energy between cones of beams to affect low-mode capsule symmetry by means of wavelength detuning,\(^{25}\) but this forward-scatter mode is unimportant for direct drive.\(^{15}\)

In direct-drive ICF, CBET backscatter typically occurs when an outbound ray is refracted into the path of the central, high-energy region of an inbound beam and the ion-acoustic phase-matching conditions are near resonance. The energy transferred between beams by CBET over a distance \( ds \) attenuates the incident laser light by \( 1 - e^{-\Delta x} \), where \( \Delta x = \zeta_{\text{pol}} P_{\text{CBET}} d_s \), the resonance function is given by \( P(\eta) = \eta - v_a / \left( \eta v_a \right)^2 + (1 - \eta^2)^2 \), and \( I_{\text{pump}} \) is the intensity of the “pump” beam. Reduction in energy transfer caused by mismatched inbound–outbound polarization is represented by the factor \( \zeta_{\text{pol}} \).

CBET is particularly effective at scattering energy from incoming rays and significantly reducing ablation pressure because the matching condition can be met over a large volume, where \( |\eta| \approx 1 \) and the resonance function \( P(\eta) \) peaks. Consider this region in the absence of wavelength detuning: the shift in a ray’s frequency resulting from the changing plasma refractive index \( n_r \) is small enough that \( \omega_{\text{pump}} = \omega_{\text{probe}} \) (Ref. 26). The mass flow in the corona is nearly radial, so CBET backscatter is greatest in a region where \( M k_a \cdot r = M \cos \theta_a - 1 \) (where \( M \) is the flow Mach number and \( \theta_a \) is the angle between \( k_a \) and \( r \)). This region is largely exterior to the Mach-1 surface where, as \( M \) increases radially outward, \( \theta_a \) decreases so the product is still approximately unity. As a result of the beam angles in the PDD configuration and the lack of usable equatorial beam ports, this resonance region occurs preferentially over the equator where repointed beams from each hemisphere overlap. The CBET power density during the drive pulse for the ignition design, but without detuning, is plotted in Fig. 152.7(a). Note that the CBET power density includes the transferred power resulting from both backscatter between incoming rays, which has little effect on the target drive, and backscatter between inbound and outbound rays, which is of primary interest here.

The target designs presented here were simulated using the 2-D radiation hydrocode DRACO.\(^{27}\) DRACO uses a 3-D ray-based inverse bremsstrahlung energy deposition model (Mazinisin).\(^{5}\) The CBET model in DRACO, Adaawam,\(^{15}\) uses an angular-spectrum representation (ASR) in which the ASR captures, for each computational zone, the accumulated intensity as a function of direction and color from all the beams that enter that cell, representing the field of pump rays. A pump ray traversing a cell interacts with the other rays, which cross the zones by means of the ASR, using the formalism above. A proportional-integral-differential (PID) predictor–corrector controller iterates until a self-consistent, energy-conserving solution is found. Since the ray-trace approach used in DRACO does not presently include the polarization state of the light, random polarization is included by setting \( \zeta_{\text{pol}} = (1 / \Delta) \left[ 1 + \left( k_{\text{pump}} \cdot k_{\text{probe}} \right)^2 \right] \) (Ref. 28). This model for CBET has been shown to accurately predict the large-scale morphology of implosions on the NIF when a constant multiplier of 1.5 is applied to \( d\tau_{\text{CBET}} \) (Refs. 9 and 15). Since experiments have yet to probe plasma conditions (density scale lengths, flow speeds, and electron temperature) relevant to NIF PDD ignition, the equations above for the attenuation caused by CBET are used here without an \textit{ad hoc} multiplier. Using a 1.5 multiplier would reduce the shell speed and require a redesign of the target and likely a reduction in shell mass and corresponding increase in the IFAR. Simulations indicate it is also possible to compensate for an increase in CBET multiplier.

Figure 152.7

The CBET (cross-beam energy transfer) power density during the drive pulse for simulation of the ignition design (a) without wavelength detuning for CBET mitigation and (b) with hemispheric wavelength detuning. With hemispherical detuning, the interaction volume is reduced in extent in the polar angle. The CBET power density includes both backscatter and sidescatter, which accounts for the higher level of power density in (b). The shell mass density is also indicated (with a radius ~500 to 600 \( \mu m \)), showing the greater convergence with wavelength detuning (right). The Mach-1 surface is indicated by the solid black circle.
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by increasing drive power, even though peak drive power is limited by optics damage considerations.

**DRACO** uses the implicit Schurtz–Nicolaï–Busquet (iSNB) nonlocal heat-transport model,\(^\text{29}\) based on the Schurtz–Nicolaï–Busquet (SNB) model.\(^\text{30}\) The SNB model computes the nonlocal heat flux using multigroup diffusion by means of a multidimensional convolution integral, which has the effect of delocalizing the Spitzer–Härm heat flux. The iSNB model improves on SNB by solving the diffusion equations implicitly for improved robustness and numerical accuracy. Because of sensitivity in direct-drive ignition to fast-electron preheat, modified mean free paths are used to bring results closer to predictions by more complex but computationally expensive nonlocal models (e.g., Ref. 11). These mean free paths are such that nonlocal electron thermal transport overwhelmingly affects the drive rather than the fuel adiabat. The iSNB model has demonstrated predictive capability for shock timing\(^\text{29}\) and shell shape\(^\text{31}\) in numerous experiments on OMEGA.

In the wavelength-detuning approach to CBET mitigation, the laser cavities are detuned slightly for different collections of beams to increase the frequency separation, which in turn alters the region over which the CBET efficiency is greatest. The detuning magnitude considered here for the designs presented here is \(\Delta \lambda \pm 12 \, \text{Å} \) (UV). Extension of these results to \(\pm 6 \, \text{Å}\) is discussed below. The designs presented here will require modifications to the NIF, including the ability to extend the wavelength tunability of the laser drive and enhanced beam conditioning such as multifrequency-modulated SSD and distributed polarization rotators. Future target designs that mitigate laser beam imprint may reduce or alter the requirements for enhanced beam smoothing. A cryogenic handling system that reduces the time between when the target is extracted from the cryostat and the start of the laser pulse will also be necessary. It should also be recognized that PDD would also enable the use of external magnetic fields that may enhance fusion performance by reducing thermal conduction losses from the hot spot and more efficiently trapping the alpha particles.\(^\text{32}\)

If the outgoing probe rays are detuned to shorter wavelengths (blue-shifted) relative to the pump field, the resonance region moves to greater Mach numbers and correspondingly larger radii, where the beam overlap and corresponding energy transfer are reduced.\(^\text{15}\) If the probe rays are red-shifted relative to the pump field, the resonance region moves radially inward, reducing the overlap between the resonance region and the region reached by the rays. Figure 152.7(b) shows the CBET power density for the ignition design but with the simpler hemispheric wavelength-detuning configuration. Over time, for red-detuned outgoing rays, this resonance region is exposed,\(^\text{15}\) reducing the CBET mitigation.

The effectiveness of this approach depends on the choice of which beam groups to “detune” and by how much. Several detuning configurations were investigated (see Fig. 152.8). The

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**Figure 152.8**

Four of the wavelength-detuning configurations explored in this study. The wavelength shift of each port is indicated by the color of the symbols (where green corresponds to a zero wavelength shift).
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This scheme greatly reduces the energy loss caused by the beams interacting across the equator, which is where the greatest scattering occurs, but does not reduce losses caused by interactions between beams from the same hemisphere. The beams on the NIF are divided into four cones for each hemisphere: two inner cones nearer the pole and two outer cones nearer the equator. The banded scheme reverses the sign of the detuning for the two inner cones of beams in each hemisphere, thereby increasing the coupling. The incident laser energy is 1.8 MJ. Both designs achieve a high implosion speed of ~400 μm/ns, sufficient to generate burn-averaged hot-spot pressures of 190 Gbar for the ignition design and 215 Gbar for the alpha burner, which is higher because of the delayed disassembly. It is important to recognize that x-ray–driven implosions on the NIF have achieved inferred hot-spot pressures well in excess of those calculated in these designs. Both designs have moderately low in-flight aspect ratios (IFAR’s), given by the maximum ratio during the implosion of the shell radius to its thickness. The IFAR is an indicator of shell stability, with lower values being less unstable. The ignition design has an IFAR of 23 and a minimum end-of-pulse, density-weighted adiabat of 2.8, and the alpha burner has a somewhat lower IFAR of 21 with a larger ablator adiabat, resulting in an end-of-pulse, density-weighted adiabat of 4.8. Both of these IFAR’s are lower than that of their flux-limited predecessor, which also used a CH ablator and was calculated to withstand the effects of laser imprint. The simulations presented here include only nonuniformities related to port geometry, repointing, CBET, and nonlocal heat transport; sensitivity to other illumination nonuniformities, such as beam power imbalance, and to target imperfections will be investigated in the near future.) The higher fuel adiabat of the alpha burner is reflected in a lower hot-spot convergence fuel ratio of 25, compared to 28 for the ignition design, and a lower peak total fuel areal density of 1.4 g/cm², compared to 1.7 g/cm² for the ignition design. The ion temperature and density of these designs are shown around the time of peak convergence in Fig. 152.10. The alpha burner achieves a total neutron yield of $1.2 \times 10^{17}$ (~320 kJ of fusion energy) and the ignition design achieves a yield of $1.8 \times 10^{19}$, with 1.8 MJ of incident energy, for a gain of 27. While the alpha burner does not ignite, it operates at a moderate adiabat for acceleration-phase stability, and the neutrons generated by bootstrap heating are over 5× that generated by compression alone. Since this design

The two designs shown here use a 194-μm DT shell with a 36-μm CH ablator and an outer radius of 1482 μm. Each design uses a triple-picket pulse shape to shape the adiabat. The tricolor configuration improves on both of these by not detuning the inner cones. The tricolor scheme is more effective than the banded scheme because the inner cones interact with both the equatorial beams in the same hemisphere and with the equatorial beams in the opposite hemisphere; by not detuning the inner cones, more energy is regained from the interaction across the equator than is lost to the intrahemispherical interaction. Figure 152.9 also shows the primary effect of nonlocal heat transport: an increase in absorption efficiency, which occurs especially near the equator where the radial thermal gradient is greater. A comparison between the implosions with and without nonlocal heat transport, shows an increase of ~30% in the absorbed laser energy for the tricolor scheme, resulting in a much higher implosion speed. All three of these schemes introduce a north–south asymmetry, as described above. This asymmetry is greatly reduced by using a fourth configuration, balanced tricolor, in which the tricolor scheme is inverted, north to south, in alternating quadrants.

Figure 152.9
Absorption efficiency is shown as a function of time for four detuning configurations for the polar-direct-drive ignition design, modeled by including the effects of CBET and nonlocal electron thermal transport. The case without detuning is also shown. The legend shows the cumulative absorption efficiency as well as the peak implosion speed. Shown in blue are the pulse shapes used in the inner-cone beams (which preferentially illuminate the polar and mid-latitude regions on the target) and the outer-cone (equatorial) beams for the ignition (blue) and alpha-burning (gray, visible only at ~400 ps) designs.

The tricolor configuration improves on both of these by not detuning the inner cones. The tricolor scheme is more effective than the banded scheme because the inner cones interact with both the equatorial beams in the same hemisphere and with the equatorial beams in the opposite hemisphere; by not detuning the inner cones, more energy is regained from the interaction across the equator than is lost to the intrahemispherical interaction. Figure 152.9 also shows the primary effect of nonlocal heat transport: an increase in absorption efficiency, which occurs especially near the equator where the radial thermal gradient is greater. A comparison between the implosions with and without nonlocal heat transport, shows an increase of ~30% in the absorbed laser energy for the tricolor scheme, resulting in a much higher implosion speed. All three of these schemes introduce a north–south asymmetry, as described above. This asymmetry is greatly reduced by using a fourth configuration, balanced tricolor, in which the tricolor scheme is inverted, north to south, in alternating quadrants.
lacks an “ignition cliff,” it is also less sensitive to drive and target nonuniformities, making this design an ideal platform for initial study; initial estimates suggest the neutron yield for the alpha burner varies approximately linearly with the implosion speed, rather than the much-steeper dependence of an ignition design.

The use of a $\Delta \lambda = \pm 12 \text{ Å} \ (\text{UV})$ detuning bandwidth would require significant modifications to the NIF laser chain as presently understood. It may be possible, however, to obtain ignition-relevant hot-spot conditions for lower values of $\Delta \lambda$. Figure 152.6 shows, for the ignition design, the dependence of ablation pressure and implosion speed on $\Delta \lambda$. (The pointing and detuning configurations are held constant.) The volume over which CBET is active changes as $\Delta \lambda$ is varied; this is the most likely cause of the second-order nonlinearities in the dependence on $\Delta \lambda$. This plot makes clear the effectiveness of detuning as a mitigation scheme; detuning by $\pm 12 \text{ Å}$ increases the drive pressure by over 50% and the implosion speed by $\sim 10\%$. As expected, the CBET efficiency increases as $\Delta \lambda$ decreases, reducing the coupling and raising $E_{\text{min}}$. This reduction in coupling may, in principle, be offset by a reduction in fuel mass at the cost of increased IFAR. Figure 152.6 also shows the IFAR that would result from reducing the fuel mass in order to obtain the original shell speed of 400 km/s. Reducing $\Delta \lambda$ to $\pm 6 \text{ Å}$ corresponds to an increase in the IFAR from 23 to 25 and may require greater beam smoothing to achieve ignition. However, the alpha burner is already on a high shell adiabat and is less sensitive to imprint. Development of a PDD alpha burner with a thinner shell and $\Delta \lambda = \pm 6 \text{ Å}$ is a natural next step.

The two designs presented here—the first of their kind—demonstrate a promising approach to generating high-energy densities on the NIF and offer a useful research platform for ICF ignition. These designs have peak equatorial intensities of $\sim 1.4 \times 10^{15} \text{ W/cm}^2$ and are likely to experience some degree of fast-electron preheat because of two-plasmon decay and stimulated Raman scattering. A solution for this preheat has already been proposed and will be explored, in which the ablator is doped with mid-Z elements, in order to raise the electron temperature and the instability threshold and increase absorption efficiency. This is likely to be far less of a design issue for the alpha burner, which already operates on a high adiabat. This is, in part, because of the lack of the ignition cliff mentioned above. It is also true because an increase of $\Delta \alpha = 1 \ (\text{where } \alpha \text{ is the adiabat})$ is a $\delta \ln \alpha = 50\%$ increase for an $\alpha = 2$ ignition design but only a $\delta \ln \alpha = 20\%$ increase for an $\alpha = 5$ alpha burner, and the fractional increase in hot-spot pressure (in the absence of alpha heating, which is relevant for achieving ignition-scale conditions needed for both target designs) is $\delta \ln p_{hs} \approx -0.9 \delta \ln \alpha$ (Ref. 22), where $p_{hs}$ is the hot-spot pressure. Furthermore, the designs presented here are modeled in 2-D, although nonaxisymmetric perturbations are expected from both the laser-port geometry (which introduces a perturbation with azimuthal mode number $m = 4$) and the detuning configuration ($m = 2$). While these low modes may be compensated by using azimuthal target “shimming,” other detuning configurations that do not introduce nonaxisymmetric modes are also being developed. As mentioned above, an embedded external magnetic field may also improve the target performance.

Figure 152.10
(a) The alpha-burning and (b) ignition designs are shown near peak compression. On the left of each contour plot is the ion temperature and on the right is the mass density.
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REFERENCES


Readout Models for General Electric BAS-MS Image Plates

Introduction
The performance of inertial confinement implosions on OMEGA is potentially limited by the production of hot electrons. The hard x-ray image-plate (HXIP) diagnostic is used to determine the energy distribution of the hot electrons by estimating the spectrum of the x rays that are generated through bremsstrahlung with the target. Simulations and data from this diagnostic agree with a bi-Maxwellian shape for the total time-integrated hot-electron spectrum, which for low-Z targets generates an x-ray energy spectrum that is approximately the sum of two exponentials.

HXIP is a time-integrated multichannel x-ray spectrometer. X rays incident on the detector are attenuated by different high-pass filters for each channel and are then recorded by a General Electric (GE) BAS-MS image plate, whose active layer is BaFBr$_{0.85}$I$_{0.15}$:Eu$^{2+}$ and is structured as described in Table 152.I. After an experiment, the image plates are removed from the detector and the photostimulated luminescence (PSL) is scanned out by a Typhoon FLA 7000 flying spot scanner. Since the initial signal level for some channels is often higher than the saturation limit of the scanner, image plates are repeatedly scanned until all signal levels are below saturation.

The signal on image plates fades both with time and through readout. The time fading of BAS-MS image plates has been characterized by Ohuchi and Hatano, who measured the signal fading given elapsed time and ambient temperature. The decay rate of the signal on an image plate based on readout number alone has been estimated several times, although never for the BAS-MS image plate with the Typhoon FLA 7000 scanner. A model covering the relationship between the depth of deposited energy and the contribution to signal was provided by Bonnet et al., who exponentially weighted the deposited energy distribution with a falloff length. Thoms provides a theoretical model of the readout process over any number of scans based on the photon diffusion equation and finds a triple exponential decay for PSL centers as a function of integrated photon flux.

Image-plate scan sequences for HXIP indicate a small variation in fade ratio (defined as the current signal value divided by the preceding signal value) between channels and a larger variation between different experiments. A data set of 200 scan sequences for HXIP with at least two scans (see, for example, Fig. 152.11) is used to determine the parameters for a new model of the readout fading.

The following sections: (1) explain the nonapplicability of existing image-plate models to the BAS-MS image plate; (2) derive two basic image-plate models; (3) present methods of determining the model parameters from the HXIP experimental data; and (4) discuss the results.

Table 152.I: Layers of a BAS-MS–type image plate. The first four columns are transcribed from vendor documents. The empirical formulas are calculated from the fifth column, using estimates for unknown values such as the composition of “plastic.”

<table>
<thead>
<tr>
<th>Layer</th>
<th>Width ($\mu$m)</th>
<th>Material</th>
<th>Density (g/cm$^3$)</th>
<th>Calculated empirical formula</th>
</tr>
</thead>
<tbody>
<tr>
<td>Surface</td>
<td>9</td>
<td>Polyethylene terephthalate (PET)</td>
<td>1.4</td>
<td>C$<em>{10}$H$</em>{8}$O$_4$</td>
</tr>
<tr>
<td>Phosphor</td>
<td>115</td>
<td>25:1 mix BaFBr$<em>{0.85}$I$</em>{0.15}$:urethane</td>
<td>3.3</td>
<td>Ba$<em>{2263}$F$</em>{2263}$Br$<em>{1923}$I$</em>{339}$C$<em>{741}$H$</em>{1730}$N$<em>{247}$O$</em>{494}$</td>
</tr>
<tr>
<td>Back</td>
<td>12</td>
<td>Plastic</td>
<td>1.4</td>
<td>C$<em>{10}$H$</em>{8}$O$_4$</td>
</tr>
<tr>
<td>Base</td>
<td>190</td>
<td>PET</td>
<td>1.4</td>
<td>C$<em>{10}$H$</em>{8}$O$_4$</td>
</tr>
<tr>
<td>Ferrite</td>
<td>80</td>
<td>MnO, ZnO, Fe$_2$O$_3$, plastic</td>
<td>3.0</td>
<td>Mn$<em>{1015}$Zn$</em>{885}$Fe$<em>{902}$C$</em>{1315}$H$<em>{1315}$O$</em>{4568}$</td>
</tr>
<tr>
<td>Back protective</td>
<td>25</td>
<td>PET</td>
<td>1.4</td>
<td>C$<em>{10}$H$</em>{8}$O$_4$</td>
</tr>
</tbody>
</table>
Image-Plate Models

1. Review of Existing Image-Plate Models

A first model for the readout process for image plates was derived by Thoms\textsuperscript{11} and later extended to include a triple-exponential fading function for the excitation centers in the image plate, which are referred to as photostimulable F-centers.\textsuperscript{12} Photon propagation in the image plate is modeled using a 3-D photon diffusion equation and solved numerically. While this may be appropriate for the Fuji ST III image plate for which parameters were determined,\textsuperscript{12} the light transport equation used requires that (1) photons are scattered isotropically, (2) the absorption length is much larger than the scattering length, and (3) the scattering length is small compared to the sensitive layer thickness so that photons can immediately be treated as diffuse. Since the estimated parameters for a BaFBr$_{0.85}$I$_{0.15}$:Eu image plate that are similar in structure to the BAS-MS indicate significant absorption and weak and anisotropic scattering of photons,\textsuperscript{14} this model is not applicable.

A variation on the Thoms model using the four-flux model\textsuperscript{15} for light transport comes from Masalovich et al., who considered KCl:Eu and KBr:Eu image plates.\textsuperscript{16} The four-flux model invokes fewer assumptions than the photon diffusion equation used by Thoms,\textsuperscript{11} although it requires the F-center distribution to be uniform along all but the depth coordinate. The four-flux model is used to estimate the readout light flux distribution and the PSL escape probability, from which the image-plate response can be calculated. Since readout light and PSL have different wavelengths, the many free parameters for the four-flux model must be specified for both wavelengths. In practice, the two functions can usually be approximated by much simpler linear or exponential functions.

Finally, Vedantham and Karellas\textsuperscript{14} construct a comprehensive image-plate model with more than 20 parameters that relies on simulations of photon transport, taking into account grain sizes and binder material within the image plate. If all the parameters were measured and the structural assumptions of the optical simulations were correct, the model would be appropriate for the BAS-MS image plate. However, many of the model parameters were measured for BaFBr:Eu and assumed to match for BaFBr$_{0.85}$I$_{0.15}$:Eu, and most information on the internals of the BAS-MS image plate has been kept proprietary.

2. Image-Plate Readout Process

Two image-plate models are introduced in the following sections. In the fixed-distribution model, it is assumed that the shape of the distribution of F-centers in the image plate at any time is unchanged and only its magnitude decreases. In the fading-distribution model, the fading of the F-center distribution as a function of time and depth is structured around the interaction of laser photons and F-centers, which generates PSL.

Figure 152.12 summarizes the readout process of a given distribution $\eta(z)$ [1/mm$^3$] of F-centers within the image plate’s sensitive (phosphor) layer. Laser photons at 650 nm pass into the sensitive layer from the front of the image plate, and the photon flux $\phi(z)$ is reduced (by absorption or scattering) deeper in the image plate. A single cross section $\sigma$ [cm$^2$] describes the interaction between laser photons and the photostimulable F-centers to produce a 390-nm photon and remove an F-center.
[Note that F(Br\textsuperscript{−}) and F(\textsuperscript{−}) centers are treated identically.] The density \( \eta(z) \) is assumed to be small enough that absorption by F-centers does not significantly alter the laser photon flux. PSL generated at a given depth \( z \) escapes the image plate and is recorded by the scanner optics with probability \( p(z) \).

Because a detailed and accurate optical analysis of the BAS-MS plates is difficult, the readout laser photon flux can instead be approximated by the function

\[
\phi(z) = I e^{-zL_R},
\]

which is defined over \( z \in [0,d] \) for the \( d \) sensitive layer depth. \( I[1/cm^2 \cdot s] \) is the average photon flux density at the surface of the sensitive layer and \( L_R \) is a length factor accounting for the readout photon scattering and absorption effects within the sensitive layer. Similarly, the escape probability can be approximated by

\[
p(z) = \frac{1}{2} e^{-zL_B},
\]

in which \( L_B \) is another length factor accounting for PSL photon scattering and absorption in the sensitive layer. The F-center distribution \( \eta(z) \) can be estimated by assuming that the density of F-centers is proportional to the energy deposited by incident high-energy particles. Separate proportionality constants may apply for each particle type.\(^{10}\)

3. Fixed-Distribution Model

Under the fixed-distribution model, the total signal \( S_1 \) for the first scan is

\[
S_1 \propto \int_0^d \phi(z) \sigma \eta(z) p(z) dz,
\]

where \( \sigma:[cm^2] \) is the cross section of photostimulated luminescence with an F-center. Because both \( p(z) \) and \( \phi(z) \) contain exponentials, defining a combined absorption and scattering length \( L = (L_{R}^{-1} + L_{B}^{-1})^{-1} \) reduces this expression to

\[
S_1 = \alpha \int_0^d \eta(z) e^{-zL} dz.
\]

The factor \( \alpha:[PSL] \) is a scale factor encompassing the cross section \( \sigma \), average laser photon flux density \( I \), illumination time \( \tau \), the factor 1/2 from \( p(z) \), and the conversion from photostimulated luminescence photons to FLA 7000 intensity units (which are also labeled “PSL”).

To account for readout fading, the fixed-distribution model uses an empirical fading formula that has been used for other scanner and image-plate combinations.\(^{7–9}\) The signal for the \( n \)th scan is then

\[
S_n = S_1 \prod_{i=2}^{n} \left[ f_{\text{max}} - (f_{\text{max}} - f_{\text{min}}) e^{-(i-2)/T_{\text{fade}}} \right].
\]

The fade ratio \( S_n / S_{n-1} \) decays exponentially with time. For the first two scans, the fade ratio is \( f_{\text{min}} \); for large \( n \) it approaches \( f_{\text{max}} \); and in between, the falloff in fade ratio has an exponential time constant \( T_{\text{fade}} \).

The controlling parameters for the fixed-distribution model are the minimum and maximum fade ratios \( f_{\text{min}} \) and \( f_{\text{max}} \), time constant \( T_{\text{fade}} \), falloff length \( L \), and scale factor \( \alpha \).

4. Fading-Distribution Model

In the fading-distribution model, the F-center distribution \( \eta(z) \) changes with readout time according to

\[
\frac{d}{dt} \eta(z) = -\phi(z) \sigma \eta(z).
\]

If \( \tau \) is the total readout time for a given area over a single scan, the change in the number of F-centers during the \( n \)th scan is

\[
\eta(z) \left[ e^{-\sigma \tau \phi(z)} - e^{-\sigma \tau \phi(z)} \right],
\]

where \( \phi(z) \) is the photon flux-density approximation from Eq. (1). Since the change in the F-centers equals the PSL generation, the \( n \)th scan signal is

\[
S_n = \beta \int_0^d \eta(z) w_1(z) e^{-(n-1)\sigma \tau \phi(z)} dz,
\]

where \( \beta:[PSL] \) is a scale factor including the PSL photon to scanner PSL conversion and a factor 1/2 from \( p(z) \), which in Eq. (2) gives the PSL escape probability.

The fading-distribution model expression for the first scan signal reduces to Eq. (4) when the number of readout photons is small \( \phi(z) \sigma \tau < 1 \). The empirical time decay formula in Eq. (5) can be derived as an approximation for small \( n \) if the shape of \( \eta(z) \) is fixed, although the expression for \( T_{\text{fade}} \) is complicated.

The fading-distribution model is parameterized by falloff lengths \( L_R \) and \( L_B \), scale factor \( \beta \), and unitless product \( I \sigma \tau \).
encompassing the average incident photon flux density \( I \) [see Eq. (1)], F-center cross section \( \sigma \), and illumination time \( \tau \).

**Parameter Estimation**

1. Scanner-Induced Image-Plate Errors

The Typhoon FLA 7000 scanner used to scan the image plates uses GE software, 50-\( \mu \)m resolution, a 650-nm laser, and a special calibration procedure as described in Williams et al.\(^{17} \) As a result of a lossy conversion from exponential scaling to the offset quadratic scale used by the scanner software, precision at low intensities is reduced. Time fading is not an issue, however, for the HXIP data set: over the 90 s used to scan a 10-cm \( \times \) 10-cm region, the change in signal caused by time fading is less than 1% for scans begun at least 20 min after exposure.

Streaking effects along one axis of the HXIP images are visible below the image plate boundaries in Fig. 152.11. One possible explanation is that stray readout laser light (possibly reflecting off the front image-plate surface or off scanner components) interacts with regions of the image plate far away from the currently illuminated spot. Only the generated PSL aligned with the light guide of the scanner is recorded, so that the signal at a given location is proportional to the F-center density at that location plus a small fraction of the F-center density along the scan line. If this hypothesis is correct, the streaking effects are linear and can be accounted for.

2. Bounds on Parameter Values

Bounds on the model parameters can be determined using extremal values from the HXIP data set. In this example and for all other parameter estimation, the signal values have been corrected for time fading (assuming a constant 20°C temperature).\(^6 \) The observed maximum fade ratio between successive scans is 0.85 for the change in channel signal from the 11th to the 12th scan of a scan sequence, and the minimum fade ratio is 0.15, measured for the least-filtered channel between the first and second scans.

If the time-fading correction does not introduce any significant errors, then for the image-plate readout models to be consistent with this information, the entire range of fade ratios from 0.15 to 0.85 must be attainable for some F-center distribution. For the fixed-distribution model, this condition requires that the parameters \( f_{\text{min}} < 0.15 \) and \( f_{\text{max}} > 0.85 \). With the fading-distribution model, the minimum-possible fade ratio \( f_{\text{min}} < 0.15 \) is obtained for an F-center distribution concentrated at the front of the sensitive layer \( \eta(z) = \delta(z-d) \) and the maximum \( f_{\text{max}} > 0.85 \) for a concentration at the back of the layer. \( \eta(z) = \delta(z-d) \). Evaluating the fading-distribution model for these two extreme cases yields \( I \tau = -\log f_{\text{min}} > 1.9 \) and \( L_R = d/\log (\log f_{\text{max}}/\log f_{\text{min}}) < 100 \mu \text{m} \). The parameter \( L_B \) cannot be determined without constraining the shape of the depth distribution of the F-centers.

3. Determining Parameters by an Error Minimization Fit

For a given model \( m \), the model parameters \( \Omega_m \) are determined by minimizing the errors between the available data and a fit to that data based on the assumption that the x-ray energy spectra incident on HXIP are sums of two exponentials. The average error \( \chi^2 \) for a specific scan sequence \( s \) is computed using

\[
\chi^2_s = \sum_{n \in \text{scans}} \sum_{c \in \text{scans}} w_c \left( \log \frac{v_{n,c}[\Psi_s;\Omega_m]}{x_{s,n,c}} \right)^2,
\]

where \( w_c \) is a weight inversely proportional to the number of scans for which a channel is neither saturated nor read out below the noise threshold. The values \( x_{s,n,c} \) are the measured signals for scan sequence \( s \), the \( n \)th scan, and the \( c \)th channel. The function \( v_{n,c}[\Psi_s;\Omega_m] \) calculates the expected channel signal as a function of scan sequence parameters \( \Psi_s \) and model parameters \( \Omega_m \) [see the Appendix (p. 190) for details]. The mean square log deviation is chosen as a measure of error because the data are always positive and often a factor of 2 distant from the fit values, so the assumption of normally distributed errors under which minimizing squared (linear) errors yields a maximum-likelihood solution does not apply.

Because the errors \( \chi^2 \) have a long-tailed distribution, the geometric mean is used to combine them into a single cost function \( C \), over which the model parameters \( \Omega_m \) can be minimized. (The arithmetic mean would overemphasize scan sequences whose errors are large for reasons unrelated to the readout model; using it shifts the final model parameter estimates only slightly). Therefore,

\[
C = \sqrt[|s|]{\prod_{s \in \text{seqs}} (\chi^2_s)} \quad \text{(11)}
\]

where \( N_s \) is the number of scan sequences used for the fit.

Due to the large number of total parameters, the fit procedure is performed in two loops: in the outer loop, the minimi-
zation over model parameters $\Omega_m$ uses a global optimization method for $C$ based on a surrogate function using linear interpolation on a Delaunay triangulation\(^\ddagger\) in the inner loop, for each scan sequence, the values $\Psi_j$ minimizing $\chi^2_j$ are found by brute force followed by gradient descent. The simplex mesh surrogate is chosen because it can handle both fine detail (near the minima) and coarse detail (for the global shape) without requiring tuning. Computations on the surrogate are efficient since they reduce to computations on linear functions.

Applying the fit procedure to the fixed-distribution and fading-distribution models over a data set of 212 image-plate scan series with at least two scans provides the model parameters and error measures given in Table 152.II. The error bars and limits on individual parameters are computed from the region of parameter space where the surrogate function is less than 1.1C. An alternative approach to estimating the error for the parameters is to minimize $C$ repeatedly, each time using a variation on the original data set perturbed according to the measurement errors of each value, and use the distribution of parameters found to compute the error for each parameter. For the HXIP data set, such Monte Carlo error estimation yields extremely narrow error bars when only accounting for statistical errors. Accurately modeling the systematic errors involved in the estimation of $v_n,e[W;\Omega_m]$ is beyond the scope of this article.

### Table 152.II: Best-fit parameters for the fixed-and fading-distribution models. See Determining Parameters by an Error Minimization Fit (p. 187) for an explanation of the error bars and bounds.

<table>
<thead>
<tr>
<th>Model</th>
<th>Parameters</th>
<th>Bounds</th>
</tr>
</thead>
<tbody>
<tr>
<td>Fixed distribution</td>
<td>$f_{\text{max}} = 0.83 \pm 0.10$</td>
<td>(0.65, 1.00)</td>
</tr>
<tr>
<td></td>
<td>$f_{\text{min}} = 0.36 \pm 0.02$</td>
<td>(0.35, 0.38)</td>
</tr>
<tr>
<td></td>
<td>$T_{\text{fa}} = 3.75 \pm 1.0$</td>
<td>(0.31, 11.0)</td>
</tr>
<tr>
<td></td>
<td>$L = 136 \pm 43 \mu m$</td>
<td>(39 $\mu m$, $\infty$)</td>
</tr>
<tr>
<td>Fading distribution</td>
<td>$L\sigma_T = 2.8 \pm 0.3$</td>
<td>(2.3, 3.3)</td>
</tr>
<tr>
<td></td>
<td>$L_B = 220 \pm 90 \mu m$</td>
<td>(200 $\mu m$, 680 $\mu m$)</td>
</tr>
<tr>
<td></td>
<td>$L_R = 51 \pm 28 \mu m$</td>
<td>(39 $\mu m$, 136 $\mu m$)</td>
</tr>
</tbody>
</table>

Plotting the minimum $C$ while fixing a single parameter at a time does not capture the complex trade-off between parameters; see Fig. 152.13 for the minimum model errors when fixing two of the parameters. The practical difference between best-fit models is shown in Fig. 152.14. The scale factors $\alpha$ and $\beta$ for the two models cannot be determined since the absolute intensity of the x rays recorded by HXIP is uncertain.

### Discussion

The fading-distribution model yields a significantly lower average error ($C = 0.133$) than the fixed-distribution model ($C = 0.170$). Since the cost function $C$ is a relative measure, neither value implies a quantifiable certainty in the parameter values. However, $C$ can be used as a proxy for the likelihood of a given combination of model parameters, assuming the model in question is correct.

For instance, with the fading-distribution model, $L_R$ is constrained to the range (39 $\mu m$, 136 $\mu m$). While lower values of $L_R$ do not significantly increase $C$, they are unphysical, especially for $L_R$ less than the phosphor particle size of 5 $\mu m$ (Ref. 4). For large $L_R$, the error increases up to the point where the fading-distribution model reduces to a simpler one with a fixed fade ratio. The other length parameter $L_B$ induces only gradual changes in error as it varies; it follows that the minimization of $C$ only weakly constrains $L_B$. To determine $L_B$ with any more precision would require a different method of estimation. The combined parameter $L\sigma_T$, which encompasses the photon flux, PSL cross section, and illumination time, has estimated bounds 2.3 $\leq L\sigma_T \leq 3.3$, outside which $L_B$ and $L_R$ can no longer effectively compensate for the extreme value of $L\sigma_T$.

The interaction between parameters for the fixed-distribution model is more complicated, but it is also more clearly influenced by features of the HXIP data set. For instance, most of the scan sequences available are between two and five scans in length. The variation in fade ratios between channels and between scan sequences is largest for the initial scans and smaller for the final scans in a sequence. As Fig. 152.13 shows, $C$ quickly increases on deviating from the minimized $f_{\text{min}} = 0.36 \pm 0.02$. The error on $f_{\text{min}}$ is small since changes to it directly alter the first fade ratio for about 200 scan sequences. Meanwhile, fade ratios approaching $f_{\text{max}}$ are obtained only for the few sequences with more than five scans, and correspondingly the change in $C$ is more gradual, yielding a larger error on $f_{\text{max}} = 0.83 \pm 0.10$. The bias toward short scan sequences also affects parameter $T_{\text{fa}}$ since for the first three or four scans in a sequence, increasing $T_{\text{fa}}$ in conjunction with $f_{\text{max}}$ only slightly increases $C$. Finally, without the influence of differing fade ratios on the length parameter (as for the fading-distribution model), the value of $C$ is largely independent of $L$, whose value, like that of $L_B$, must be determined by another method.
Figure 152.13
Plot of the cost function $C$ for the (a) fading-distribution model and (b) fixed-distribution model when two parameters are fixed and the rest are minimized. The black line is the contour of the surrogate function at $1.1C$, and $\times$ marks the minimum. The contour is omitted where it is too convoluted.

Figure 152.14
Normalized weight functions $w_n(z)$ for the two models defined so that the equation $S_n = \sigma_n \int w_n(z) \eta(z) dz$ describes the nth scan signal $S_n$ as a function of the F-center distribution $\eta(z)$ in the image plate, where $\sigma_n$ is a scale parameter. For the fixed-distribution model, all normalized weight functions are identical.
Not all inferred model parameters agree with previous work on the BAS-MS image plate. For instance, the fading-distribution model parameter $L_R = 51 \pm 28 \, \mu m \in (39 \, \mu m, 136 \, \mu m)$ implies strong attenuation of the incident readout laser. Although the fading-distribution and fixed-distribution models have different equations, the distorting effect of the depth-dependent fading is small enough for the first scan only that the parameters of the two models can be compared to find $L = (L_R^{-1} + L_B^{-1})^{-1}$, which implies $L \leq L_R$. Both the length $L = 125 \pm 35 \, \mu m$ from Boutoux et al.\textsuperscript{19} and the length $L = 222 \pm 72 \, \mu m$ found by Bonnet et al.\textsuperscript{18} are higher than expected given the value of $L_R$. Moreover, the error plot in Fig. 152.13 indicates that while smaller $L_R$ values are plausible, larger values of $L_R$ significantly increase the average error.

Conclusion

In summary, a simple model for BAS-MS image plates that accounts for readout fading has been proposed and found to improve the Bonnet et al. model\textsuperscript{18} combined with an empirical treatment of readout fading. Model parameters are inferred that minimize errors on a collection of experimental HXIP data. The model implies that the response of the second image-plate scan may not be proportional to the response of the first scan, especially when large variations in the depth distribution of F-centers within the image plate are present. The procedure to determine model parameters accurately determines fading parameters but yields little certainty on values related to the transportation of PSL photons and cannot provide any of the scale factors involved in an absolute calibration.

ACKNOWLEDGMENT

The authors thank D. Edgell and C. Stoeckl for insightful discussions.

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Appendix: Estimating Signal Values for a Specific Scan Sequence

The expected signal for channel $c$ of scan $n$ of sequence $s$ given model parameters $\Omega_m$ and spectrum parameters $\Psi_s$, $v_{n,c}[\Psi_s;\Omega_m]$ are defined as a product of operators,

$$v_{n,c}[\Psi_s;\Omega_m] = j \times M_n[\Omega_m] \times R_{c,s} \times [f(E;\Psi_s)], \quad (A1)$$

where $j$ approximates the additional signal from F-centers in different regions of the image plate caused by stray readout light in the scanner. This correction is performed on estimated signals rather than the image plates themselves in order to avoid negative values. $M_n[\Omega_m]$ provides the signal recorded by the scanner for a given depth distribution of F-centers, using the formulas derived in Image-Plate Models (p. 185). $R_{c,s}$ converts an arbitrary x-ray spectrum to the resulting depth distribution $\rho(z)$ within the region of the image plate corresponding to the channel. It is calculated for each HXIP configuration using a Geant\textsuperscript{20} simulation of the detector. The simulation accounts for backscattering from the aluminum image-plate holder and Compton-scattered photons from the channel filters, which contribute to the F-center distribution in other channels’ regions of interest. Finally, $f(E;\Psi_s)$ is the approximated x-ray energy spectrum given by

$$f(E;\Psi_s) = A_1 e^{-kt_1} A_2 e^{-kt_2}, \quad (A2)$$

where $A_1$, $A_2$, $kt_1$, and $kt_2$ are the spectrum parameters encompassed by $\Psi_s$.

REFERENCES


A Time-To-Frequency Converter for Measuring the Shape of Short Optical Pulses

Introduction
First noted in the 1960s, a mathematical equivalence exists between paraxial beam diffraction and dispersive pulse broadening. This equivalence, known as space–time duality, has led to the development of temporal analogs of several optical devices. An important component of such devices is the time lens, which is designed to impose a time-dependent parabolic phase across an optical pulse passing through it; just as a traditional lens provides a parabolic phase in space. The development of such a time lens has led to applications such as temporal imaging, spectral phase conjugation, and temporal cloaking.

Most modern time lenses produce the required parabolic phase using nonlinear effects such as four-wave mixing (FWM), which requires a highly nonlinear waveguide and careful control of the pump dispersion and timing. Using an electro-optic phase modulator driven by a phase-locked sinusoidal radio-frequency (rf) signal, we must adjust the timing of our test pulse within only one cycle of the rf signal, a task that can be accomplished with commercially available phase shifters. Additionally, the electro-optic phase shift does not have the intensity dependence of FWM and can be used for test pulses of any energy.

One of LLE’s diagnostic needs is to measure the shape of infrared (λ = 1053-nm) pulses with durations in the range of 1 to 30 ps. Pre-shot characterization of such short-pulse beams is important for preventing damage to the system. Optical streak cameras have been used at LLE for this purpose; however, several challenges to streak cameras limit their use. First, time-of-flight broadening occurs because of variations in the kinetic energy of the generated photoelectrons. These variations in kinetic energy lead to different electron velocities and, therefore, different amounts of time to reach the other end of the streak tube. For an infrared-sensitive Ag-O-Cs photocathode (designated S-1 photocathode) used in a streak camera, this leads to impulse responses of several picoseconds in width. Second, space-charge effects cause the electrons generated from short, intense pulses to repel each other. This produces broadening of the electron pulse in the drift region of the streak tube, which causes the measured pulse to be longer. The space-charge effects can be reduced by using lower-power pulses, but lower powers lead to signal-to-noise issues. The combination of these two factors means that streak cameras are not particularly well suited to measuring pulses of durations <10 ps. Finally, and perhaps most importantly, recent experience at LLE has shown that current optical streak tubes based on S-1 photocathodes have such a limited lifetime that the long-term costs of operating such streak cameras are not realistic. Therefore, it would be beneficial to develop new diagnostic techniques as alternatives to the streak cameras.

Temporal imaging systems are of particular interest because they can be run in both single-shot and averaging modes without changing the aperture and resolution of the time lens. They are also well suited to imaging picosecond to tens-of-picosecond pulses. In particular, electro-optic phase modulators driven by GHz-rf signals can have apertures in the tens of picoseconds. As a proof of concept, we have developed a pulse-imaging system that uses an electro-optic phase modulator as a time lens in a time-to-frequency converter configuration. Our device maps the pulse shape onto the spectrum, allowing us to record the pulse shape with an optical spectrum analyzer. In the following sections, we address the design of our system, compare its performance to streak-camera and autocorrelator traces, and discuss how the system can be scaled up to cover a range of 1 to 30 ps.

Theory and System Design
We use a time lens in a time-to-frequency conversion system in which the input pulse first propagates inside an optical fiber before passing through the time lens. For a linear system, the electric field at the output of the dispersive medium of length L can be related to the input electric field in the frequency domain as

\[ E_1(\omega) = E_0(\omega) \exp \left[ i\beta_2 L \left( \omega - \omega_0 \right) \right]. \]  

(1)

where \( E_1(\omega) \) is the Fourier transform of the output electric field and the dispersion effects inside the fiber are included by...
the second derivative, $\beta_2 = \frac{d^2 \beta}{d\omega^2}$, of the modal propagation constant $\beta$ at the central frequency $\omega_0$ of the pulse spectrum. The parameter $\beta_2$ takes into account the group-velocity dispersion (GVD) and affects both the duration of the input pulse and its chirp. The pulse duration and chirp after the fiber are determined by the group-delay dispersion (GDD), given by $D_1 = \beta_2 L$.

When an electro-optic phase modulator driven by a sinusoidal voltage is used as a time lens, the phase shift applied to the pulse has the form

$$\Phi(t) = \phi_0 \cos(2\pi v_m t), \quad (2)$$

where $\phi_0$ is the amplitude of the phase modulation and $v_m$ is frequency of the rf signal used to drive the modulator. The phase amplitude is determined by $\phi_0 = \pi V/V_\pi$, where $V$ is the amplitude of the rf voltage used to drive the modulator and $V_\pi$ is the voltage required for the modulator to produce a phase shift of $\pi$, a known quantity for commercial modulators. The electric field after the phase modulator is then related to the electric field at the output of the dispersive medium as

$$E_2(t) = E_1(t) \exp[-i\Phi(t)]. \quad (3)$$

In close analogy to the focal length of a traditional lens, a focal GDD is used to describe a time lens; it is defined as

$$D_f = \left(\frac{(2\pi v_m)^2 \phi_0}{\phi_0^2} \right)^{-1}. \quad (4)$$

For a time-to-frequency converter, the length of the dispersive medium is chosen such that the GDD of the medium, $D_1 = \beta_2 L$, is equal to the focal GDD of the time lens, $D_f$ (Ref. 7). Therefore, the required length in our case is $L = D_f / \beta_2$. When this condition is satisfied, the output-pulse spectrum maps the temporal shape of the input pulse according to the scaling relation

$$t = D_f (\omega - \omega_0). \quad (5)$$

As with a traditional lens, it is useful to define a time aperture and temporal resolution for our time lens. The time aperture is the longest Gaussian pulse that can be imaged by our system without significant distortion of its measured full width at half maximum (FWHM) and has the form

$$\Delta T = \frac{1}{2\pi v_m}. \quad (6)$$

The time resolution is the shortest FWHM pulse duration that can be generated by compressing a pulse that fits within the time lens aperture and is given by

$$\delta t = \frac{4\ln(2)}{2\pi v_m \phi_0}. \quad (7)$$

A schematic of the experimental setup is shown in Fig. 152.15. A mode-locked laser (High Q femtoTrian IC-1053-400 fs Yb) producing 150-fs pulses at 1053 nm with a 38-MHz
A Time-to-Frequency Converter for Measuring the Shape of Short Optical Pulses

Using $\nu_m = 10$ GHz and a maximum phase amplitude of $\phi_0 = 16$ rad in Eqs. (6) and (7) gives a time aperture of $\Delta T = 15.9$ ps and a resolution of $\delta t = 2.75$ ps. The minimum focal dispersion for the time lens is then found using Eq. (4) to be $D_f = 15.8$ ps$^2$. To calibrate the time lens, the sinusoidal phase modulation is scanned across the pulse using a phase shifter and the amplitude of the rf voltage is adjusted until the peak of the pulse spectrum oscillates over a 1.2-nm range. To create the input GDD, a single-mode fiber (Corning HI1060) was used. Using the value of $\beta_2 = 23.8$ ps$^2$/km at 1053 nm, 667 m of this fiber was required to give the input dispersion of $D_1 = 15.8$ ps$^2$. The chirped pulse was then sent through the phase modulator and the spectrum was recorded using an optical spectrum analyzer (Yokogawa AQ6370D).

Because the laser source produces pulses shorter than the resolution of the time lens, a spectral filter must be applied to the laser signal to broaden the pulse in time. A volume Bragg grating (VBG) with a 0.5-nm bandwidth was used to filter the spectrum. The VBG was used in a double-pass configuration to better attenuate the wings of the spectrum, resulting in a final spectral bandwidth of 0.254 nm.

The second arm of the rf line was filtered to 38 MHz and was used as a clock for a digital delay generator (Stanford Instruments DG645), which triggered an acousto-optic modulator (AOM) and a Rochester Optical Streak System (ROSS). The AOM was used to gate the pulse train to achieve a 0.1-Hz repetition rate to prevent damage to the photocathode of the ROSS and to allow only a single pulse to be captured in the streak camera image, thereby eliminating jitter. The ROSS then captured images of the pulse shape, which were used as a comparison for the time-lens measurements.

Although the time lens has a theoretical aperture of 15.9 ps, this value was found, based on the FWHM of the measured pulse, to be largely the same as the FWHM of the actual pulse. However, even if pulses shorter than the time aperture are used, the wings of the pulse, which extend outside of the time aperture, can still see significant distortions. To explore this effect, numerical simulations of the pulse shape measured by a time-to-frequency converter with the same parameters as our experimental time lens were performed. The input dispersion and phase modulation were modeled using Eqs. (1) and (3), respectively. The frequency axis is scaled to the time axis using the relation given in Eq. (5).

Figure 152.16 shows the results for Gaussian input pulses with a FWHM of (a) 15 ps, (b) 12 ps, and (c) 10 ps. The pulse shape is plotted on a logarithmic scale to better show the behavior in the pulse wings. Comparing the three plots, it is shown that the 12-ps and 10-ps pulses in Figs. 152.16(b) and 152.16(c), respectively, are well imaged in the wings, while the 15-ps pulse has significant errors. Note that the wings are beginning to distort for the 12-ps pulse in Fig. 152.16(b), so the effective aperture is close to 12 ps. Also note that the FWHM of the 15-ps pulse is largely unchanged, with the errors arising from a suppression of the wings. A similar problem occurs for the time resolution, with the simulations showing that the resolution is closer to $\delta t = 3$ ps.

![Figure 152.16](E26582JR)

Numerical simulations showing the pulse shapes at the input (dashed red curve) and output (solid blue curve), respectively, of the time-to-frequency converter. Initial pulse widths are (a) $T_{\text{FWHM}} = 15$ ps, (b) $T_{\text{FWHM}} = 12$ ps, and (c) $T_{\text{FWHM}} = 10$ ps. The time axis for the output pulse was obtained using the scaling from Eq. (5).
Experimental Results

Three experimentally recorded spectra (dashed red curves) are shown in Fig. 152.17. The wavelength axis has been converted to a time axis by first converting wavelength to frequency and then using the focal GDD, $D_f = 15.8 \text{ ps}^2$, as a conversion factor to time. A Gaussian fit (solid blue curves) for each pulse provides a measure of the FWHM duration of the pulse. Our measurements show a typical pulse width of around 7.2 ps, with a few traces showing FWHM pulse durations near 7.32 ps, as in Fig. 152.17(a).

![Figure 152.17](image)

Figure 152.17
Measurement of the pulse shape produced by the volume Bragg grating using the time-to-frequency converter (dashed red curves) for (a) $T_{\text{FWHM}} = 7.32 \text{ ps}$, (b) $T_{\text{FWHM}} = 7.13 \text{ ps}$, and (c) $T_{\text{FWHM}} = 7.24 \text{ ps}$. The time axis is obtained using the scaling from Eq. (5). A Gaussian fit to the data is shown as the solid blue curves.

We first compare these measurements to an autocorrelation trace of the pulse as shown in Fig. 152.18(a). Because the autocorrelation signal was very weak as a result of the low peak intensity of our filtered pulses, the oscilloscope trace was averaged over 512 traces. The autocorrelation was then fitted with a Gaussian profile and found to have a FWHM duration of 10.31 ps. Using the known decorrelation factor of 0.707 for Gaussian pulses, a pulse width of $T_{\text{FWHM}} = 7.29 \text{ ps}$ was obtained, which agreed very well with the FWHM calculated from the time-lens measurements in Fig. 152.17 of 7.20±0.08 ps. To obtain this value, it was assumed that the input pulse shape was approximately Gaussian. While this was a good assumption in this case, the decorrelation factor can change drastically for different pulse shapes, taking a value of 0.65 for sech-shaped pulses. This is the reason why autocorrelation is a useful technique for pulses of unknown shapes.

For our time-lens technique, no assumptions are necessary for the pulse shape; as a result, it can be used for pulses of arbitrary shapes. Finally, the pulses measured by our time lens have a small asymmetric peak located near $t = -9 \text{ ps}$. This peak is likely caused by a secondary reflection in the VBG. It is not present in the autocorrelation trace because autocorrelation involves the overlap of two copies of the same pulse, resulting in a symmetric trace. The ability of the time lens to measure asymmetries in the pulse shape is a major benefit of our technique over an autocorrelation-based technique.

Figure 152.18 compares these results to measurements from the ROSS. Figure 152.18(b) shows the image produced by the streak camera using the fastest sweep of 1.75 ns, where the region of the image corresponding to the signal pulse is zoomed in on. Each image is averaged along the displacement axis to produce a temporal profile for a single pulse. Figure 152.19 shows the resulting temporal profiles from three independent streak-camera traces, using 0.85 ps per pixel for the 1.75-ns sweep. After multiple shots, the pulse width measured on the streak camera was found to be 9.29±0.76 ps, which does not agree with the 7.20±0.08-ps value deduced from the time-to-frequency converter. There are several reasons behind this discrepancy: First, the signal is inherently noisier compared

![Figure 152.18](image)

(a) Autocorrelation of pulses used to generate Fig. 152.17 has a width of $T_{\text{FWHM}} = 10.32 \text{ ps}$; (b) streak-camera image of the same pulses. The signal is labeled in the image.

![Figure 152.19](image)

ROSS measurements of the pulse used to generate Fig. 152.17 (dashed red curves) and a Gaussian fit to the data (solid blue curves) for three different shots: (a) $T_{\text{FWHM}} = 9.01 \text{ ps}$, (b) $T_{\text{FWHM}} = 9.63 \text{ ps}$, and (c) $T_{\text{FWHM}} = 9.38 \text{ ps}$. 
to the other two techniques owing to the single-shot nature of the streak camera. Second, the peak power of the pulses was considerably reduced to prevent the space-charge effects, which also decreased the signal-to-noise ratio. The primary reason the streak camera measurements are consistently over a picosecond longer than those found with either the time lens or the autocorrelator is related to the time-of-flight broadening. For an impulse response $\tau$, the measured pulse width will be $T_m = \sqrt{T_0^2 + \tau^2}$. The measured impulse response of 5.78±0.63 ps accounts for the observed discrepancy between the streak camera and the other two methods of measuring the pulse duration. Furthermore, even though the streak camera can potentially see the asymmetries in the pulse shape, the impulse response of the camera is longer than the asymmetry, which can no longer be resolved. Therefore, the same consistent bump in the pulse shape that could be clearly seen in the time-lens measurements is no longer seen.

**Conclusions**

A time-to-frequency converter using an electro-optic phase modulator acting as a time lens was built. Such a device was used to record the temporal shape of infrared pulses at a wavelength of 1053 nm (width of ~7 ps) and these measurements were compared to those made by using both a streak camera and an autocorrelator. Although the proof-of-concept system has successfully demonstrated the time-to-frequency conversion process, several improvements must be made before its use becomes practical. The most-challenging task is to expand the range of pulse durations that can be successfully imaged to cover the entire 1- to 30-ps range. As seen earlier, the current time lens can only accurately image pulses of less than 12 ps, so the time aperture must be expanded by lowering the drive frequency, $v_m$. To reliably image the wings of 30-ps pulses, the drive frequency must be less than $v_m = 4$ GHz, with a time aperture of $\Delta T = 40$ ps.

By lowering the rf frequency, the time resolution will be expanded by the same factor according to Eq. (7). Therefore, the peak phase modulation must also increase by the same factor to maintain the same time resolution. However, it is also desired to lower the resolution from 3 ps to 1 ps, and the amplitude of the phase modulation must be increased to $\phi_0 = 120$ rad. The first step to accomplishing this will be to use a phase modulator designed for the 1053-nm wavelength. Such modulators can produce phase amplitudes of up to $\phi_0 = 30$ rad. If four of these modulators are connected in series, the required phase amplitude can be achieved. With the new frequency and amplitude for the time lens, the focal dispersion becomes $D_f = 13.2$ ps$^2$, requiring only 554 m of single-mode fiber. This gives a theoretical time resolution of 0.92 ps. However, numerical simulations show that the actual resolution of the time lens is closer to 2 ps. This is likely caused by the aperture of the time lens being smaller than predicted in Eq. (6). Therefore, the resolution will actually be larger than the predicted value according to Eq. (7). Indeed, a 2-ps pulse should be well imaged by this time lens. To image pulses with durations down to 1 ps, it would be necessary to double the total phase amplitude. This would require eight phase modulators connected in series, and the insertion losses would begin to affect the signal-to-noise ratio.

For the OMEGA EP Laser System in particular, this limitation can be worked around because the longer pulses are formed by chirping a shorter pulse with diffraction gratings. Because optical fiber has the opposite GVD as that of the diffraction gratings, the 30-ps pulse will actually recompress during propagation through the input fiber, allowing it to fit within the aperture of the time lens. The drive frequency can therefore be increased to $v_m = 7$ GHz to obtain a time aperture of $\Delta T = 21$ ps, while keeping the phase amplitude at $\phi_0 = 120$ rad. Using these parameters, our simulations show that the full 1- to 30-ps range can be well imaged with only slight errors for the 1-ps pulse. Therefore, we can image much longer pulses with a considerably weaker time lens if the long pulses are properly chirped.

Once the time-to-frequency converter is able to image the proper range of pulses, the system can be converted to a single-shot mode. This can be accomplished by feeding the output of the time lens into a single-shot spectrometer. For the above system with $v_m = 7$ GHz and $\phi_0 = 120$ rad, the spectrometer must be able to resolve spectral widths as small as 160 pm. Using a diffraction grating with a line density of 1200 g/mm and a charge-coupled device with 13.5-$\mu$m pixels, this resolution can be achieved with a spectrometer that is less than 30 cm long.

**ACKNOWLEDGMENT**

This material is based upon work supported by the Department of Energy National Nuclear Security Administration under Award Number DE-NA0001944, the University of Rochester, the New York State Energy Research and Development Authority, and the National Science Foundation under award ECCS-1505636.

**REFERENCES**

The Ninth Omega Laser Facility Users Group Workshop

Introduction
The Ninth Omega Laser Facility Users Group (OLUG) Workshop was held at the Laboratory for Laser Energetics (LLE) on 26–28 April 2017. It was attended by 110 researchers, including scientists, postdoctoral fellows (postdocs), and students (Fig. 152.20). The attendees represented institutions from five countries, including the U.S., UK, France, Spain, and Hungary. As has been the case for previous workshops, postdocs and students received travel support to attend the workshop from the Department of Energy’s (DOE’s) National Nuclear Security Administration (NNSA).

The Workshop Program
The OLUG program included the following four invited talks: “Exploring the Structure of Extra Solar Planets Using the OMEGA Laser,” by Tom Duffy (Princeton University) (Fig. 152.21); “Systematic Fuel Cavity Asymmetries in
Directly Driven ICF Implosions,” by Rahul Shah [Los Alamos National Laboratory (LANL)] (Fig. 152.22); “Fast-Electron Transport in Warm and Hot Dense Plasmas,” by Farhat Beg [University of California, San Diego (UCSD)] (Fig. 152.23); and “Using Multi-Hohlraum Arrays for Studying the Pillars of Creation,” by David Martinez [Lawrence Livermore National Laboratory (LLNL)] (Fig. 152.24). DOE’s NNSA perspective was presented by ICF Program Director Njema Frazier. Other highlights included an evening tutorial, “X-Ray Imaging at OMEGA,” offered by Chuck Sorce (LLE); a facility talk, “Omega Facility Update and Progress on OLUG Recommendations,” by Sam Morse (LLE); research talks by representatives from LLE and the national laboratories [Mike Campbell, LLE; Peter Celliers, LLNL; Kirk Flippo, LANL; and Kyle Peterson, Sandia National Laboratories (SNL)]; a lunch round-table discussion on career opportunities in high-energy-density science, the student and postdoc panel; and a discussion of OLUG’s Findings and Recommendations with LLE management. In addition, LLE staff organized tours of the OMEGA and OMEGA EP lasers (Fig. 152.25). The lunch round-table discussion on career opportunities followed up on the talks from laboratory representatives on Thursday morning and gave the students and postdocs an opportunity to engage in a relaxed and productive discussion with laboratory researchers (Fig. 152.26). The round-table discussion was a new activity of the workshop introduced this year. It was very well received by
A lunch round-table discussion on career opportunities in high-energy density followed up on talks by national laboratory representatives.

Students and postdocs, and it will be continued and expanded in future workshops to include scientists working in universities and industry.

Student, postdoc, scientist, and facility posters comprised a total of 62 poster presentations that were organized in three poster sessions. Of the total number, 46 posters were presented by graduate students and postdocs. In addition, six posters were presented by undergraduate students, and two posters were presented by high school students who had participated in LLE’s 2016 Summer High School Research Program (Figs. 152.27–152.29).

**Student and Postdoc Poster Awards**

In an effort to promote and reward excellence in young researchers, the posters presented at the OLUG Workshop by students and postdocs were reviewed and ranked by a committee of scientists. As a result, honorable mentions and prizes were

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**Figure 152.26**

A lunch round-table discussion on career opportunities in high-energy density followed up on talks by national laboratory representatives.

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**Figure 152.27**

Amina Hussein (University of Michigan) discusses her poster on particle-in-cell simulations of laser-accelerated electrons in underdense plasmas.

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**Figure 152.28**

Rui Hua (UCSD) presented a poster on broadband proton radiography of self-generated fields in strongly shocked, low-density systems.

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**Figure 152.29**

Sapna Ramesh from Pittsford Mendon High School and her mentor Kenneth Marshall (left) and Leah Xiao from Webster Schroeder High School and her mentor Stephen Craxton (right) presented posters based on their summer internships at LLE. In the center is OLUG chair Roberto Mancini.
awarded to those posters at the top of the ranking. The following are the awards granted during this OLUG Workshop (Fig. 152.30).

**Undergraduate Students**
First place ($250): Katelyn Cook and Micah Coates, Houghton College, “Measurement of the $^6$He Decay Produced by the $^9$Be(n,$\alpha$)$^6$He Reaction”

Honorable Mention ($75$): Hannah Harrison, Hannah Visca, David Chin, and Praveen Wakwella, State University of New York (SUNY), Geneseo, “Characterizing Neutron Diagnostics on the nTOF Line at SUNY Geneseo” (Fig. 152.31)

**Graduate Students**
First place ($250$): Hong Sio, Massachusetts Institute of Technology (MIT), “Probing Kinetic and Multi-Ion Fluid Effects in ICF Implosions Using DT and D$^3$He Reaction Time-Histories on OMEGA”

Second place ($175$): Archie Bott, University of Oxford, “Proton Imaging of Stochastic Magnetic Fields”


Honorable mentions ($75$): Derek Nasir, Ohio State University, “Enhanced Laser Plasma Interaction Using Micro-Structured Targets”


Samuel Totorica, Stanford University, “Plasmoid Formation and Particle Acceleration in Laser-Driven Magnetic Reconnection”

**Postdoctoral Fellows**
First place ($250$): Hans Rinderknecht, LLNL, “Measurements of Shock-Front Structure in Multi-Species Plasmas”

Honorable mention ($75$): Edward Marley, LLNL, “Development of a Buried Layer Platform at the OMEGA Laser to Study Non-Equilibrium Coronal Plasmas”

**Nominations and Election**
In the winter of 2017, a nominating committee was established to request nominations for the election of three new executive committee (EC) members according to the guidelines of OLUG’s bylaws. The nominating committee was comprised of Mark Koepke [West Virginia University (WVU), Chair] (Fig. 152.32), Ray Leeper (LANL), and Chris McGuffey (UCSD). The nominations for the three new members of OLUG’s EC were to include one representative from a U.S. university/small business, one representative from a national laboratory/major business, and one representative from non-U.S. researchers. Once again, we had an excellent group of nominees who agreed to put their name on the ballot and were willing to serve on the EC if elected. The election resulted in the selection of Maria Gatu Johnson from MIT, Channing Huntington from LLNL, and Alexis Casner from the University of Bordeaux, France, as new members of the EC. Taking
Mark Koepke (WVU) explained the annual nominations and election process carried out in the winter of 2017 that led to the election of three new members of OLUG’s Executive Committee. The newly elected members are:

- U.S. university/small business: Roberto Mancini [University of Nevada, Reno (UNR), Chair], Mark Koepke (WVU, Vice Chair), Maria Gatu Johnson (MIT), and Johan Frenje (MIT)
- National laboratory/major business: Peter Celliers (LLNL), Channing Huntington (LLNL), and Mingsheng Wei [General Atomics (GA)]
- Junior researcher: Alex Zylstra (LANL)
- Non-U.S. researcher: Alexis Casner (University of Bordeaux, France)
- LLE, ex-officio: Jim Knauer

into account the newly elected members as well as those who continue from the previous year, the EC of OLUG for the year April 2017–April 2018 will comprise the following members:

- U.S. university/small business: Roberto Mancini [University of Nevada, Reno (UNR), Chair], Mark Koepke (WVU, Vice Chair), Maria Gatu Johnson (MIT), and Johan Frenje (MIT)
- National laboratory/major business: Peter Celliers (LLNL), Channing Huntington (LLNL), and Mingsheng Wei [General Atomics (GA)]
- Junior researcher: Alex Zylstra (LANL)
- Non-U.S. researcher: Alexis Casner (University of Bordeaux, France)
- LLE, ex-officio: Jim Knauer

The three new EC members replaced Paul Drake (University of Michigan), Kirk Flippo (LANL), and Peter Norreys (Rutherford Laboratory, UK) who stepped down from the EC after completing their terms. The OLUG EC is very grateful to Paul Drake, Kirk Flippo, and Peter Norreys for their service in the EC and their contributions to the success of OLUG (Fig. 152.33).

### Summary of Findings and Recommendations

An important outcome of OLUG’s annual workshop is the list of Findings and Recommendations that OLUG submits for consideration to LLE’s management every year. The 2017 Findings and Recommendations are summarized below, including those put forward by the student and postdoc panel (Fig. 152.34).

1. An increase in the NLUF shot allocation to advance fundamental high-energy-density science and student/postdoc training.
2. An opposing beam configuration for OMEGA EP.
3. Development of the capability for an absolute measurement of Raman backscattered light.
4. A distributed phase plate (DPP)—smoothed, nanosecond-duration beam with small focal spot on OMEGA EP.
5. Additional heated tritium-fill cells for filling glass capsules.
6. A facility-owned ten-inch manipulator (TIM)—mounted, DMX-type instrument to spectrally characterize x-ray drives on OMEGA EP.
7. More options of spectral range coverage for the imaging x-ray Thomson spectrometer (IXTS) and high-resolution spectrometer 2 (HRS2).

8. An optical Thomson-scattering capability for OMEGA EP.


10. A Dante radiation temperature analysis and time-history result available during shot day.

11. Several upgrades and improvements of the active shock breakout (ASBO) and streak optical pyrometer diagnostics, including (a) suppressing the “wiggles” in OMEGA EP ASBO, (b) a smaller field of view to accommodate smaller targets, (c) an absolute calibration of OMEGA EP ASBO, (d) faster timing combs, and (e) a centralized server to archive and make available all necessary calibration details.

12. A three-wavelength VISAR (velocity interferometer system for any reflector) for OMEGA EP.

13. Improving the beam combiner optic lifetime and/or replacement capability to support year-round interleaved joint OMEGA–OMEGA EP shots.


15. Investigating the extension of the duration of OMEGA EP UV beams to 15 to 20 ns.

16. Investigating the feasibility of splitting one of the OMEGA EP short-pulse beams into two focal spots.

17. Continuing the work to improve the Principal Investigator (PI) portal and web-based resources, in particular, with emphasis on data permission access.

18. A web-based system and better microphones for pre-shot briefings so offsite attendees can improve their involvement and participation in the discussions.

19. The addition of a web-based meeting option to Monday morning’s experiment briefings so PI’s who want to join the meeting after their experiment can do so.

ACKNOWLEDGMENT

This OMEGA Laser Facility Users Group Workshop was made possible in part by the generous support of the National Nuclear Security Administration of the U.S. Department of Energy for travel expenses of students and postdocs; by the Physics Department at the University of Nevada, Reno; and by the Laboratory for Laser Energetics at the University of Rochester for the use and availability of critical resources and support. In addition, OLUG thanks the LLE management for their responsiveness to our Findings and Recommendations. For capturing through his lens the workshop ambiance, OLUG thanks Eugene Kowaluk. Roberto Mancini is the editor for this Proceeding.
During the summer of 2017, 11 students from Rochester-area high schools participated in the Laboratory for Laser Energetics’ Summer High School Research Program. The goal of this program is to excite a group of high school students about careers in the areas of science and technology by exposing them to research in a state-of-the-art environment. Too often, students are exposed to “research” only through classroom laboratories, which have prescribed procedures and predictable results. In LLE’s summer program, the students experience many of the trials, tribulations, and rewards of scientific research. By participating in research in a real environment, the students often become more excited about careers in science and technology. In addition, LLE gains from the contributions of the many highly talented students who are attracted to the program.

The students spent most of their time working on their individual research projects with members of LLE’s technical staff. The projects were related to current research activities at LLE and covered a broad range of areas of interest including laser physics, computational modeling of implosion physics, experimental diagnostic development, laser system diagnostics, physical chemistry, cryogenic target characterization, and web-based data analysis (see Table 152.III).

Table 152.III: High School Students and Projects—Summer 2017.

<table>
<thead>
<tr>
<th>Name</th>
<th>High School</th>
<th>Supervisor</th>
<th>Project Title</th>
</tr>
</thead>
<tbody>
<tr>
<td>Viknesh Baskar</td>
<td>Webster Schroeder</td>
<td>J. P. Knauer and C. J. Forrest</td>
<td>Ion Temperature Analysis of Neutron Time-of-Flight Data</td>
</tr>
<tr>
<td>Nikhil Bose</td>
<td>Pittsford Sutherland</td>
<td>M. J. Guardalben</td>
<td>Compensation for Self-Focusing on OMEGA EP by Use of Frequency Conversion</td>
</tr>
<tr>
<td>Benjamin Chaback</td>
<td>Byron Bergen</td>
<td>J. P. Knauer and C. J. Forrest</td>
<td>Modeling and Analysis of Cherenkov Radiation Detectors</td>
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<tr>
<td>Meshach Cornelius</td>
<td>Gates Chili</td>
<td>T. Walker and G. Brent</td>
<td>Characterization and Detection of the Deterioration of Electrical Connectors in a Flash-Lamp System</td>
</tr>
<tr>
<td>Griffin Cross</td>
<td>Pittsford Sutherland</td>
<td>W. T. Shmayda</td>
<td>Studying the Hydrogen-Palladium System at Low Temperatures</td>
</tr>
<tr>
<td>Matthew Galan</td>
<td>Fairport</td>
<td>R. W. Kidder</td>
<td>Data Services for Scientific Analysis on OMEGA and OMEGA EP</td>
</tr>
<tr>
<td>Claire Guo</td>
<td>Penfield</td>
<td>A. Bose and R. Epstein</td>
<td>Analysis of Asymmetries of the Hot Spot Using Synthetic X-Ray Images</td>
</tr>
<tr>
<td>Jonathan Moore</td>
<td>Pittsford Sutherland</td>
<td>M. D. Wittman and A. Kalb</td>
<td>Predetermination of DT Fuel Mass in Cryogenic Target Capsules from Any Viewing Angle</td>
</tr>
<tr>
<td>Arian Nadimzadah</td>
<td>Brighton</td>
<td>W. T. Shmayda</td>
<td>Modifying Stainless-Steel Surfaces by Electropolishing</td>
</tr>
<tr>
<td>Yujia Yang</td>
<td>Brighton</td>
<td>R. S. Craxton</td>
<td>Improving the Uniformity of Revolver Designs for the National Ignition Facility</td>
</tr>
</tbody>
</table>
The students attended weekly seminars on technical topics associated with LLE’s research. Topics this year included laser physics, fusion, holography, nonlinear optics, atomic force microscopy, laser focusing, and pulsed power. The students also received safety training, learned how to give scientific presentations, and were introduced to LLE’s resources, especially the computational facilities.

The program culminated on 30 August with the “High School Student Summer Research Symposium,” at which the students presented the results of their research to an audience including parents, teachers, and LLE staff. The students’ written reports will be made available on the LLE Website and bound into a permanent record of their work that can be cited in scientific publications.

Three hundred and sixty-four high school students have now participated in the program since it began in 1989. This year’s students were selected from approximately 60 applicants.

At the symposium LLE presented its 21st annual William D. Ryan Inspirational Teacher Award to Mrs. Lois Houlihan, a chemistry teacher at Pittsford Mendon High School. This award is presented to a teacher who motivated one of the participants in LLE’s Summer High School Research Program to study science, mathematics, or technology and includes a $1000 cash prize. Teachers are nominated by alumni of the summer program. Mrs. Houlihan was nominated by Sapna Ramesh, a participant in the 2016 program. Sapna wrote, “Mrs. Houlihan cares deeply about all her students and does everything she can to help them succeed… She goes above and beyond to encourage her students… Her teaching style is also very unique and practical. Whenever we went over a new topic in class, Mrs. Houlihan would start off by asking us about the practical uses of the concept, such as in medicine or industry. She makes a point of learning about each student’s interests and background… Instead of teaching to the test, Mrs. Houlihan wants to spark an interest in science and technology in her students.” Sapna acknowledged her personal debt to Mrs. Houlihan: “In terms of encouraging students, Mrs. Houlihan is the reason that I applied to the summer internship at the Laser Lab… I feel blessed to have had Mrs. Houlihan as a teacher, because she was the first teacher I really connected with. The thing with Mrs. Houlihan is that she has high expectations for each and every student she teaches, but also helps everyone individually to push their limits and reach her expectations.” Sapna concluded by saying, “All in all, Mrs. Houlihan has inspired me and many of her other students to love chemistry and science in general. Not only that, but she has opened my eyes to the practical applications of chemistry in the world. I think the best thing about Mrs. Houlihan is that it is obvious that she enjoys teaching, but her greatest joy is seeing her students succeed in college and beyond.” Ms. Houlihan also received strong support from Mr. Karl Thielking, Principal of Pittsford Mendon High School, who described her as a caring and dedicated teacher with a passion for chemistry.
During FY17, the Omega Laser Facility conducted 1353 target shots on OMEGA and 785 target shots on OMEGA EP for a total of 2138 target shots (see Tables 152.IV and 152.V). OMEGA averaged 10.7 target shots per operating day with Availability and Experimental Effectiveness averages for FY17 of 95.7% and 94.4%, respectively.

OMEGA EP was operated extensively in FY17 for a variety of internal and external users. A total of 773 target shots were taken into the OMEGA EP target chamber and 12 joint target shots were taken into the OMEGA target chamber. OMEGA EP averaged 8.7 target shots per operating day with Availability

Table 152.IV: OMEGA Laser System target shot summary for FY17.

<table>
<thead>
<tr>
<th>Laboratory</th>
<th>Planned Number of Target Shots</th>
<th>Actual Number of Target Shots</th>
<th>ICF</th>
<th>Shots in Support of ICF</th>
<th>Non-ICF</th>
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<tr>
<td>Calibration</td>
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<td>25</td>
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<td>—</td>
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<tr>
<td>Total</td>
<td>1250</td>
<td>1353</td>
<td>93</td>
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<td>858</td>
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</tbody>
</table>

Table 152.V: OMEGA EP Laser System target shot summary for FY17.

<table>
<thead>
<tr>
<th>Laboratory</th>
<th>Planned Number of Target Shots</th>
<th>Actual Number of Target Shots</th>
<th>ICF</th>
<th>Shots in Support of ICF</th>
<th>Non-ICF</th>
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<tr>
<td>LLE</td>
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<td>546</td>
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and Experimental Effectiveness averages for FY17 of 95.8% and 96.6%, respectively.

**Highlights of Achievements in FY17**

1. 100-Gbar Campaign

The OMEGA Laser System is the preeminent direct-drive laser facility. It capitalizes on the benefits of a spherically symmetric laser configuration and high-uniformity focal spots to conduct implosion and high-energy-density–physics experiments. In FY16, LLE embarked on a campaign to seek higher implosion pressures through improved laser power balance and a cryogenic fill-tube target system. In FY17, significant progress was made in these two areas.

LLE has continued the effort to temporally balance the energy over 100-ps sections of the pulse shape. To achieve this, each of the beamlines of OMEGA must balance passive transmission losses and active gain between each of the 60 beams. During earlier characterization of transmission, amplifier losses were directly linked to observed surface scattering. In FY17, 23 amplifier disks were precision cleaned, resulting in a direct increase in transmission and improved balance of passive transmission in the amplifier stages. After a study of the frequency-conversion process was concluded, the second tripler optics were removed since the current three-color-cycle smoothing by spectral dispersion (SSD) does not require a dual-tripler setup. This removal eliminated a source of loss and imbalance in the system. A set of 15 rover calorimeters has been deployed to expedite System Science measurements of the gain and loss of each stage. With dedicated system time, the power imbalance has been cut in half to 3%. Numerous efforts have begun to expand our capability to characterize the system: LLE is working on a full-beam-in-tank (FBIT) diagnostic to observe the focal spot after UV transport; a modification to the streak-camera diagnostic that will improve signal integrity by eliminating fluorescence in the fiber; stage-F digital alignment cameras (to characterize losses caused by pre-shot optics damage); and a passive IR beam-transmission diagnostic that will not require amplified laser shots.

In FY17 significant strides were made on two additional efforts required to reach our goal of 100-Gbar pressure. Achieving the highest uniformity is dependent on the target placement at the time of the implosion. Vibrations have been a major source of target offsets. The moving cryostat transfer carts were outfitted with a vibration isolation stage that uses eddy current feedback to actively damp vibrations. It is difficult to quantify the effect of just these isolators because other changes were also made but the overall improvement is sufficient to achieve ~60% of all targets positioned to less than 10 μm (see Fig. 152.36).

Early theoretical predictions showed that the targets must have nonpermeable capsules to optimize the ablator. A cryogenic fill-tube project is underway to provide for this need on OMEGA and to augment the permeation fill system. In FY17, a DT fill-tube target was successfully layered in the laboratory. By FY20, this system will be able to fill, characterize, and deliver a target to OMEGA.

2. Cross-Beam Energy Transfer Mitigation Study

Longer-term improvements to the laser–plasma interaction physics will require mitigation of cross-beam energy transfer (CBET). This phenomena must be characterized to fully understand how to design a system that minimizes the coupling loss to direct-drive capsules that diverts ~30% of the hydrodynamic drive of implosion capsules. LLE chose to utilize OMEGA EP to produce a tunable UV beam and inject it into the OMEGA target chamber to study CBET. This is an efficient design because it utilizes many of the existing OMEGA EP subsystems with little or no modifications. This effort consists of two main subprojects: (1) The installation and activation of a tunable wavelength source laser injected into the optical parametric amplifier (OPA) system on Beamline 1. The tunable laser was installed in Q3 and is being activated. (2) Transport of the UV beam from the OMEGA EP Bay to port P9 in the OMEGA target chamber. Final design of this project was completed in FY17, and fabrication and installation are underway. Closely tied to this effort, a gas-jet target system has been deployed on OMEGA to produce highly uniform, low-density plasma. This produces the optimum environment to study the interactions of OMEGA beams with the new tunable laser.

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**Figure 152.36**

Cryogenic DT target positioning accuracy for FY15–FY17, including position data for the five most-recent shots. TCC: target chamber center.
3. Improvements to the Laser Systems

The OMEGA EP short-pulse diagnostic package was augmented with a new ultrafast temporal diagnostic able to measure pulse widths between best compression and 100 ps. This feat—achieved with a phase-diversity technique in a fiber pulse stacker\(^2\)—has increased the permissible energy on target for pulses between best compression and 10 ps. This diagnostic is utilized in shot preparations for ensuring proper configuration and also for an on-shot measurement with improved accuracy. Other front-end improvements include the expansion of a time-multiplexed pulse-shaping system to OMEGA EP. Currently deployed on Beamlines 3 and 4, this subsystem allows a single higher-resolution waveform generator to feed the independent OMEGA EP beams and will minimize jitter between pulses. On OMEGA, an improved SSD spectrometer is being used to monitor bandwidth up to the time of shot and to ensure that the phase modulators are operating to specification.

4. Improvements to Target Diagnostics

Diagnostic improvements continue to expand the capabilities of the laser systems. On OMEGA EP, a high-resolution spectrometer has been deployed with the ability to measure time-resolved x-ray spectra over a range that includes the Cu \(K\alpha\) lines. This diagnostic is housed in a \(4\pi\) lead enclosure to optimize the signal-to-noise ratio (SNR) by preventing scattered x-ray signals from reaching the sensor. The \(4\omega\) probe diagnostic has been upgraded with an interferometry arm.

A single-line-of-sight, time-resolved x-ray imager (SLOS-TRXI), developed in conjunction with General Atomics, Lawrence Livermore National Laboratory, Sandia National Laboratories, and Kentech Instruments Ltd., has been deployed on OMEGA to measure hot-spot self-emission with a temporal resolution of 40 ps and spatial resolution of 10 \(\mu\)m. Unique in its ability to take multiple pinhole images along a single axis, this hardware represents the first phase in development with improvements in the areas of throughput and spatial resolution to follow.

LLE continues to design and develop improved diagnostics for characterizing experiments. A neutron time-of-flight diagnostic is being deployed on the H10 port with significant improvement to the SNR. The powder x-ray diffraction diagnostic is being upgraded to acquire time-resolved images.

REFERENCES


Under the facility governance plan implemented in FY08 to formalize the scheduling of the Omega Laser Facility as a National Nuclear Security Administration (NNSA) User Facility, Omega Facility shots are allocated by campaign. The majority (68.1%) of the FY17 target shots were allocated to the Inertial Confinement Fusion (ICF) Campaign conducted by integrated teams from Lawrence Livermore National Laboratory (LLNL), Los Alamos National Laboratory (LANL), Naval Research Laboratory (NRL), Sandia National Laboratories (SNL), and LLE; and to the High-Energy-Density (HED) Campaigns conducted by teams led by scientists from the national laboratories, some with support from LLE.

The Fundamental Science Campaigns accounted for 24.2% of the Omega Laser Facility target shots taken in FY17. Nearly 62% of these shots were dedicated experiments under the National Laser Users’ Facility (NLUF) Program, and the remaining shots were allotted to the Laboratory Basic Science (LBS) Program, comprising peer-reviewed fundamental science experiments conducted by the national laboratories and by LLE.

The Omega Laser Facility was also used for several campaigns by teams from the Commissariat à l’énergie atomique et aux énergies (CEA) of France. These programs are conducted at the facility on the basis of special agreements put in place by Department of Energy (DOE)/NNSA and participating institutions.

In this section, we briefly review all the external user activity at the Omega Laser Facility during FY17.

**FY17 NLUF Program**

During the first quarter of fiscal year 2017 (FY17), the Inertial Fusion Office of DOE/NNSA completed a solicitation, review, and selection process for NLUF experiments to be conducted at the Omega Laser Facility during calendar years (CY’s) 2017 and 2018. Twenty-eight proposals were submitted in response to the call for proposals, and the shot requests totaled 60.5 shot days at the Omega Laser Facility. The proposals were peer reviewed by an independent review committee, and ICF/NNSA selected 13 proposals for funding and shot allocation for CY17–CY18.

CY17 was the first of a two-year period of performance for these 13 NLUF projects (Table 152.VI). In addition, several NLUF campaigns completed experiments during FY17 that had been approved during the FY15–FY16 NLUF cycle. In total, 319 target shots were taken for NLUF projects during FY17. The NLUF experiments conducted at the facility during FY17 are summarized in this section.

A critical part of the NLUF program is the education and training of graduate students in high-energy-density physics. During the year, 33 graduate students from nine universities participated in experiments conducted under the NLUF program at the Omega Laser Facility (Table 152.VII).

**Transport of Relativistic Electrons in Cylindrically Imploded Magnetized Plasmas**

Principal Investigator: F. N. Beg [University of California, San Diego (UCSD)]
Co-investigators: P. Forestier-Colleoni, M. Dozières, and C. McGuffey (UCSD); M. S. Wei and C. M. Krauland [General Atomics (GA)]; P. Gourdain, J. R. Davies, and E. M. Campbell (LLE); S. Fujioka (University of Osaka); and J. J. Santos and D. Batani (University of Bordeaux)

In the fast-ignition (FI) scheme of ICF, fuel compression to high densities and temperatures is achieved in separate processes. A high-energy (≥100-kJ), high-intensity (≥10²⁰-W/cm²) short-pulse (~10-ps) laser is first used to create high-energy (~MeV) electrons (or ions), which then heat the precompressed fuel plasma to initiate ignition. One critical issue is the knowledge of the energy and number of relativistic electrons that can reach, and effectively heat, the core plasma. This unresolved issue warrants a new approach to observe the spatial energy deposition of relativistic electrons.

The objective of the UCSD NLUF project in collaboration with GA, LLE, the University of Bordeaux, and the University
of Osaka is to systematically investigate the propagation and energy deposition of relativistic electrons in a preassembled cylindrical plasma under controlled conditions of density and temperature with and without an external magnetic field. Understanding the role of an external magnetic field in relativistic electron transport and energy deposition is important for several applications including ICF FI, isochoric heating, and the study of warm dense matter.

Our first NLUF experiment in 2017, which measured the time-dependent plasma conditions of an imploded cylinder with and without an external magnetic field, was successfully performed in the OMEGA chamber. This will be useful for the second experiment, which will reveal the temporal evolution of the plasma conditions with fast-electron energy deposition. In the first experiment, 36 beams (0.3 TW/beam, 1.5-ns square pulse) of the OMEGA laser were used to compress a Cl-doped CH foam cylinder to reach densities close to 7 g/cm$^3$. This implosion was characterized by two main diagnostic techniques: proton deflectometry and Cl spectroscopy. A schematic of the experimental layout is shown in Fig. 152.37. The cylinder (600-μm outer diameter and 540-μm inner diameter) was filled with 0.1 g/cm$^3$ of CH foam doped with 1% of Cl. In addition, one Cu foil and one Zn foil were attached to the cylinder’s surface to decrease the magnetic mirror effect and allow K$_\alpha$ emissions during the second shot day. The protons used in proton deflectometry were created by a compressed D$^3$He capsule producing two energy populations: 3.5 MeV and 13 MeV. These protons, collected by CR-39, provided an image of the imploded cylinder deformed by the presence of the magnetic field. The Cl spectroscopy focused on the x-ray absorption of Cl 1s→2p transitions detected on axis. The buck-

<table>
<thead>
<tr>
<th>Principal Investigator</th>
<th>Institution</th>
<th>Title</th>
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<tbody>
<tr>
<td>A. Battacharjee</td>
<td>Princeton University</td>
<td>Dynamics of Magnetic Reconnection in High-Energy-Density Plasmas</td>
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<tr>
<td>F. N. Beg</td>
<td>University of California, San Diego</td>
<td>Transport of Relativistic Electrons in Cylindrically Imploded Magnetized Plasmas</td>
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<tr>
<td>R. P. Drake</td>
<td>University of Michigan</td>
<td>Experimental Astrophysics on the OMEGA Laser</td>
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<tr>
<td>T. S. Duffy</td>
<td>Princeton University</td>
<td>Phase Transitions and Crystal Structure of Tin Dioxide at Multi-Megabar Pressures</td>
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<tr>
<td>R. Jeanloz</td>
<td>University of California, Berkeley</td>
<td>High-Energy-Density Chemical Physics and Planetary Evolution</td>
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<td>H. Ji</td>
<td>Princeton University</td>
<td>Particle Acceleration Resulting from Magnetically Driven Reconnection Using Laser-Powered Capacitor Coils</td>
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<td>D. Q. Lamb</td>
<td>University of Chicago</td>
<td>Properties of Magnetohydrodynamic Turbulence in Laser-Produced Plasmas</td>
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<td>R. Mancini</td>
<td>University of Nevada, Reno</td>
<td>Development of a Photoionized Plasma Experiment on OMEGA EP</td>
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<tr>
<td>R. D. Petrasso</td>
<td>Massachusetts Institute of Technology</td>
<td>Explorations of Inertial Confinement Fusion, High-Energy-Density Physics, and Laboratory Astrophysics</td>
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<td>A. Spitkovsky</td>
<td>Princeton University</td>
<td>Study of Magnetized Collisionless Shocks in Laser-Produced Plasmas</td>
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<td>D. Stutman</td>
<td>Johns Hopkins University</td>
<td>Demonstration of Talbot–Lau X-Ray deflectometry Electron Density Diagnostic in Laser–Target Interactions</td>
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<td>M. S. Wei</td>
<td>General Atomics</td>
<td>Hot-Electron Generation with 10$^{16}$ W/cm$^2$ Infrared Lasers in Shock-Ignition-Relevant Conditions</td>
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Table 152.VII: Graduate students participating in NLUF experiments in FY17.

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<th>Name</th>
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<tr>
<td>Rui Hua</td>
<td>UCSD</td>
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<td>Shu Zhang</td>
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<td>Adrianna Angulo</td>
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<td>Patrick X. Belancourt</td>
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<td>Graeme Sutcliffe</td>
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<td>Cole Holcomb</td>
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Figure 152.37
Schematic of (a) the proton deflectometry setup of the experiment in the OMEGA chamber and (b) the Cl-spectroscopy setup of the experiment in the OMEGA chamber; (c) a cylinder target with cones (made by GA) to prevent plasma interaction with the cylinder's surface. XRFC: x-ray framing camera; MIFEDS: magneto-inertial fusion electrical discharge system; ZVH: zinc von Hamos detector.

lighter for the absorption measurement was a gold foil irradiated by two pairs of stacked (0.3-TW/beam, 1.5-ns) OMEGA beams, creating a 3-ns x-ray source.

Our results show a cylindrical compression of the cylinder by the 36 OMEGA beams. The most important result is the impact of the external magnetic field on the compression. Figure 152.38 shows the temporally and spectrally resolved x-ray signal from the plasma with and without an external magnetic field. The compression time seems to be 0.5 ns later with a magnetic field, and we infer the electron temperature from the spectra of Cl. Another spectrometer [zinc von Hamos (ZVH)] was used to detect the Cu and Zn foil emissions with and without an external magnetic field. Furthermore, the x-ray pinhole charge-coupled device (CCD) shows a more-homogenous x-ray emission for the compression in the presence of an external magnetic field, indicating a more-homogenous compression. For proton deflectometry, the protons interacted with the cylinder at two different delays: 0.3 ns before and 2 ns after the
time of laser incidence on the cylinder (the latter approximately corresponds to the implosion time). Figures 152.39(a) and 152.39(b) show the proton radiography of the cylinder deflected by the external magnetic field (estimated ~5.4 T), 0.3 ns before laser incidence and 2 ns after the laser pulse (after implosion time); the thermoelectric effect appears at the surface of the cone [Figs. 152.39(c) and 152.39(d)]. At this time, the cylinder's diameter can be observed and is directly related to the plasma density, revealing a difference in compression with or without the external magnetic field.

Figure 152.38
Temporally and spectrally resolved absorbed Cl signal (a) with backlighter only, (b) without an external magnetic field, and (c) with a magnetic field.

Figure 152.39
Proton deflectometry results for two delays (~0.3 and +2 ns from the laser) and with or without an external magnetic field: (a) 0.3 ns before the laser without a magnetic field, (b) 0.3 ns before the laser with a magnetic field, (c) 2 ns after the laser with a magnetic field, and (d) 2 ns after the laser without a magnetic field.
The data analysis and simulations are ongoing. Magneto-hydrodynamic (MHD) simulations will be used to retrieve the compression history by matching the diagnostics at different times, and atomic physics codes will be used to determine the time-resolved plasma temperature inside the cylinder. Together, these will determine the best delay for the OMEGA EP beam on the second shot day.

**Experimental Astrophysics on the OMEGA EP Laser**

Principal Investigator: R. P. Drake (University of Michigan)

We study hydrodynamic processes by creating long-duration (tens of nanoseconds) steady shocks using the long-pulse UV laser beams on the OMEGA EP laser. Shocks encountering an interface with a density gradient generate shear or vorticity, depending on the geometry, which will induce either the Kelvin–Helmholtz instability (KH) or the Richtmyer–Meshkov (RM) process. Previously, we have independently studied the KH and RM processes; however, both processes will be induced if the interface is at an oblique angle. The contribution from each process depends on the interface angle. Our initial work aims to minimize the RM growth so that the KH growth dominates over time. We will do this by maximizing the parallel (shear) velocity and minimizing the perpendicular velocity.

The experimental target is shown in Fig. 152.40. Three UV laser beams are incident on an ablator package that includes a plastic layer and a plastic-doped bromine thermal insulator layer. This layer absorbs any laser preheat so it cannot adversely affect the target. The laser beams have a 1.1 full-width-at-half-maximum (FWHM) laser spot and a total energy of ~12 J; also, each of the 10-ns pulses is stacked in time to create an almost 30-ns laser pulse for an overall irradiance of ~4.2 \times 10^{13} \text{ W/cm}^2. This long, steady laser pulse creates a strong shock in the ablator package, is then driven into a 100-g/cm\(^3\) carbon foam, and is finally incident on an oblique interface in a plastic material. The interface has a precision-machined sinusoidal pattern that is either a single mode (\(\lambda = 100 \text{ \mu m}, a = 5 \text{ \mu m}\)) or dual mode (\(\lambda_1 = 100 \text{ \mu m}, a_1 = 2.5 \text{ \mu m}, \lambda_2 = 50 \text{ \mu m}, a_2 = 5 \text{ \mu m}\)). Our goal is to compare the developing structure with a single mode (where no merger is expected) and the dual-mode case. We also have experimental targets with a planar interface.

The evolution of the vertical structures was imaged with a spherical crystal imager (SCI) using Cu K\(_\alpha\) radiation at 8.0 keV. Figure 152.41 shows an example of a high-resolution, high-

![Figure 152.40](image-url)

**Figure 152.40**

Schematic of vortex merger target showing the unstable interface between the low- and high-density material at 40°.

![Figure 152.41](image-url)

**Figure 152.41**

(a) X-ray radiograph of an experimental target using a single-mode initial condition at an oblique interface. (b) The radiograph has been cropped and processed to better visualize the unstable growth at the interface.
Phase Transitions and Crystal Structure of Tin Dioxide at Multimegabar Pressures
Principal Investigator: T. S. Duffy (Princeton University)
Co-investigators: R. F. Smith and F. Coppari (LLNL);
J. K. Wicks (Princeton University); and T. R. Boehly (LLE)
Graduate Students: D. Kim and R. Dutta (Princeton University)

Silica is the most abundant oxide component of terrestrial mantles and serves as an archetype for the dense, highly coordinated silicates of planetary interiors. Understanding the behavior of silica at ultrahigh pressure is necessary to model the structure and dynamics of large rocky exoplanets known as super-Earths. Pressures in the mantles of such exoplanets may exceed 1 TPa, which is beyond the range of standard static high-pressure experimental techniques. Dynamic compression using the OMEGA laser offers an alternative means to explore structures and equations of state of planetary materials at these extreme conditions.

SiO$_2$ is one of a family of dioxides whose high-pressure behavior has been of strong interest because of their extensive polymorphism, highly coordinated structures, and varied transition pathways. The challenge in structure determination of silicates and oxides at exoplanet interior pressures is to obtain sufficient x-ray diffraction intensity from weakly diffracting, low-symmetry phases at ultrahigh pressures. The use of analog materials that undergo similar phase transition sequences at lower pressures has a long history in geoscience and high-pressure research and provides a useful pathway for the exploration of high-pressure structures. In this study, we have examined the behavior of SnO$_2$ at exoplanetary conditions using OMEGA and OMEGA EP. In addition to its role as an analog material, tin oxide is also of interest since its high-pressure phases have been predicted to be potential ultra-incompressible materials.

Our ramp compression experiments use a target package consisting of a thin sample foil sandwiched between a diamond pusher and a LiF window. The OMEGA laser is used to ablate the diamond front surface driving a ramp compression wave into the sample. A quasi-monochromatic x-ray source is generated by irradiating a metal foil to create predominately He-like x rays. In-situ x-ray diffraction is performed using the powder x-ray diffraction image-plate (PXRDIP) diagnostic. The pressure is determined from measurements of the free-surface velocity of the sample/lithium fluoride interface.

We have carried out density-functional-theory calculations that indicate that tin oxide is expected to transform from the orthorhombic cotunnite-type phase to the hexagonal Fe$_2$P-type structure above 200 GPa. Our measured x-ray diffraction data for SnO$_2$ are $\sim$300 GPa and are shown in Fig. 152.42. We observe multiple diffraction lines from the sample, indicating that we have the potential to constrain high-pressure structures and equations of state at multi-megabar pressures for SnO$_2$. Analysis to distinguish among the possible crystal structure at this pressure is ongoing. In future work, we will extend these measurements to higher pressure as well as study the behavior of other dioxides including GeO$_2$ and SiO$_2$.

Dynamics of Magnetic Reconnection in High-Energy-Density Plasmas
Principal Investigators: W. Fox (Princeton Plasma Physics Laboratory); D. Schaeffer and A. Bhattacharjee (Princeton University); G. Fiksel (University of Michigan); and D. Haberberger (LLE)

We have developed and conducted experiments on the OMEGA and OMEGA EP Laser Systems to study the related phenomena of magnetic reconnection and collisionless shocks. Magnetic reconnection occurs when regions of opposite directed magnetic fields in a plasma can interact and relax to a lower-energy state; it is an essential plasma-physics process that governs the storage and explosive release of magnetic energy in systems such as the Earth's magnetosphere, the solar corona, and magnetic-fusion devices. The energy thereby liberated can produce heat flows and can enable the acceleration of a large number of particles to high energies. During previous NLUF reconnection experiments on OMEGA EP, we also unexpectedly observed the formation of magnetized collisionless shocks. Like reconnection, collisionless shocks are common in space and astrophysical systems. Magnetized shocks form from the nonlinear steepening of a magnetoacoustic wave and convert highly kinetic supersonic inflows to high-pressure subsonic outflows. In the process, they can also accelerate particles to...
extremely high energies. Those experiments were recently published in Ref. 1 and are under review in Ref. 2 as part of an invited talk at the 58th APS DPP meeting.

In this campaign we successfully carried out two experimental shot days on OMEGA EP and one shot day on OMEGA. The experiments on OMEGA EP utilized a new flat-foil platform for studying Biermann-mediated reconnection, while the experiments on OMEGA ported the magnetized colliding plasmas platform developed previously by our group and first published in Ref. 3.

The experiments on OMEGA EP used magnetic fields generated by the Biermann battery effect as the seed field for reconnection. Two oppositely directed Biermann fields were driven by the interaction of two lasers with a flat-foil target. Unlike previous Biermann-mediated experiments, our experiments utilized a flat target with a gap between the laser spots to provide a low-density region where the two laser plumes collide. The experiments on OMEGA used a single MIFEDS-driven coil but were otherwise similar to our MIFEDS platform developed for our previous NLUF campaigns on OMEGA EP.

The first shot day (March 2017) on OMEGA EP used proton radiography and the recently developed single-channel electron spectrometer (SC-ESM) to study the interaction between expanding plasma plumes and the resulting electron acceleration. We obtained spectacular proton radiography images that measure the magnetic field (see Fig. 152.43). Proton radiographs showed significant generation of Weibel filamentation and plasmoid formation in the gap region, which indicates fast reconnection. The SC-ESM also observed a population of energized electrons. This population was consistent with reconnection-accelerated particles but could also be caused by laser–plasma interaction (LPI) effects.

The second shot day (July 2017) on OMEGA used the Thomson-scattering diagnostic to measure the temperature and density of a single plasma plume interacting with a background magnetic field and ambient plasma (see Fig. 152.44). These measurements complement previous efforts on OMEGA EP that used proton radiography and angular filter refractometry (AFR) to diagnose large-scale magnetic topology and density profiles. They also provided information on conditions necessary for magnetized collisionless shock formation. We successfully obtained data at a range of times, distances from the target, ambient plasma conditions, and magnetic-field conditions. The results filled in several previous gaps. The measurements directly give plasma parameters and agree well with the 2-D radiation–hydrodynamics code DRACO. They also indicate...
that the formation of magnetized shocks is sensitive to the configuration of the main laser beams. The experiments, and their extension to counter-streaming plumes, will be continued on an upcoming NLUF shot day.

The third shot day (August 2017) on OMEGA EP continued the previous experiments on Biermann-mediated reconnection. We used AFR, which was unavailable on the previous shot day, proton radiography, and the SC-ESM to study the interaction of the colliding plumes. We obtained excellent AFR images, which show a complex interplay between magnetic turbulence and reconnection (see Fig. 152.43). The combination of proton and AFR images are currently being analyzed and compared to particle-in-cell simulations. We saw similar energized electron populations as on the first OMEGA EP day, but we also confirmed using the sub-aperture backscatter (SABS) diagnostics that LPI may play a significant role in energizing particles. This will be pursued in upcoming NLUF experiments.

![Figure 152.43](image)

*Figure 152.43*

Interaction of two Biermann-generated magnetic plumes 1.8 ns after laser ablation on OMEGA EP. The gap in the flat-foil target can be seen between the laser spots. (a) Proton radiography reveals the formation of Weibel filaments at the edge of the Biermann fields and plasmoids in the gap region. (b) AFR also shows the formation of filamentation at the plume edges and in the gap region.

![Figure 152.44](image)

*Figure 152.44*

Streaked Thomson-scattering measurements of the interaction of a piston plasma with a magnetized ambient plasma: (a) electron plasma wave feature and (b) ion-acoustic wave feature. The scattered spectra show that the ambient ions are swept up by the piston and form a hot, high-density downstream region.
Magnetized Accretion Shocks and Magnetospheres in the Laboratory
Principal Investigator: P. Hartigan (Rice University)
Co-investigators: C. C. Kuranz, G. Fiksel, J. Levesque, and R. Young (University of Michigan); J. Foster and P. Graham [Atomic Weapons Establishment (AWE)]; A. Frank (University of Rochester); A. Liao (Rice University); C. K. Li and R. D. Petrasso [Massachusetts Institute of Technology (MIT)]; and D. H. Froula (LLE)

The goal of our NLUF campaign is to use the magnetoinertial fusion electrical discharge system (MIFEDS) on OMEGA to create analogs of magnetized, hypersonic plasma flows that are ubiquitous in astrophysics. In previous experiments we explored several related topics, including the dynamics of magnetized star-forming clouds, magnetized supersonic Kelvin–Helmholtz instabilities, stellar magnetospheric infall, and planetary magnetospheres. Our current experiments relate most closely to stellar wind interactions with atmospheric outflows from magnetized exoplanets. The experimental setup drives a supersonic flow that impinges upon a current-carrying wire, producing a bow shock. Ablation flow from the wire encounters the supersonic flow, creating a working surface where the flows meet. Altering the amount of current in the wire changes the strength of the magnetic field in the ablation flow.

Last year we succeeded in producing a layer of concentrated magnetic flux embedded in the strongly shocked plasma of the working surface. However, the optical imaging diagnostics we used did not spatially resolve this layer well, and parts of the shocked layers were optically thick, preventing us from directly probing this layer where the two flows collide. To improve upon this design, in FY17Q1 we deployed two new diagnostic techniques that both penetrate into the working surface. First, we used spatially resolved Thomson scattering as a means to map how the velocities and densities vary across the shocked layers. In the spectra shown in Fig. 152.45(a) both the ablation flow from the wire (left side of the image in blue) and the incident flow that is driven from the laser target (right side of image in blue) are traced as a function of position. The two flows collide to generate a shocked layer with high temperature and density. According to theoretical models, the field in the wire should be swept by the ablation flow into a compressed layer that coincides with the working surface.

To diagnose the field geometry and strength, we employed proton radiography (Fig. 152.46). The results, which are quite striking, show the two main characteristics predicted by the simulations: (1) a voided area caused by protons being deflected away from the wire by the toroidal field caused by the current in the MIFEDS wire, and (2) a caustic produced by the

Figure 152.45
(a) Electron wave spectra resolved across the magnetized shock and (b) derived quantities. The density profile indicates a shock at the $d = 1$-mm position of $\sim 100$-$\mu$m thickness at half maximum.
accumulation of magnetic flux in the working surface. These observations allowed us to characterize the thickness and density of the magnetized flux sheet. Interestingly, the flux density we derived from the proton radiography was significantly less than that predicted by pressure equilibration; diffusion of the magnetic field resulting from inertial ions is a likely cause of these differences.

This experiment will serve as a template for future studies of magnetized shock layers. Following the temporal evolution of the system is a promising area of research, and with small changes to the experimental setup we should be able to further improve the spatial resolution and possibly measure the degree of mixing between the two fluids within the working surface.

Our NLUF research currently supports the thesis preparation of three graduate students in laboratory astrophysics: A. Liao (Rice University), and J. Levesque and R. Young (University of Michigan).

**Influence of Plasma Density on the Generation of Hundreds of MeV Electrons via Direct Laser Acceleration**

Principal Investigators: A. E. Hussein, T. Batson, K. Krushelnick, and L. Willingale (University of Michigan); A. V. Arefiev (UCSD); P. M. Nilson, D. H. Froula, R. S. Craxton, A. Davies, and D. Haberberger (LLE); and H. Chen and G. J. Williams (LLNL)

The OMEGA EP Laser System was used to study the acceleration of electrons to many times the ponderomotive energy by high-energy, picosecond-duration laser pulses interacting with an underdense plasma target. A high-intensity picosecond pulse propagating through underdense plasma will expel electrons along its path, forming a positively charged plasma channel. Electrons that are injected into this channel can gain significant energy through direct acceleration by the laser field. This acceleration mechanism is known as direct laser acceleration (DLA). Experiments on the OMEGA EP laser employed four of the chamber beams to study, optimize, and diagnose the influence of plasma density on the DLA mechanism. The existence of an optimal plasma density for the generation of high-energy, low-divergence electron beams was demonstrated. This result is consistent with results from 2-D particle-in-cell (PIC) simulations using the code *EPOCH*.

A schematic of the experiments on OMEGA EP is given in Fig. 152.47. A long-pulse UV beam (2.5 ns, 1200 J) ionized a CH flat-foil target to generate an expanding plasma plume. The backlighter beam (1 ps, 400 J) interacted with the plasma plume in an oblique geometry to generate a channel and accelerate an electron beam. The plasma density was varied by changing the timing between the long-pulse UV beam and the short-pulse (backlighter) beam. The sidelighter beam (1 ps, 200 J) was focused onto a Cu foil to generate a proton probe via target normal sheath acceleration (TNSA) for imaging electromagnetic fields onto a stack of radiochromic film (RCF). Shadowgraphy imaging, polarimetry of magnetic-field formation, and plasma density measurements by angular filter refractometry (AFR) were fielded using the 4ω optical diagnostic probe. The electron spectrum along the axis of the short-pulse beam was measured.

![Figure 152.46](a) Real and (b) synthetic 3-MeV proton radiographs with a fiducial, indicating the actual location of the wire (blue). The synthetic image (b) was constructed by placing a flux layer in the deprojected magnetic-field model with the position and thickness indicated by the electron plasma wave–derived shock structure (see Fig. 152.45). The apparent position of the flux layer (brimmed hat) defines the magnetized flux layer coincident with the shocked gas.

![Figure 152.47](Experimental configuration. A long-pulse beam generated an expanding CH plasma plume as the target of the short-pulse beam for channel formation and electron acceleration by direct laser acceleration.)
using an absolutely calibrated magnetic spectrometer [electron positron proton spectrometer (EPPS)].

Previous experiments explored the formation and evolution of plasma channels on an RCF film stack. In these images, the upward deflection of DLA-accelerated electrons was observed as a result of the refraction of the laser pulse in the plasma-plume density gradient. AFR measurements of plasma-plume expansion were used to extract a 2-D Gaussian density profile. This density profile agreed reasonably with density profiles of a CH expanding plasma plume obtained using the 2-D hydrodynamic code SAGE.

Two-dimensional PIC simulation conditions were designed to match the OMEGA EP Laser System. A 1.053-μm-wave-length laser, with coincident focal spots of 3.4 μm and 17 μm, with intensities of $3.776 \times 10^{19}$ W/cm$^2$ and $2.81 \times 10^{18}$ W/cm$^2$, respectively, was simulated. The plasma density profile was defined by the Gaussian profile from AFR measurements. Simulations were performed with a resolution of 30 cells per wavelength in the longitudinal direction ($x$) and six cells per wavelength in the transverse direction ($y$). The laser was linearly polarized in $y$ and propagated in $x$. Ions were treated as mobile, and the plasma (electron) density was varied to simulate the temporal evolution of the plasma plume in experiments. Simulated densities ranged between $0.1 n_c$ and $0.008 n_c$, where the quoted value corresponds to the peak density along the laser trajectory and $n_c$ is the critical density.

Simulated electron spectra are shown in Fig. 152.48(a) and reveal the existence of an optimal plasma density for electron acceleration to energies exceeding 230 MeV. Experimental electron spectra are shown in Fig. 152.48(b). The quoted density in this figure refers to the peak electron density along the trajectory of the main interaction beam, as estimated by SAGE simulations and AFR measurements. Experimental measurements revealed an enhancement of peak electron energy at a peak density of $0.01 n_c$. The enhancement of peak electron energy, up to nearly 600 MeV, was significantly higher than predicted by 2-D simulations.

The proton-probe diagnostic captured electromagnetic-field structures of the plasma channel. Although the initial time $t_0$ of the short-pulse interaction cannot be exactly determined because of timing jitter of the order of 20 ps, relative timing between each film in the RCF stack can be calculated from proton time-of-flight calculations. An example radiograph is given in Fig. 152.49(a), where clear formation and filamentation can be observed using a normal film pack, 8 cm from the interaction region. The centroid of the resultant electron beam on another RCF stack was found to deflect above the axis of laser propagation, as shown in Fig. 152.49(b). This is consistent with upward refraction of the laser pulse in the density gradient. Information about beam pointing and angular divergence as a function of plasma density could be extracted from the on-axis RCF stack. Beam divergence tends to decrease as a function of plasma density, with some evidence of channel filamentation, creating multiple beamlets, at lower densities.

These experiments demonstrated the existence of an optimal plasma density for the generation of high-energy electron beams by the interaction of a high-intensity picosecond pulse with underdense plasma. Experimental results also indicate a relation between plasma density and beam divergence. Continued work will focus on the role of quasi-static channel fields on
electron energy enhancement, beam pointing, and divergence to elucidate the mechanisms and action of DLA at different plasma densities.

**High-Energy-Density Chemical Physics and Planetary Evolution**

Principal Investigator: R. Jeanloz (University of California, Berkeley)

Co-investigators: M. A. Millot, D. E. Fratanduono, P. M. Celliers, and J. H. Eggert (LLNL); S. Brygoo and P. Loubeyre (CEA); and T. R. Boehly, G. W. Collins, and J. R. Rygg (LLE)

During FY17, our international research team conducted two campaigns with diamond-anvil cell targets on the OMEGA laser (DirectDAC17A and 17B) for a total of 16 shots. The configuration is a direct-drive geometry (Fig. 152.50), with up to 12 beams delivering up to 6 kJ in 1-ns square pulses to the 1-mm aperture in the tungsten carbide seats holding the diamond anvils. VISAR (velocity interferometer system for any reflector) velocimetry and streaked optical pyrometry monitored the shock propagation in the sample pressure chamber to diagnose the pressure–density equation of state and the optical properties (reflectivity, absorption coefficient) using a quartz reference. Most of the shots were dedicated to the study of the metallization of hydrogen (deuterium) using the combination of high precompression (6 to 13 GPa) and double shock.

![Proton probe image of a plasma channel. (b) Radiochromic film along the axis of short-pulse (backlighter) beam propagation can be used to study electron beam pointing and divergence as a function of plasma density.](image1)

![Typical VISAR velocimetry record and (b) sketches of the re-shock experimental configuration showing the arrival of the elastic wave at the diamond/sample interface ($t_1$), followed by the arrival of the inelastic wave ($t_2$) and the catchup of these two waves inside the D$_2$ sample, transforming it in an opaque fluid. (c) Impedance mismatch at the arrival of the shock wave at the sample/diamond interface induces the generation of a second shock wave into the deuterium sample ($t_4$). The bright fringes after $t_4$ indicate that deuterium becomes a good reflector upon reshock, indicative of high electrical conductivity.](image2)
compression to access high-density states in the vicinity of the predicted first-order transition from the insulating molecular fluid to a metallic atomic fluid.

Our team also successfully demonstrated the feasibility of using diamond-anvil cell targets on the OMEGA EP laser during a one-day campaign. Preliminary analysis of the data indicates that OMEGA EP makes it possible to reach higher shock pressures using longer pulse durations. Future experiments might also benefit from the excellent pulse-shape capability available on OMEGA EP to generate multishock compression and map out the metallization transition of several key planetary constituents.


**Particle Acceleration Resulting from Magnetically Driven Reconnection Using Laser-Powered Capacitor Coils**

Principal Investigator: H. Ji (Princeton University)

Co-investigator: L. Gao (Princeton Plasma Physics Laboratory)

Magnetic reconnection is a ubiquitous astrophysical phenomenon in which magnetic energy is rapidly converted into plasma kinetic energy in the form of flow energy and thermal energy as well as nonthermal energetic particles. Energy particles are often regarded as an observational signature of the magnetic reconnection, which can be a more-efficient generation mechanism than other competing processes such as collisionless shocks. Despite its long history, most laboratory work in this area has focused on the mechanisms of fast reconnection as well as the generation of plasma flow and thermal energy during magnetic reconnection, mostly resulting from the limitations in either experimental setups or diagnostic capabilities. The goal of our research is to build an effective new platform to achieve and measure conspicuous particle acceleration by magnetically driven axisymmetric reconnection using laser-powered capacitor coils. In FY14, our team successfully measured and reported the first direct measurement of the magnetic fields generated by these laser-powered capacitor coils. With strong magnetic fields approaching the MG level and by tuning plasma parameters, we will be able to access magnetically driven, collisionless reconnection for efficient particle acceleration.

In FY17, we successfully carried out one experimental shot day on OMEGA EP to study the above-mentioned physics goals. A schematic of the experimental setup on OMEGA EP is shown in Fig. 152.51. The main interaction target is comprised of two parallel copper plates connected by two copper wires. Two OMEGA EP 2.5-kJ, 1-ns laser pulses pass through the laser entrance holes on the front plate and are focused on the back foil, generating a beam of superthermal hot electrons. The hot electrons stream onto the front plate and build up an electrical potential between the plates. This in turn drives

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Figure 152.51

Experimental setup for the recent reconnection experiments on OMEGA EP based on the laser-powered capacitor coils. BL: backlighter; NTA: near-target arm; TIM: ten-inch manipulator; TPS: target positioner system; TTPS: TIM target positioner system.
large currents in both wires and creates magnetic reconnection because antiparallel magnetic-field lines exist in the middle plane. Another OMEGA EP 50-J, 1-ns laser pulse irradiates a plastic target and generates a tenuous background plasma for the reconnection region.

A suite of existing OMEGA EP diagnostics were fielded to identify quantitative confirmation of reconnection. Ultrafast proton radiography was utilized to probe the reconnection process at various times with high spatial and temporal resolutions. The new OMEGA EP high-resolution x-ray spectrometer (HiRes) was fielded to monitor Cu Kα emission from both front and back Cu plates. Three fixed port x-ray pinhole cameras viewed the time-integrated x-ray emission from the entire target. The Osaka University electron spectrometer was used to measure the particle spectrum from reconnection. The HiRes data showed strong Cu Kα emission from the front Cu plate, indicating hot electrons generated during the main interaction streaming onto the front plate and exciting these K-shell emissions. This measurement is a direct confirmation of the charging mechanism for building the electrical potential between the two plates. The hot electrons therefore have energies of at least 8 keV. Detailed atomic physics analysis is ongoing to understand the spectrum of hot electrons generated when the two main laser pulses hit the back copper foil. This will facilitate a better understanding of the mechanism for creating these strong magnetic fields. Our electron spectrometer data consistently showed a peak in the electron power spectrum with maximum electron energy ~1 MeV from reconnection. These experimental results are being compared with our PIC simulations.

Principal Investigators: K. Krushelnick, P. Campbell, L. Willingale, and G. Fiksel (University of Michigan); and P. M. Nilson and C. Mileham (LLE)

Recent experiments were conducted on the OMEGA EP Laser System to study a magnetic-reconnection geometry established by firing a short-pulse laser alongside a long-pulse UV beam onto solid targets. A 1-ns, 1250-J UV beam was focused to an intensity of $2 \times 10^{14}$ W/cm² onto a 25-μm-thick copper target. Misalignment of temperature and density gradients in the ablated plasma plume generated azimuthal Biermann battery magnetic fields (of the order of MG). As this long-pulse–produced plasma developed, a 10-ps pulse containing 500 J was focused to relativistic intensity ($I > 10^{18}$ W/cm²) in close proximity. In contrast to the slowly expanding Biermann battery fields ($V \approx c_s$), relativistic currents driven by the short-pulse laser generate a strong azimuthal magnetic field (of the order of 10 to 100 MG) that spreads radially with a velocity near the speed of light. This dramatic difference in scales yields a highly asymmetric field geometry, with the rapidly expanding short-pulse–generated field driving into a quasi-static Biermann battery field.

Proton radiography was implemented to diagnose the magnetic-field dynamics of the interaction. A second short-pulse laser (300 J in 1 ps) accelerated protons via the TNSA mechanism from 20-μm-thick gold targets to energies exceeding 60 MeV. A stack of RCF detected the deflection of the proton beam by the electromagnetic fields of the target. The time of flight for a proton depends on the kinetic energy. Because of the Bragg peak in the proton stopping power, each layer in the RCF stack detects a different energy and therefore a different time in the interaction.

As shown in Fig. 152.52, the proton radiography captures the rapid expansion of the short-pulse–generated magnetic field...
into the quasi-static Biermann field. At \( t_0 \), the Biermann field has evolved for 750 ps when the short-pulse IR beam arrives on target and the focal-spot separation is 1.25 mm. After 11.6 ps, there is a local enhancement in proton flux along the edge of the Biermann field structure as the fields meet and interact, indicating a strengthening of the field gradient. At later times, there is evidence of outflow structures emanating from the magnetic-field interaction.

In future experiments, the relative beam timing and focal-spot separation will be tuned to optimize and study these features of the interaction. The copper targets will be replaced with a lower-\( Z \) material to improve the quality of the proton radiography, and the dimensions will be increased to mitigate target-edge effects. In addition to the experiments, our future work will include a quantitative analysis of the proton-radiography results, moving toward full 3-D PIC simulations.

**Properties of Magnetohydrodynamic Turbulence in Laser-Produced Plasmas**
Principal Investigator: D. Q. Lamb (University of Chicago)

During the first shot day (19 April 2017) of our experiments (TDYNO NLUF Campaign) we used the OMEGA laser to study the turbulent dynamo amplification of magnetic fields, a ubiquitous process in astrophysical systems. The experiments employed a platform (Fig. 152.53) similar to the one we fielded on OMEGA for our very successful first TDYNO Campaign (FY15–FY16), during which we demonstrated nonlinear amplification by turbulent dynamo for the first time in a laboratory environment. The main goal of our first shot day was to map the time history of the magnetic-field amplification, probing the various phases of turbulent dynamo: the kinematic phase, when fields are weak with respect to the turbulent motions; the nonlinear phase, when the Lorentz force back-reacts on the plasma’s momentum; and the saturation phase, when the magnetic energy reaches a sizable fraction of the kinetic reservoir and amplification stops (Fig. 152.54). We designed the experimental platform using numerical simulations on one of the nation’s leadership supercomputers (Fig. 152.55). The platform is uniquely suited to generating turbulent plasmas in the large magnetic Reynolds numbers regime, where the dynamo can operate. The configuration consists of two diametrically opposed foil targets that are backlit with temporally stacked beams, delivering 5 kJ of energy on each side in a 10-ns span. The beams drive a pair of counter-propagating plasma flows that carry seed magnetic fields generated by the Biermann battery effect. The flows propagate through a pair of grids that destabilize the flow and define the driving scale of the turbulence. The flows meet at the center of the chamber to form a hot, turbulent interaction region where the weak seed magnetic fields are amplified to saturation values.

![Figure 152.53](U2283JR)

**Figure 152.53**
Experimental platform for the NLUF campaign to study turbulent dynamo amplification, a ubiquitous astrophysical process. The target assembly consists of two polystyrene foils and a pair of meshes, held together by a pair of cylindrical shields and a “tuning-fork” stalk. The foils and meshes were carefully designed to optimize the conditions in the interaction region for turbulent field amplification. The shields and flaps protect the interaction region, the imploding \( \text{D}^3\text{He} \) capsule, and the diagnostics from direct view of the laser spots.

![Figure 152.54](U2386JR)

**Figure 152.54**
Simulated time history of the magnetic field’s strength (in gauss) in the interaction region, showing the kinematic (exponential), nonlinear, and saturation phases of the dynamo for root-mean-square (rms) and peak values of the magnetic field. The arrows denote the different times at which we fired the proton radiography diagnostic: the 1.5-ns cadence allowed us to collect enough experimental data to temporally resolve the rise of the magnetic field.
With a modest redesign of our original experimental platform, we were able to develop a faster and more-accurate alignment procedure; this enabled us to perform 15 shots during our first shot day—a record number given the complexity of our platform. The shots yielded a wealth of experimental data. The diagnostics we fielded allowed us to fully characterize the turbulent interaction region and quantify its plasma properties in space and time. More specifically, x-ray imaging enabled us to directly visualize the formation and evolution of the turbulent region. From the x-ray intensity fluctuations we reconstructed the density power spectrum of the magnetized turbulence and inferred its power law. Moreover, the spatially resolved spectrum from the Thomson-scattering diagnostic yielded ion-acoustic and electron features—at different times—that allowed us to characterize the plasma properties, including the ion and electron temperatures; the bulk flow and turbulent velocities; and the electron density in a 1.5-mm field of view. Therefore, we were able to probe for the first time both the turbulent interaction region and the inflowing plasma.

Finally, to recover the time history of the magnetic-field amplification, we fielded the proton radiography diagnostic in all shots (Fig. 152.56). Utilizing the proton-radiography data and the novel magnetic-field mapping techniques we have developed, we are able to reconstruct the strength and topology of the magnetic field during all phases of dynamo amplification, with a 1.5-ns cadence (Fig. 154.54). This plenitude of experimental data is under analysis and promises to greatly expand our understanding of the puzzle that is astrophysical turbulence.

**Creation of a Magnetized Jet Using a Hollow Ring of Laser Beams**

Principal Investigator: E. Liang (Rice University)

Progress toward the objectives of this NLUF project as listed in the original application far exceeded expectations in its second year. We carried out a one-day joint OMEGA and OMEGA EP laser experiment in November 2016, using 20 OMEGA beams to form a hollow ring focal pattern to create a magnetized jet from a flat plastic target. Some of the targets were doped with 2% Fe. The hollow ring radius varied from 800 to 1200 μm. Twelve shots were successfully completed, half of which were joint shots with OMEGA EP. The Thomson-scattering (TS) diagnostic was used to measure the on-axis electron and ion density, temperature, and flow velocity at 2.5 mm from the laser target for each shot. The TS results confirmed the predictions of new FLASH 3-D simulations, after allowance was made for the TS probe beam’s heating of the electrons. The on-axis densities, temperatures, and velocities are comparable for the 800- and 1200-μm radius rings, suggesting that the optimal OMEGA ring radius lies between 800 and 1200 μm. In six of the shots, we used 3-MeV and 14-MeV monoenergetic protons from D³He capsule implosions to measure the magnetic fields in the jet via proton radiography. For the other six shots we used continuum OMEGA EP protons to probe the magnetic fields. Both sets of proton radiography images gave consistent results, which compare favorably with 3-D FLASH simulation predictions.
of approximately megagauss axial fields for both the 800- and 1200-μm cases. This shows that the plasma properties of the hollow ring jet, including its collisionality and MHD properties, will be significantly impacted by the self-generated magnetic fields. We used an x-ray framing camera to take x-ray images of the jet at 1-ns time intervals. The images show well-collimated laminar outflows for both the 800- and 1200-μm cases. This is in excellent agreement with 3-D FLASH predictions. When we compare the x-ray images of the undoped and 2% Fe-doped jets, we find that the Fe-doped jet is narrower and the x-ray emission stronger, as predicted by FLASH. A poster on the 2016 results was presented at the Omega Laser Facility Users Group (OLUG) Workshop in April 2017. The Principal Investigator (PI) also gave invited talks on the NLUF results at the NNSA symposium in Chicago in April 2017 and at LANL in July 2017. Updated results on both the OMEGA data and FLASH simulations will be presented at the American Physical Society Division of Plasma Physics (APS DPP) meeting in Milwaukee in October 2017. A National Ignition Facility (NIF) Discovery Science proposal based on the OMEGA results and FLASH simulations was submitted to the NIF in September 2017.

**Development of a Photoionized Plasma Experiment on OMEGA EP**

Principal Investigator: R. C. Mancini (University of Nevada, Reno)

Experiments on basic high-energy-density science on OMEGA EP provide a unique opportunity to create states of matter at extreme conditions of temperature, density, and radiation flux relevant to astrophysics. The focus of this project is to study the fundamental atomic and radiation physics of plasmas driven by a broadband intense flux of x rays; i.e., photoionized plasmas. Most laboratory work performed to date on high-energy-density laboratory plasmas pertains to collisional plasmas; i.e., those where electron collisional processes play a dominant role in the plasma ionization and atomic kinetics. However, relatively little attention has been paid to studying and understanding the fundamental properties of laboratory photoionized plasmas, where both photoionization and photoexcitation driven by a broadband x-ray flux are dominant. These plasmas are important for understanding a myriad of astrophysical sources including x-ray binaries, active galactic nuclei, and the accreting disks formed in the vicinity of black holes. The information that we obtain on these objects is based on the analysis of spectroscopic measurements recorded by orbiting telescopes such as Chandra and XMM-Newton. Yet, the analysis of the spectra relies on sophisticated atomic and radiation physics models that have been developed only on a best-theory effort. Therefore, there is a critical need for performing systematic photoionized plasma laboratory experiments to benchmark theory and modeling codes.

We are developing a silicon photoionized plasma experiment on OMEGA EP in which a plastic-tamped silicon foam is ionized by the 30-ns-duration, broadband x-ray flux produced by the “gatling-gun” radiation source. This source is comprised of three copper hohlraums that are driven by three OMEGA EP beams, each delivering 4 kJ of UV energy in a 10-ns square pulse shape. The laser beams sequentially illuminate one hohlraum at a time, thereby producing an x-ray drive characteristic of 90-eV radiation temperature for a time period of 30 ns. The silicon sample is placed at a distance of 7 mm from the source. It has an initial mass density of 100 mg/cm³ and a thickness of 0.1 mm and is coated with two 1-μm-thick layers of plastic. Heated by the x-ray flux, the silicon sample expands by a factor of 10 and ionizes into the L-shell range of silicon ions, thereby producing a photoionized plasma with an atom number density of a few times 10¹⁹ atoms/cm³ and a relatively uniform spatial distribution.

In the first phase of the experiment’s development, we performed OMEGA EP shots in which the expansion and ionization of the tamped silicon foam was monitored with a gated imaging x-ray spectrometer that recorded the L-shell self-emission of the plasma. The radiation temperature of the gatling source was measured with a VISAR diagnostic. The expansion to 1-mm thickness and ionization of the silicon sample were both confirmed by the observations during the experiment.

Figure 152.57 displays L-shell emission from line transitions in B- and Be-like silicon ions recorded at t = 6 ns and t = 9 ns. No measurable line emission in these ions is noted before t = 6 ns. The x-ray flux starts at t = −15 ns and lasts until t = +15 ns. Therefore, these observations are taken during the second half of the x-ray drive duration. This is reasonably consistent with the pre-shot expectation based on radiation–hydrodynamics modeling of the experiment. This observation is now being used to refine the numerical simulation. The spectra demonstrate the formation of a highly ionized silicon plasma driven by the x-ray flux. It also provides data to extract the electron temperature.

The next step in the experiment’s development is to perform x-ray K-shell transmission spectroscopy of the silicon photoionized plasma with a streaked instrument and a separate titanium backlighter driven by the fourth beam of OMEGA EP. The transmission spectroscopy will permit the extraction of the silicon charge-state distribution and an independent check on the temperature from the L-shell emission spectra analysis.
Explorations of Inertial Confinement Fusion, High-Energy-Density Physics, and Laboratory Astrophysics
Principal Investigators: R. D. Petrasso, C. K. Li, and J. A. Frenje (MIT)
Co-investigators: F. H. Séguin and M. Gatu Johnson (MIT)
Graduate students: N. Kabadi, B. Lahmann, H. Sio, R. Simpson, G. Sutcliffe, and C. Wink (MIT)
Undergraduate Student: M. Manzin (MIT)

MIT work in FY17 included a wide range of experiments applying proton radiography, charged-particle spectrometry, and neutron spectrometry methods developed by MIT and collaborators to the study of laboratory astrophysics, high-energy-density physics (HEDP), and ICF plasmas. This was an outstanding year for the HEDP Division’s scientists and students and their work on NLUF-related research. Based on NLUF work resulting in the development of the multiple-monoenergetic-particle source (MMPS) and its application to a wide range of physics experiments involving the observation and measurement of laboratory plasma phenomena and associated electromagnetic fields through radiography and other means, Drs. C. K. Li, R. D. Petrasso, and F. H. Séguin (Fig. 152.58) were chosen as recipients of the APS 2017 John Dawson Award for Excellence in Plasma Physics Research. The MMPS has been used by several MIT students in research critical to their outstanding Ph.D. theses. These include Drs. M. Manuel and M. Rosenberg, who received the APS Rosenbluth Outstanding Doctoral Thesis Award in 2014 and 2016, respectively, and Dr. A. Zylstra, who has now been nominated for the same award. In addition to the HEDP Division’s own students, investigators from many other institutions have enlisted MIT’s collaboration in their HEDP experiments on OMEGA to gain the unique information supplied by radiography with the MMPS. Those institutions include LANL, the University of Rochester, LLNL, the University of Chicago, Princeton University, UCSD, the University of Michigan, and the University of Oxford.

The MMPS is a laser-driven capsule containing D$_3$He fuel that produces monoenergetic charged-fusion products including 3.0-MeV protons, 14.7-MeV protons, and 3.6-MeV alpha particles during a 0.1-ns time interval, used either as a backlighter for multiple-monoenergetic-particle radiography or as a source of monoenergetic particles for other nonimaging experiments. The many subjects MIT has studied with the MMPS during the NLUF program include ICF experiments, plasma jet propagation, and magnetic reconnection, utilizing the MMPS as a backlighter for radiography, and quantitative studies of ion stopping and ion-electron equilibration in plasmas. These NLUF-developed techniques have also recently been ported to the NIF (LLNL).

The MIT scientists shared the Dawson Award with three scientists at other institutions (A. MacKinnon, Lawrence Livermore National Laboratory; M. Borghesi, The Queen’s University, Belfast; and O. Willi, Heinrich Heine University, Düsseldorf). Those scientists had applied a different kind of proton source to radiography of plasmas, using TNSA, in which a laser pulse strikes a planar target and generates a strong electric field, charge separation, and a resultant picosecond-duration proton beam with a continuous energy spectrum.
In the meantime, other important NLUF accomplishments have included experimental measurements of nuclear reactions relevant to stellar and big-bang nucleosynthesis using high-energy-density plasmas on OMEGA and extensive experiments on kinetic physics. Many of these experiments have been enhanced by MIT student H. Sio’s development of the new particle x-ray temporal diagnostic (PXTD) on OMEGA for simultaneous time-resolved measurements of several nuclear products as well as the x-ray continuum produced in HEDP. The PXTD system makes it possible, for the first time, to take accurate and simultaneous measurements of x-ray emission histories, nuclear reaction histories and their time differences along with measurements of \( T_i(t) \) and \( T_e(t) \) for studies of kinetic, multi-ion effects, and ion-electron equilibration rates in ICF plasmas.

**Study of Magnetized Collisionless Shocks in Laser-Produced Plasmas**

Principal Investigator: A. Spitkovsky (Princeton)
Co-investigator: C. M. Huntington (LLNL)

The FY17 MagShock EP Campaign was dedicated to the study of collisionless magnetized shocks in ablated plasma flows and exploration of new experimental concepts. Collisionless shocks commonly form in supernova remnants and in the heliosphere. The shock thickness is determined by the Larmor radius of the incoming protons, and the collisional mean free path must be much longer. The setup is shown in Fig. 152.59(a). The experiments used the OMEGA EP Laser System in which a 3-D–printed Helmholz coil powered by MIFEDS was inserted; two targets were mounted on MIFEDS. A 400- to 800-J, 1-ns pulse was used to ablate plasma that propagated along the coil’s magnetic field (this component is called “background” plasma). A 1.3-kJ, 1-ns pulse was used to drive a fast-flow orthogonal to the magnetic field [this component is called “piston” plasma; see Fig. 152.59(b)]. The interaction between the flows was expected to drive a compression in the background plasma and the magnetic field [see Fig. 152.59(c)]. At a strong enough drive, this compression becomes a collisionless shock. This compression was diagnosed using proton radiography with TNSA from a short 1-ps laser pulse. The protons were recorded on CR-39 film, which was our primary diagnostic. On some shots, the \( 4\omega \) optical probe was also utilized.

The experiments resulted in the detection of a magnetized collisionless shock propagating through the plasma. The main feature of the magnetic compression in the data was the appearance of a white band in the proton image, indicating additional deflection of the protons. The band was followed by a sharp caustic of enhanced proton concentration [Figs. 152.60(a)–152.60(c); shots separated by 2 ns]. The band and the caustic propagated at 300 km/s. The thickness of the band made it possible to constrain the magnetic compression ratio to 2.3, and the caustic was interpreted as the signature of the contact discontinuity between the piston and compressed background plasma [Fig. 152.60(d) lineout along the blue line in Fig. 152.60(b)]. This compression ratio corresponds to a Mach-3.4 shock. The shock is in the collisionless regime since the mean free path of the background protons is larger than the size of the plasma. These results were confirmed on several shots that performed the time-series study. Several time offsets were also tried before the piston plasma was launched, presumably probing different background densities. Extensive numerical simulations of the experiment were performed with 3-D PIC simulations [Figs. 152.59(c) and 152.59(d)], including simulated proton radiography through the fields of the simulation. The experimental results agree quite well with the predictions of the simulations. These findings will be presented at APS DPP Meeting in October 2017 as an oral contribution.
In a second campaign in August, we attempted to probe the configuration through the orthogonal direction, resulting in rotation of the entire apparatus. The results are currently being analyzed. A clear shock signature was not seen, most likely because the background plasma density was lower than expected.

We thank the OMEGA EP personnel for their assistance in planning and executing this campaign.

**Demonstration of Talbot–Lau X-Ray Deflectometry**

**Electron Density Diagnostic in Laser–Target Interactions**

Principal Investigator: D. Stutman (Johns Hopkins University)

An experiment aimed at demonstrating Talbot–Lau x-ray deflectometry (TXD) in high-energy-density plasmas was performed on the Multi-Terawatt (MTW) laser, where a laser-produced x-ray backlighter was used to illuminate a Talbot–Lau interferometer and obtain electron density maps of solid targets. These experiments confirmed that the TXD technique has the potential to become a basic and widespread diagnostic for HEDP experiments. To benchmark the technique in the HED environment, the interferometer must be tested in the presence of a plasma target. To this aim, a CH foil will be radiated by three UV OMEGA EP beams. A fourth beam will be used to obtain an x-ray backlighter source to illuminate the Talbot–Lau deflectometer, which will provide an electron density map of the ablated foil. Simulation and theory have failed to accurately predict the electron density profile from the ablation dynamics. TXD will provide electron density information in ranges not available today (<$n_c$) and will therefore advance the field of HED.

The proposed two-year research includes the implementation of a TXD on the OMEGA EP laser (FY18Q4). Preparatory experiments have been performed on MTW to test and optimize the x-ray backlighter. The first year has focused on diagnostic conceptual design and preparation of the experiment in collaboration with S. P. Regan and C. Stoeckl from LLE. The diagnostic design and implementation in the OMEGA EP laser are underway and being directed by C. Sorce and C. Mileham from the LLE Experimental Support Group. The conceptual design is shown in Fig. 152.61. Similar to the MTW setup, all the gratings will be mounted on a common optical base or rail. The common rail will allow us to accurately pre-align the gratings on site during pre-shot setup. The experiments will be defined and optimized in collaboration with P. A. Keiter (University of Michigan), so that the main goal of benchmark-
ing TXD as an electron density diagnostic for HEDP will be compatible with the goal of advancing the understanding of plasma ablation of plasma-irradiated foils.

**Hot-Electron Generation with 10^{16}.W/cm^2 Infrared Lasers in the Shock-Ignition–Relevant Conditions**

Principal Investigator: M. S. Wei (GA)
Co-investigators: C. M. Krauland (GA); S. Zhang, J. Li, and F. N. Beg (UCSD); C. Ren, W. Theobald, D. Turnbull, D. Haberberger, C. Stoeckl, R. Betti, and E. M. Campbell (LLE); and J. Trela and D. Batani (CELIA)

Shock ignition (SI) is an alternative ICF scheme that achieves ignition with a strong convergent shock launched by a high-intensity (~10^{16}.W/cm^2) laser spike at the end of the low-intensity (~10^{14}.W/cm^2) assembly pulse. SI spike-pulse energy coupling to the fusion target is uncertain because of laser–plasma instability (LPI) such as filamentation, stimulated Brillouin scattering (SBS), stimulated Raman scattering (SRS), two-plasmon decay (TPD), and the resultant hot-electron generation. Therefore, it is important to characterize the LPI and the hot-electron beam energy, temperature, and divergence in SI-relevant conditions to further assess the SI scheme.

This GA-led NLUF HotEScaleEP-17A Campaign conducted on OMEGA EP in August 2017 is an extension of our previous experiments (HotEScaleEP-15A), where we successfully demonstrated the propagation of an infrared (IR) laser beam (1.053 μm, 100 ps, up to 2.5 kJ, ~10^{16} W/cm^2) over a long-scale plasma (1.4-keV, 450-μm scale length) and the generation of a hot-electron beam with small divergence and moderate energies (T_{hot} ~ 90 keV). This year we continued using the IR laser with two beams (up to 2.3 kJ) in the co-propagating (co-prop) geometry to extend the pulse duration to 200 ps. The IR pulse (i.e., 200 ps with the co-prop beams or 100-ps single beam) had a nominal vacuum laser intensity of ~2 \times 10^{16} W/cm^2 at the quarter- or tenth-critical density (n_c) surface. We used three-layer (25 μm CH/20 μm Cu/50 μm Al) solid disk targets designed to stop hot electrons up to 200 keV. A hot, large-scale CH plasma was created by a 2-ns, 2.2-kJ UV laser (B3 or B4) with 750-μm distributed phase plates (DPP’s) (I ~ 2 \times 10^{14} W/cm^2) to mimic the large corona plasma in SI. At 1.5 ns after the start of the UV pulse, the single or co-prop IR beams were injected. A suite of optical (the 10-ps 4ω probe and the newly available OMEGA EP SABS) and x-ray diagnostics [e.g., SCI, ZVH, bremsstrahlung MeV x-ray spectrometer (BMXS), and hard x-ray detector (HXRD)] were utilized to characterize LPI and the hot-electron energy, spectral, and angular distribution.

In our HotEScaleEP-17A experiments, we found that hot-electron generation by the co-prop IR beams (200 ps) was sensitive to the plasma temperature as shown in Fig. 152.62. Such dependence was not observed in the single IR beam (100-ps) interaction. With the higher-temperature (1.5-keV) plasma, the Cu K_{α} photon yield produced by the co-prop IR beams increased 170% compared to the data with 1.0-keV plasma. Meanwhile, the measured hard x-ray signal was also increased about 130% to 180% in the 17- to 200-keV energy range. HXRD data suggested a hot-electron temperature of less than 100 keV. The AFR image obtained from the 4ω probe diagnostic showed that the co-prop IR beams propagated beyond the n_c/4 in high-temperature plasma and produced a bright self-emission spot at the n_c surface as shown in Fig. 152.63. Resonance absorption may have contributed to hot-electron generation at the n_c surface. The bright self-emission spot was not observed from the interaction of either the co-prop IR beam with the 1.0-keV plasma or the single IR beam (100 ps) with the same density scale plasma with a temperature from 1.0 to 1.4 keV. The AFR image also captured the density perturbation between n_c/10 and n_c/4 surface as the result of strong nonlinear LPI’s. OMEGA EP SABS recorded spectrally resolved sidescattering of the IR beam in the 400- to 750-nm spectrum range showing SRS. The 4ω self-emission observed at the n_c surface indicated that SRS and TPD were saturated since the co-prop IR beams propagated beyond the n_c/4 surface (TPD and SRS boundary). The data analysis and PIC modeling are ongoing.

![Figure 152.62](U2252JR)

The Cu K_{α} yield (normalized to the IR beam energy) from the interaction of the single and two co-prop IR beams with low- (1.0-keV) and high- (1.4- to 1.5-keV) temperature plasmas was measured using a zinc von Hamos x-ray spectrometer.
FY17 Laboratory Basic Science (LBS) Program

Sixteen LBS projects previously approved for FY17 target shots were allotted Omega Laser Facility shot time and conducted a total of 199 target shots at the Omega Laser Facility in FY17 (see Table 152.VIII). The FY17 LBS experiments are summarized in this section.

During FY17, LLE issued a solicitation for LBS proposals to be conducted in FY18. A total of 28 proposals were submitted. An independent committee reviewed and ranked the proposals; on the basis of these scores, 16 proposals have been allocated 21 shot days at the Omega Laser Facility in FY18. Table 152.IX lists the approved FY18 LBS proposals.

Exploring Pair Plasmas and Their Applications
Principal Investigator: H. Chen (LLNL)

In FY17, an LLNL/SLAC/University of Michigan team continued this project on OMEGA EP with one LBS shot day. These experiments use the short-pulse beams to produce jets of electron–positron antimatter pairs. The FY17 campaign focused

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on measuring the pair yield enhancement with nanostructured targets. The experiments successfully demonstrated that the laser-positron energy conversion can be improved by using novel structured targets. A total of 12 shots were performed.

The OMEGA EP short-pulse beams (~1 kJ in 10 ps) irradiated 1-mm-thick Au targets with and without the nanostructure on the laser-interaction surface. It was found that for the same laser energy, positron yields and acceleration both were increased dramatically by using a nanostructure. This finding is important to future experiments and applications using laser pair jets. Previous experiments used primarily gold targets and showed that quasi-monoenergetic relativistic positron jets are formed during high-intensity irradiation of thick gold targets, and also that these jets can be strongly collimated using the magneto-inertial fusion electrical discharge system (MIFEDS). The external field produces a 40-fold increase in the peak positron and electron signal. The positron yield was found to scale as the square of the laser energy, and the yield also increases with the Z of the target material. Together with the nanostructured target yield enhancement, these favorable scalings are expected to enable the laboratory study of relativistic pair plasmas to aid one’s understanding of some of the most exotic and energetic systems in the universe.

**Shock-Ignition Timing Measurements on OMEGA**


The objective of the Shock-Ignition Timing Campaign was to measure the ablation pressure and hot-electron preheat from the ignitor spike. Small [650-μm outer diam (OD)] CH shells were irradiated with two pulse shapes—one with a spike and one without a spike—as shown in Fig. 152.64(a). The compression of the shell resulting from the ignitor shock was measured by the x-ray framing camera (XRFC). This is illustrated in Fig. 152.64, where (a) represents the ablation-front trajectories of the implosions with and without a spike, while (b) is an image from the XRFC showing how the ablation front and hot-

### Table 152.IX: LBS experiments approved for target shots at the Omega Facility in FY18.

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spot radius are measured. A maximum ignitor shock ablation pressure of \( \sim 180 \text{ Mbar} \) was inferred from a post-shot LILAC simulation approximately constrained to the data.

One way to experimentally determine preheat by hot electrons is to measure the premature expansion of the shell before the ignitor shock breaks out. This preheat causes the hot spot to emit earlier than it would without preheat. In Fig. 152.64(a), the implosion with the ignitor shock starts emitting earlier than the no-spike case. Understanding the relative importance of hot-electron preheat versus ignitor shock breakout time on the earlier emission is the subject of an ongoing investigation.

Measuring the Nerst Effect and the Thermal Dynamo
Principal Investigator: J. R. Davies (LLE)

The objective of this experiment was to observe the dynamo effect caused by cross-field heat flow. The dynamo effect occurs because of sheared rotation of plasma around magnetic-field lines, which acts to twist the field lines, generating a component in the direction of the rotation. The dynamo effect has been extensively studied in the context of astrophysical jets.

A magnetic field in a plasma moves not only with the bulk plasma motion but also with the electron heat flow, which can be ascribed to the field being preferentially frozen to the electrons responsible for heat flow because of their lower collision frequency. In magnetized plasma there exists heat flow perpendicular to the temperature gradient and the magnetic field, known as cross-field heat flow. Cross-field advection of the field has largely been ignored, but it could readily lead to a dynamo effect in the absence of plasma rotation.

To generate a thermal dynamo, we placed a carbon disk inside a MIFEDS coil, giving an axial magnetic field, and irradiated the edge of the disk with ten OMEGA beams, five from each side, creating a radial temperature gradient, which will lead to an azimuthal cross-field heat flow and the generation of an azimuthal magnetic field. To probe the magnetic field, we used \( \text{D}_3\text{He} \) protons passing through a grid on one side of the target.

Unfortunately, MIFEDS started to trigger erratically and only one magnetized shot was obtained, at maximum field, which deflected all of the protons out of the detection angle, as shown in Fig. 152.65. The grid is not visible, which may be a result of being attached to the target or simply being too close to the target or backlighter. Although this one shot indicates that the thermal dynamo effect does occur, shots with lower magnification, a lower magnetic field, and a working grid setup would be required to confirm and quantify the effect.
**Characterizing Pressure Ionization in Ramp-Compressed Materials with Electron-Induced Fluorescence**

Principal Investigator: S. Jiang (LLNL)
Co-investigators: Y. Ping, R. F. Smith, A. Jenei, and J. H. Eggert (LLNL)

This campaign used one day on OMEGA EP to measure ionization in compressed materials as a function of density, using K-shell fluorescence spectroscopy. The K-shell line emissions were induced by hot electrons generated through short-pulse laser–solid interactions. The high pressure was achieved by ramp compression using the long-pulse drivers while keeping the temperature low. A large, thick target was used to avoid heating from the short pulse, as shown in Fig. 152.66(a). The configuration is intended to compress the material up to 1.5\(\times\) to 2\(\times\) its original density without raising the thermal ionization effect. There is still little consensus on pressure ionization under these conditions despite extensive theoretical and experimental efforts.

A schematic of the experimental setup is displayed in Fig. 152.66(a). The main diagnostic used in this campaign was the high-resolution imaging x-ray Thomson spectrometer (IXTS). The Cu K\(\alpha\) and Co K\(\beta\) fluorescence lines were successfully observed with a high signal-to-noise ratio, as can be seen in Fig. 152.66(b). Under the designed experimental conditions, pressure ionization has a negligible effect on the innermost electron shells (K,L), but it can affect the M shell. Therefore, the Cu K\(\alpha\) line is not subject to an energy change and can be used as a reference. On the other hand, the Co K\(\beta\) line is more prone to a shift. Figure 152.66(b) shows the measured Co K\(\beta\) peaks under different driver energies and time delays. While some differences can be observed between the undriven and driven conditions, the energy shifts are small enough that they are close to the resolution of the spectrometer (~4 eV). In a future experiment, we will increase the driver energy to reach a higher pressure and also probe other materials.

**Time-Resolved Measurement of the Radiative Properties of Open L-Shell Zinc**

Principal Investigator: E. V. Marley (LLNL)

This campaign was designed to measure the emitted L-shell zinc spectrum from a well-characterized and uniform plasma for comparison to atomic kinetic models. Recent studies have shown a discrepancy between atomic kinetic models and high-Z M-shell spectral data. This study was done to test the accuracy of models for L-shell emission.

Planar buried layer targets were illuminated evenly on both sides (Fig. 152.67) to heat the sample composed of Ti and Zn. The sample was buried between two 6.7-\(\mu\)m-thick layers of Be, which inertially tamps the sample, slowing its expansion.

![Figure 152.66](image)

(a) A schematic of the experimental configuration. The inset shows the target geometry. (b) Measured IXTS spectra of Co K\(\beta\) (normalized with the peak intensity). The lineouts are from different driver conditions. An example of the raw image is shown in the inset in (b).
Time-resolved 2-D images of the target’s x-ray emission, viewed both face-on and side-on, were recorded using pinhole cameras coupled to framing cameras. The K-shell spectrum from the Ti was used to determine the electron temperature of the plasma. The time-resolved spectrum was recorded using a crystal spectrometer coupled to a framing camera to give a temporally and spatially resolved spectral measurement of the zinc L shell. This measurement will help to verify the uniformity of the plasma at different times during the experiment. All of the framing cameras, used for imaging as well as for spectroscopy, were co-timed on the shot day so that plasma conditions could be determined for the measured zinc L-shell spectra.

Two different pulse shapes were used: a 3.2-ns square pulse and a 3.0-ns square pulse with a 100-ps picket preceding it by 1 ns. The second pulse tested whether creating a pre-plasma before the main pulse would create a smoother interface, allowing for a more-efficient coupling of energy into the target. The results look promising. A complete set of data from all six co-timed diagnostics was recorded for both pulse shapes during the campaign at sample temperatures \( \pm 1 \text{ keV} \).

### Optimizing Backlighters for Imaging Low-Density Plasmas for Eagle Pillar Studies

Principal Investigator: D. A. Martinez (LLNL)
Co-investigators: J. Kane and R. F. Heeter (LLNL); and B. Villette and A. Casner (CEA)

The LBS Eagle Pillar experiments were designed to optimize the backlighter for imaging the plasma plume created by an ablated CH solid-density cone target in conditions similar to counterpart NIF Discovery science experiments. The CH target was driven for 30 ns using three Cu hohlraums (Fig. 152.68) heated in succession, which will eventually create
Equation of State, Structure, and Optical Properties of Silicates at Multi-Mbar Pressures for Super-Earth Mantle and Accretion Modeling

Principal Investigator: M. A. Millot (LLNL)
Co-investigators: F. Coppari, D. E. Fratanduono, and S. Hamel (LLNL); N. Dubrovinskaia and L. Dubrovinsky (Bayreuth University, Germany); and R. Jeanloz (University of California, Berkeley)

During FY17, our international research team conducted two campaigns at the Omega Laser Facility to investigate the equation of state, structure, and optical properties of silicates at multi-Mbar pressures using shock compression.

Following our previous study on the melting line of SiO$_2$ using stishovite crystals, we conducted a shock compression study of MgSiO$_3$ bridgmanite (perovskite) samples synthesized at Bayreuth University, Germany, to investigate the pressure–density–temperature equation of state and the melting temperature of MgSiO$_3$, representative of the material making up the interiors of rocky planets and the cores of gas giants. Preliminary analysis of the velocity interferometry (VISAR) and streaked optical pyrometry (SOP) data indicates that we successfully observed reflecting shocks up to 15 Mbar.

Another set of experiments used OMEGA EP pulse-shaping capabilities to launch a carefully timed series of two steady shocks into stishovite samples to make it possible to determine the atomic structure of SiO$_2$ at Neptune core conditions using x-ray diffraction. Composite three-beam pulse shapes enabled us to generate the complex drive conditions, while focusing Beam 1 on a Ge foil was used to generate a 1-ns pulsed x-ray source. Preliminary analysis indicates that x-ray diffraction (XRD) patterns (Fig. 152.70) for single- and double-shock compressed silica were obtained up to 9 Mbar.

(a) 0 ns (b) 30 ns (c) 35 ns

Figure 152.69
Time sequence of the target imaged through a pinhole array backlit with a U backlighter. Multiple images were obtained to remove the backlighter structure. The (a) 0-ns image had lower filtering but the (b) 30- and (c) 35-ns images had 200 $\mu$m Be.
**Laser-Driven Collisionless Shock Acceleration of Ions (E-Shock)**  
Principal Investigator: A. Pak (LLNL)  
Co-investigators: A. J. Link (LLNL); D. Haberberger and D. H. Froula (LLE); S. Tochitsky, C. Joshi, and F. Fiuza (University of California, Los Angeles)

This LBS shot day explored the acceleration of ions into narrow energy distributions caused by the reflection from the strong electrostatic field of a collisionless shock wave. As Fig. 152.71(a) indicates, a near-critical-density CH target was first produced by the x-ray drive generated from the ablation of a gold foil by a 1-ns-long laser pulse (UV Beamline 3). After waiting for the CH target to expand to a peak density of $+10 \times 10^{21} \text{cm}^{-3}$, the ultra-intense backlighter beam was used to irradiate the target and drive the collisionless shock wave. Nearly simultaneously, the sidelighter beam produced a beam of protons through target normal sheath acceleration (TNSA) to radiograph the shock-formation process.

In these experiments, the pinhole of the Thomson parabola ion energy (TPIE) diagnostic was replaced with a new $5$-mm-wide $\times 0.25$-mm-long slit to extend the angular acceptance of the diagnostic. The electric bias of the TPIE was turned off, and differential filtering of $50 \mu m$ and $300 \mu m \text{Al}$ was used at the image-plate detector to differentiate between ion species. TPIE data from this configuration are shown in Fig. 152.71(b).

In this experiment, narrow energy distributions of both protons and carbon ions were observed to be accelerated to similar velocities of $\sim 0.25 c$. The acceleration of disparate charge to mass ratio ion species to similar velocities is consistent with acceleration from the moving near-relativistic electric field associated with a collisionless shock wave. A distribution of protons centered at $36 \text{MeV}$ (velocity $= 0.28 c$), and with an energy range $\Delta E/E$ of $\sim 30\%$ was observed, as well as a distribution of $\text{C}^{6+}$ ions centered at $308 \text{MeV}$ (velocity $= 0.23 c$) and $\Delta E/E$ of $\sim 12\%$. The difference in velocity between the two ion species is thought to arise from the remaining sheath field of the expanded target, which preferentially accelerates the lighter species. Analysis of the radiography and accelerated beam profile data is in progress.

**Astrophysical Collisionless Shock Experiments with Lasers**  
Principal Investigator: H.-S. Park (LLNL)  
Shot Principal Investigator: G. F. Swadling (LLNL)

Experiments ACSEL-17A and 17B investigated the physical processes that lead to the formation of astrophysical collisionless shocks. These shots continued a broad, long-running, cross-institutional collaboration. A total of 26 target shots were completed in two shot days, primarily investigating interactions between beryllium ablation outflows; this material was selected to provide a low-Z, single-species plasma, which greatly simplifies Thomson-scattering interpretation and analysis, while maintaining the large collisional scale lengths required to observe the development of interpenetrating flow instabilities. In these experiments, the OMEGA beams heat the surfaces of a pair of opposed planar disk targets (see Fig. 152.72), ablating counter-propagating plumes of high-velocity (up to $1.5 \times 10^6 \text{ms}^{-1}$), high-temperature ($\sim \text{keV}$) plasma. The outflow parameters are such that the coulomb mean-free path for interflow collisions is long, but the interaction of the flows is still susceptible to the growth of the interstream instabilities that are believed to mediate the formation of collisionless shocks.

This year, experiments focused on taking spatially resolved Thomson-scattering measurements across the interaction...
region to quantitatively investigate the development of the ion-Weibel instability. Thomson scattering was combined with proton radiography measurements; a D³He exploding-pusher capsule provided a dichromatic (3.3- and 14.4-MeV) proton source for radiography, probing the plasma at two separate times during each experiment. Images were recorded on CR-39, with processing and analysis of the CR-39 plates carried out by collaborators at MIT.

The OMEGA Thomson-scattering diagnostic records both ion-acoustic and electron plasma wave features of the Thomson-scattering spectrum. Analysis of the detailed shape of these spectra made it possible to extract information about the spatial variation in electron temperature, electron density, ion temperature, and flow velocity across the interaction of the two flows. The primary goal in FY17 was to make a direct measurements of spatial density modulations of the ion flows and the underlying electron density of the plasma. High-quality data were recorded and are expected to provide a wealth of data on the development of the ion-Weibel instability in these experiments; examples of the spectrograms are shown in Fig. 152.72. Striations in the intensity of the ion-acoustic feature along the probe beam are characteristic of the development of the ion-Weibel instability, while the presence of modulations in the electron plasma wave data suggests that the ion-Weibel instability is in the nonlinear growth phase.

**Measuring Strong Plasma Shock-Front Structure Using Thomson-Scattering Imaging**

Principal Investigator: H. G. Rinderknecht (LLNL)
Co-investigators: H.-S. Park and J. S. Ross (LLNL); and D. H. Froula (LLE)

This series of shots was designed to directly measure for the first time the spatial structure of a strong shock front in a plasma. The experiments were intended to develop a platform for kinetic plasma studies using the new gas-jet system on OMEGA to quantify collisional phenomena in high-energy-density plasmas and to benchmark high-fidelity physics codes. These experiments were also the first use of the gas-jet system on OMEGA.

The KineticShockLBS-17A Campaign on 24 August used the new gas-jet system to inject a column of hydrogen or neon gas into the OMEGA target chamber. A 1-μm Si₃N₄ foil positioned near the gas-jet nozzle was driven by ten beams with 2.5 kJ in 0.6 ns, exploding the foil to drive a strong shock into the low-density (~5 × 10¹⁸ cm⁻³) gas. A 526.5-nm probe beam with 40 J in a 100-ps impulse was injected normal to the ablator foil at a 4- to 6-nsec delay from the drive beams, and Thomson-scattered light from the probe was imaged along the probe axis. The imaged region was 3.25 to 4.75 mm from the ablator, with a resolution of 20 μm. Despite operational difficulties with the gas-jet system on its first shot day, six shots were completed and excellent data collected on all shots for which the gas-jet successfully operated.

Figure 152.73 shows ion-acoustic wave (IAW) and electron plasma wave (EPW) Thomson-scattering images from shot 86801. Fits to the EPW data demonstrate the characteristics of strong shock formation: heating of electrons in the pre-shock region, followed by an increase in density as the ion shock forms. IAW data appear to show streaming protons in advance of the shock front, heating, and slowing down on the pre-shocked plasma; analysis of these results is underway. These exciting results will be presented in an invited talk at the APS DPP meeting in October 2017 to demonstrate the high value of this platform for future kinetic plasma studies on OMEGA.
Electron Energization During Magnetic Reconnection in High-Energy-Density Plasmas
Principal Investigator: M. J. Rosenberg (LLE)

The MagReconnection-17A shot day on OMEGA (14 March 2017) through the LBS program successfully diagnosed the spectrum of energetic electrons generated in laser-plasma experiments in which magnetic reconnection was driven. These experiments utilized a well-established platform for studying the generation, interaction, and reconnection of magnetic fields in plasmas created by the interaction of multiple laser-produced plasma plumes adjacent to each other using foil targets. The energization of particles during the annihilation of magnetic fields is a common process in astrophysical plasmas, but it is poorly understood and has rarely been investigated in the laboratory. The new, compact single-channel electron spectrometer microscope (SC-ESM) obtained electron spectra over the energy range of ~50 to 300 keV.

Spectra obtained perpendicular to the foil and parallel to the reconnection current sheet that supports the magnetic fields are shown in Fig. 152.74. Experiments with two beam spots [Figs. 152.74(b) and 152.74(c)] drove magnetic reconnection, while experiments with only one beam spot [Fig. 152.74(d)] did not. Energetic electron spectra were measured, with characteristic temperatures of ~30 to 50 keV. Notably, a single-beam experiment [Fig. 152.74(d)] generated energetic electrons, suggesting that additional mechanisms beyond magnetic reconnection, such as laser–plasma instability (LPI), may be producing the energized particles. In addition, Thomson-scattering measurements were successfully obtained to diagnose plasma conditions in the reconnection region, and monoenergetic proton radiography was used to confirm the interaction and reconnection of magnetic fields, as have been observed in previous experiments. Another LBS shot day has been awarded to this campaign, during which we will attempt to determine the source of energetic electrons and obtain spectra unambiguously from magnetic reconnection by eliminating LPI.

Absolute Equation-of-State Measurements from Spherically Converging Shock Waves on the OMEGA Laser
Principal Investigator: A. M. Saunders (LLNL)
Co-investigators: T. Doeppner and R. Nora (LLNL); W. Theobald (LLE); and A. Jenei, D. Swift, J. Nilsen, and R. W. Falcone (Lawrence Berkeley National Laboratory, University of California, Berkeley)

X-ray Thomson scattering (XRTS) is an experimental technique that directly probes the physics of warm dense mat-
ter by measuring electron density, electron temperature, and ionization state. XRTS in combination with x-ray radiography offers a unique ability to measure the absolute equation of state (EOS) of material in extreme conditions.

The OMEGA GbarIPD-17A Campaign took XRTS and x-ray radiography measurements from directly driven carbon-containing spheres compressed to electron densities of the order of $1 \times 10^{24} \text{ cm}^{-3}$ and temperatures of $\sim 30$ eV. X-ray radiography measurements were obtained for both plastic (CH) and high-density carbon (HDC) spheres. Fifty-two beams compressed the spheres, and six beams drove a foil backlighter. The x-rays from the foil backlighter were observed in transmission through the sphere using a gated x-ray framing camera (Fig. 152.75). They show that the shock front travels inward as predicted by simulations. The radial lineouts make it possible to obtain the shock velocity. A more-complicated analysis of post-shock density at each time step will also be performed; the combination of density and shock velocity will allow for an absolute measurement of the EOS.

In conjunction with the radiography measurements, XRTS spectra were obtained from HDC spheres. A zinc He\textsubscript{α} x-ray source was used to scatter x-rays from the imploding spheres at a scattering angle of $135^\circ$. The scattered x-rays were collected by a crystal spectrometer in conjunction with a gated x-ray framing camera. Figure 152.76 shows an example of the raw data collected and a lineout of one of the strips. The

Figure 152.74
(a) A sample raw image-plate scan from SC-ESM; [(b)–(d)] spectra from various laser drive conditions.

Figure 152.75
X-ray radiography measurements of shock-compressed CH spheres (shot 85691). (a) Example of the raw data obtained from the imploding CH spheres and [(b),(c)] radial lineouts from several of the radiography images, each at a different point in the implosion time.
XRTS data will provide an independent measurement of the mass-averaged electron temperature of the imploding sphere. The temperature measurement will further constrain the EOS measurement obtained from the radiography analysis. In summary, the data obtained in this campaign shed light on the EOS of matter under compression and support EOS measurements previously taken on the NIF.

**Determining the High-Pressure Properties of Silicon Carbide Using Decaying Shocks in In-Situ X-Ray Diffraction**

Principal Investigators: R. F. Smith (LLNL) and J. K. Wicks (Johns Hopkins University)

The goal of this campaign was to determine the high-pressure properties of single-crystal SiC along the Hugoniot, using a combination of shock decay and nanosecond x-ray diffraction techniques. Silicon carbide is an important material in geology and planetary science. It may be a host of reduced carbon in the Earth’s interior since it is found in rocks from the mantle and in inclusions in deep diamonds. It also occurs in meteorites and impact sites. The target design in Fig. 152.77 is modeled off previous campaigns on OMEGA. The raw active shock breakout (ASBO) (shock velocity) and streaked optical pyrometer (SOP) (shock front thermal emission) data make it possible to determine the pressure–temperature onset of melt (see Fig. 152.78). During this shot day the powder x-ray...
diffraction image-plate (PXRDP) diagnostic was also used to determine the high-pressure crystal structure (pre-melt).

**Experimentally Constraining the High-Pressure Thermal Conductivity of Iron**
Principal Investigator: R. F. Smith (LLNL)
Co-investigator: J. K. Wicks (Johns Hopkins University)

High-pressure thermal conductivity is one of the most important and yet most difficult to measure physical property of materials. Within the Earth’s interior the thermal conductivity \( k \) of Fe and Fe-rich alloys at core pressure–temperature conditions (135 to 360 GPa, 2500 to 5000 K, respectively) is a key parameter for heat transport models and plays an important role in determining the temperature profile and energy balance of our planet. The thermal conductivity of the Earth’s core remains poorly constrained because of the extreme difficulty in making thermal transport measurements under the relevant pressure and temperature conditions. Two experimental studies published in Nature in 2016 report values of \( k \) for Fe that vary by a factor of 7 at \( \sim 130 \) GPa (34 → 225 W/mK) (Refs. 28 and 29). The goal of the OMEGA experiments was to constrain the thermal conductivity at high pressures, using a ramp-compression platform previously developed on OMEGA\(^3\) (Fig. 152.79), where (1) stagnating plasma simultaneously launches a ramp-compression and heat wave in the sample, (2) a “cool” ramp-compression wave runs ahead of the heat wave (the sample pressure can be constrained using VISAR), and (3) LLE’s SOP makes it possible to measure the heat-wave transit time. The raw ASBO/SOP data provide velocity and thermal transport information through stepped Fe samples (Fig. 152.80). Analysis is underway to translate this data into a measurement of thermal conductivity.

![Figure 152.79](image1)

(a) Target design used to (b) measure heat flow in a ramp-compression stepped Fe sample.

![Figure 152.80](image2)

(a) ASBO and (b) SOP data for 0.84-1.94-/2.74-\( \mu \text{m} \) Fe step samples provide sufficient information to constrain the high-pressure thermal conductivity of Fe.
Recovery of Dynamically Compressed Samples
Principal Investigator: C. E. Wehrenberg (LLNL)
Co-investigators: S. Zhao and M. Meyers (UCSD); and B. Remington and A. Krygier (LLNL)

This LBS campaign studied the deformation response of a variety of materials to shock compression. Sample materials were mounted onto the front of steel recovery tubes that were in turn mounted on a ten-inch manipulator (TIM). A single OMEGA beam is used to drive a shock into the sample material, and the sample remains in the recovery tube after the shot so that it can be recovered for further ex-situ study. Two recovery tubes and a VISAR target were fielded for each shot. Consequently, a large data set of 12 recovered samples and six VISAR traces was produced using only a half-day of shots.

A wide variety of samples were recovered during this campaign. Previous iterations of this campaign have been very successful in studying the deformation response of semiconductors (Si and Ge) to shock compression, producing a series of high-profile papers on pressure-shear–induced amorphization. The FY17 campaign studied GaAs, graphite, and olivine and generated the first dynamic compression data on a new class of materials—high-entropy alloys. These samples will be taken to Oak Ridge National Laboratory for TEM (transmission electron microscopy) study.

Charged-Particle Stopping Power and Scattering Measurements in a Warm Dense Plasma
Principal Investigator: A. B. Zylstra (LANL)

The dE/dx Campaign is developing a platform to perform high-precision measurements of charged-particle stopping power (dE/dx). Stopping power in dense plasmas is important for ICF self-heating and propagating burn, particularly for particles near and below the peak in dE/dx (“Bragg peak”). While the stopping power has been measured in hot-spot–relevant plasmas using an exploding pusher platform, the current data cannot distinguish between models of interest; for example, the Maynard–Deutsch and Brown–Preston–Singleton theories. This campaign uses shocked-foam targets probed by a separate source of fusion particles to achieve higher precision. The dE/dx-17A shot day demonstrated the viability of this platform, shown in Fig. 152.81. Good data using 3-MeV protons from the D–D fusion reaction were acquired, shown in Fig. 152.81(b). A similar downshift is observed in both shocked (warm) and undriven (cold) foam, which is expected from stopping theory. Future experiments will be modified to use the lower-velocity particles, particularly D3He, which are more sensitive to dE/dx and will be able to differentiate between stopping models in this regime.

FY17 LLNL Omega Facility Experimental Programs

In FY17, LLNL’s Indirect-Drive Inertial Confinement Fusion (ICF-ID) and High-Energy-Density (HED) Physics Programs conducted numerous campaigns on the OMEGA and OMEGA EP Laser Systems. Overall these LLNL programs led
413 target shots in FY17, with 282 shots using only the OMEGA Laser System and 131 shots using only the OMEGA EP Laser System. Approximately 27% of the total number of shots (78 OMEGA shots and 35 OMEGA EP shots) supported the ICF-ID Campaign. The remaining 73% (204 OMEGA shots and 96 OMEGA EP shots) were dedicated to experiments for HED Physics. Highlights of the various ICF-ID and HED Campaigns are summarized in the following reports.

In addition to these experiments, LLNL Principal Investigators (PI's) led a variety of LBS Campaigns using OMEGA and OMEGA EP, including 85 target shots using only OMEGA and 70 shots using only OMEGA EP.

Overall, LLNL PI's led a total of 568 shots at LLE in FY17. In addition, LLNL PI's also supported 30 NLUF shots on OMEGA and 46 NLUF shots on OMEGA EP, in collaboration with the academic community.

**Indirect-Drive Inertial Confinement Fusion Experiments**

*Hydrodynamic Response from Nonuniformities in Plastic, High-Density Carbon, and Beryllium*

Principal Investigator: S. J. Ali
Co-investigators: P. M. Celliers, S. W. Haan, S. Baxamusa, M. Johnson, H. Xu, N. Alexander, H. Huang, V. A. Smalyuk, and H. F. Robey

The goal of the Capseed Campaign (comprising Capseed 17A, 17B, and 17C) is to measure shock-front velocity nonuniformities in ICF ablator materials and quantify the level of nonuniformity caused by intrinsic effects. This is done using the OMEGA high-resolution velocimeter (OHRV) to obtain velocity maps of the optically reflecting shock front following release of the ablator material into either PMMA [poly(methyl methacrylate)] for the warm experiments or cryogenic deuterium for the cryogenic experiments. For three half-days in FY17 the focus was twofold: (1) complete measurements on the impact of oxygen heterogeneity and oxygen mitigation layers for glow-discharge polymer (GDP); and (2) begin measuring velocity nonuniformities on deep release from Be, GDP, and high-density carbon (HDC) into D\(_2\) with improved velocity sensitivity.

Performance and yield from fusion capsules at the National Ignition Facility (NIF) are highly dependent on the uniformity of the capsule implosion, and hydrodynamic instabilities are a significant source of performance degradation during the implosion. A possible explanation for unexpectedly large in-flight modulations observed during NIF capsule implosions was a surface oxygenation of GDP; laboratory tests of GDP samples under controlled conditions confirmed the heterogeneous surface oxygenation effect. In FY16 the OHRV was used to test this idea further by obtaining 2-D velocity maps for both oxygen-modulated and unmodulated samples. Modulated samples showed clear evidence of the propagation of a rippled shock wave as a result of the photo-induced oxygen heterogeneity. To mitigate this effect, the target fabrication team proposed depositing a 20-nm oxygen barrier layer of alumina. Tests of the mitigation in the 17A and 17B campaigns determined, via OHRV measurements on warm GDP samples, that this barrier layer introduced no additional perturbations in the shock velocity. In 17C the velocity roughness on deep release from GDP into D\(_2\) with and without this barrier layer was also measured and no significant difference was determined. The velocity nonuniformities in both samples were close to the detection limit of the diagnostic, as described in Fig. 152.82.

The remaining shots in 17A were used to measure velocity nonuniformities slightly below the first shock level in HDC and

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*Figure 152.82*

Velocity maps from (a) alumina-protected and (b) unprotected glow-discharge polymer (GDP) releasing into D\(_2\). Root mean square (rms) velocity roughness was 8.6 m/s for (a) and 6 m/s for (b), with a diagnostic detection limit of 6 m/s.
at the first shock level in sputtered Be, both releasing into D₂. The velocity nonuniformity for the Be shots varied significantly (35±6 m/s and 24±6 m/s) but in both cases, was ~3 to 5× larger than would be expected from the surface roughness. Nonuniformity in HDC was measured at 7 and 10 Mbar and was found to decrease with increasing pressure (from 124±17 m/s for 7 Mbar to 85±17 m/s for 10 Mbar) but was again a few times larger than predicted from surface roughness alone. These experiments are continuing into FY18, with further cryogenic measurements planned.

**Diamond Sound-Speed Measurements Between 8 and 14 Mbar**

Principal Investigator: A. Fernandez-Panella  
Co-investigators: D. E. Fratanduono and P. M. Celliers

This half-day on the OMEGA laser was designed to collect high-quality data on the sound speed of diamond in the multi-Mbar range, where currently little data exist, for the purpose of constraining equation-of-state models. It was the continuation of the DiamondSS-15B Campaign, where two data points were obtained at 10 Mbar.

The DiamondSS-17A Campaign used planar targets and a direct-drive configuration with a CH ablator, a quartz pusher, and two targets side by side [quartz (standard) and diamond]. The velocity profiles at the free surface were recorded using VISAR as the primary diagnostic (Fig. 152.83). Throughout the day, the laser drive energy was changed in order to probe different shocked pressure states in diamond. The laser pulse was designed to produce a small pressure perturbation that propagated through both the standard and the sample. The wave propagation analysis described in Ref. 34 was applied to extract the sound speed of both materials, quartz and diamond, from their velocity profiles (Fig. 152.84).

The results indicate that sound-speed measurements are sensitive enough to constrain EOS models. The Livermore LEOS table 9061 shows good agreement with the data. Further measurements in an extended pressure range are desirable in order to better constrain the models. The results of this campaign are being used to optimize the design for the DiamondSS-18A Campaign, where Be ablators will be used instead of CH to reach higher pressure states; a different pair of etalons will be chosen to increase the accuracy of the velocity measurements to enable extraction of not only the sound speed but also the Grüneissen coefficient.

![Diamond Sound-Speed Measurements](image)

Thomson-Scattering Measurements from Foam-and Gas-Covered Gold Spheres

Principal Investigator: G. F. Swadling  
Co-investigators: J. S. Ross, M. Rosen, K. Widmann, and J. D. Moody

The FoamCoSphere-17A Campaign performed foam- and gas-covered high-Z sphere experiments illuminated in direct-drive geometry to investigate atomic physics models, radiative properties of the laser-spot plasma, and the interpenetration of...
multi-ion species plasmas relevant to ICF indirect-drive–ignition hohlraums. These experiments use laser irradiation at $10^{14}$ to $10^{15}$ W/cm$^2$, similar to the intensities found in hohlraums fielded on the NIF.

For the foam experiments the Au sphere was embedded in a low-density (3.8-mg/cm$^3$) CH foam [Fig. 152.85(a)]. For the gas-covered experiments, the spheres were located inside a gas bag filled to 1 atm of propane, or 1 atm of a 70/30 mix of propane and methane, to achieve initial electron densities of 4.0% of the critical density of the 3ω drive beams and to mimic the interaction of the hohlraum Au wall with the low-density hohlraum fill gas. Significant target development work was required to prepare the open-geometry foam-covered targets; we look forward to leveraging this target development work in future shot days.

The plasma temperature and density at various radial positions in the blowoff plasma are characterized using Thomson scattering, while x-ray flux from the gold sphere was recorded using the Dante and DMX soft x-ray spectrometer diagnostics. The laser beams use a shaped laser pulse (1-ns square foot, 1-ns square peak) designed to pre-ionize the gas/foam before the main drive pulse.

The electron temperature and density, the plasma-flow velocity, and the average ionization state are measured by fitting the theoretical Thomson-scattering form factor to the observed data. An example of the Thomson-scattering data from ion-acoustic fluctuations is shown in Fig. 152.85. Continued data analysis and simulations are in progress to better understand the plasma evolution and heat transport.

**Study of Interpenetrating Plasmas on OMEGA**
Principal Investigator: S. LePape

This campaign is designed to study the dynamics of plasma interpenetration in an environment relevant to hohlraums used on the NIF to drive HDC capsules. The question this campaign is trying to answer is whether a fluid description of plasma flows in a low-gas-fill hohlraum (<0.3 mg/cm$^3$ of helium gas in the hohlraum) is accurate, or if a kinetic description of the flows must be used. Following the first series of shots in 2016 that looked at the time evolution in one point in space in the gap between the two flows, this series of shots focused on 1-D spatially resolved Thomson-scattering data in addition to looking at the effect of helium gas density on the flows’ interaction. During this series of shots the helium gas density was changed as well as the ring material (carbon, aluminum, and gold) and the laser energy used to drive the target.

Two main diagnostics are fielded on these experiments: (1) a soft x-ray time-resolved imager looking at the self-emission of the plasmas along the ring axis (Fig. 152.86) and (2) a spatially resolved Thomson-scattering diagnostic (Fig. 152.87) to diagnose electron and ion temperature, flow velocities, and densities.

Time-resolved x-ray images [Figs. 152.86(c) and 152.86(d)] indicate that the helium gas holds the plasma expansion and as time goes by when helium is present, a bright layer appears, presumably being the helium compressed between the low-Z and high-Z plasma.

Figure 152.87 shows a 1-D spatially resolved spectrum acquired on a carbon/carbon shot without helium, providing

![Figure 152.85](image)

(a) Au spheres (1-mm diameter) centered in a 5-mm-diam foam. (b) Example of the Thomson-scattering ion feature, where the Thomson-scattering $k$ vector is directed radially, giving sensitivity to the ablation velocity of the Au plasma.
Figure 152.86
(a) Schematic of a gas-bag target. The ring and rod with which the laser interacts are enclosed into a gas bag that holds the helium gas. (b) VISRAD calculation of the laser-deposited energy on the ring and high-density carbon (HDC) rod. [(c),(d)] Time-resolved x-ray images of the plasma self-emission. The inner feature is the carbon rod; the outer feature is the low- or high-Z ring. In (c), gold ablates into HDC without a helium fill; in (d), gold ablates into HDC with 0.2 mg/cm³ of helium gas.

Figure 152.87
One-dimensional spatially resolved Thomson-scattering data of a carbon plasma flowing into a carbon plasma.
information on plasma temperature, densities, and plasma stagnation length. High-quality data were obtained. The results of this campaign are being analyzed.

**Measurements of Anisotropy in Non-LTE Low-Density Iron–Vanadium Plasmas**

Principal Investigator: L. C. Jarrott

Accurate characterization of the effects of geometrical anisotropies on K-shell line emission is very important for improving line-ratio–based temperature measurements in low-density, non-LTE (local thermodynamic equilibrium) plasmas. OpticalDepth-17A built on the OpticalDepth-16A Campaign, which established a working platform for accurately characterizing low-density, mid-Z, non-LTE plasmas. Specific goals of this platform included a characterization of heat conduction in the nominal target point design by comparing a layered sample target material to a mixed sample target material. Additionally, this campaign attempted to improve on-target laser drive efficiency by virtue of a picketed pulse shape to improve the laser–target interaction interface. Both OpticalDepth Campaigns used three laser–target configurations, varying target angle to verify the accuracy of data acquired. The primary target was a 10-μm-thick, 1000-μm-diam beryllium tamper with an embedded volumetrically equal mixture of iron and vanadium, 0.2 μm thick and 250 μm in diameter. The second target type was identical to the primary target except that the sample material was layered (50 nm Fe/100 nm V/50 nm Fe) rather than mixed. The third target type was a null target where the beryllium tamper contained no sample material. Three beam-target orientations were used over the course of 14 experimental shots (Fig. 152.88).

![Image](img.png)

Figure 152.88
(a) X-ray framing-camera images (all from TIM-3) using our three laser/target orientation configurations, showing a face-on view (LC-1), side-on view (LC-2), and 45° view (LC-3). (b) Spectrum measured by the MSPEC spectrometer of the x-ray emission from K-shell transitions in highly charged vanadium and iron using laser/target configurations 1 and 2.
In the first configuration (LC-1), the multipurpose gated x-ray spectrometer (MSPEC) situated in TIM-2 had an edge-on view of the target, while an identical MSPEC in TIM-6 had a face-on view. In the second configuration (LC-2), the target orientation with respect to TIM-2 and TIM-6 was reversed compared with the first configuration. In the third configuration (LC-3), all primary ten-inch-manipulator (TIM)-based diagnostics had a viewing angle of 45° with respect to target normal. Using multiple target-beam orientations resulted in an in-situ cross-calibration of the spectrometers and pinhole imagers. The data included simultaneous measurements of (1) time-resolved iron and vanadium K-shell spectra viewed from both the target edge and the target face, and (2) time-resolved images of the expanding plasma, viewed from both the target edge and target face, to infer plasma density for both layered and mixed sample target materials. The K-shell spectral data provided time-resolved electron temperature measurements of the expanding plasma, with preliminary analysis implying a plasma temperature above 2 keV.

**Imaging Electric-Field Structure in Strong Plasma Shocks**

Principal Investigator: H. G. Rinderknecht
Co-investigators: H.-S. Park, J. S. Ross, S. C. Wilks, and P. A. Amendt

This series of shots was designed to measure the electric fields produced by strong shocks in single- and multi-species plasmas. Strong shocks ($M > 1.6$) produce electric fields as a result of streaming of the electron precursor ahead of the ion shock, which is not accurately reproduced in hydrodynamic models. This experiment is intended to quantify electric-field strength and position in a shock platform previously characterized in FY16 (Ref. 35) to be used to constrain high-fidelity physics codes.

The KineticDynamics-17A Campaign performed proton radiography of a shock-tube target platform developed in the KineticShock-16 series. Ten beams with 2.5 kJ in a 0.6-ns impulse drove a 2-μm-thick SiO₂ ablator attached to one end of the tube, launching a strong shock into the 1 atm of gas contained in the tube. A D³He-filled backlighter capsule was imploded to produce 3.0- and 14.7-MeV protons, which transited perpendicular to the shock front and were recorded using CR-39. Deflections of the protons caused by electric fields were recorded as changes in the proton fluence with position on the detector, as shown in Fig. 152.89. Gas fills of H₂ (two shots), Ne (three shots), and H₂ + 2% Ne (six shots) were probed while varying the backlighter timing relative to the shock.

Radiographs were collected on all shots, demonstrating electric-field formation at both the shock front and the ablator/gas interface. X-ray framing-camera images recorded perpendicular to the radiography axis confirm that the flow velocity is consistent with previous experiments (450 μm/ns) and identify

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**Figure 152.89**

(a) Experimental configuration, (b) 3-MeV proton radiograph, and (c) lineout of proton fluence for shot 84107.
the position of the SiO$_2$ ablator at the radiography sample time on each shot. Analysis is ongoing and results were presented at the APS DPP Meeting in October 2017.

**Measuring Thermal Conductivity of D$_2$**
Principal Investigators: S. Jiang and Y. Ping
Co-investigators: P. M. Celliers and O. L. Landen

This campaign successfully developed a platform of refraction-enhanced x-ray radiography for planar cryogenic targets and obtained usable radiographs. The campaign fielded two half-days on OMEGA during FY17. The series of cryogenic shots was designed to measure the thermal conductivity of liquid D$_2$ at 1 to 10 eV, which helps to benchmark different transport models used in ICF ignition target design. Figure 152.90 shows the experimental configuration. The CH and liquid D$_2$ were heated to different temperatures by the x rays generated from heating lasers incident on a thin Zn foil. This builds up a density gradient because of thermal conduction. The evolution of the CH/D$_2$ interface was measured using refraction-enhanced x-ray radiography with high spatial resolution. The temperature can be constrained by the wave velocities in CH/D$_2$. A vanadium foil backlighter was used to generate the probing x rays. The backlight x-ray images were collected with a four-strip framing camera. Useful radiographs at different delays were recorded (shown in Fig. 152.90), including undriven and driven data at 0.5 ns, 1.5 ns, 2 ns, and 2.5 ns. Analysis of the measured data is in progress.

**Characterization of Laser-Driven Magnetic Fields**
Principal Investigator: B. B. Pollock
Co-investigators: C. Goyon, G. J. Williams, D. Mariscal, G. F. Swadling, J. S. Ross, S. Fujioka, H. Morita, and J. D. Moody

BFieldLoop-17A and -17B continued the laser-driven magnetic-field experimental campaign on OMEGA EP. The goal for 17A was to study the impact on the magnetic-field–generation process of modifying the target geometry, while the primary

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**Figure 152.90**
(a) The experimental configuration, target geometry, and (b) raw radiographs from an x-ray framing camera (XRFC). The delays of both the framing camera and the x-ray backlighter (BKL), with respect to the heating pulses, are shown in the table next to each radiograph.
objective for 17B was the extension of the probing time for the field. Figure 152.91 shows the general experimental geometry. The field is produced by currents flowing in a U-shaped gold foil target. One or two long-pulse beams are directed through holes in the front side of the target to produce a plasma on the interior rear side of the target. Plasma produced in this region expands toward the front side, setting up a voltage across the target that drives the current. The fields are measured by proton deflectometry, where the protons are produced through target normal sheath acceleration by the OMEGA EP backlighter beam incident on a separate, thin Au foil.

![Experimental configuration for BFieldLoop experiments. RCF: radiochromic film.](image)

Previous experiments showed that for a 1-TW long-pulse laser drive, fields of ~200 T can be produced in the interior of the loop portion of the target. The 17A experiment modified the front of the target by removing the bottom portion where the laser holes are located. This target modification showed a difference in the resulting proton images recorded on the radiochromic film (RCF), and modeling is ongoing to quantify the modification to the field strength and topology.

For 17B, a gold shield was added between the main target and the backlighter target to protect the rear surface of the backlighter and allow for late-time probing of the magnetic fields. Prior to 17B the ability to measure magnetic fields after the long-pulse beams turned off had been limited. By using a 750-ps long-pulse B-field drive beam, the addition of this shield allowed the probe time to be extended from 750 ps (the end of the drive beam) to 16 ns. Preliminary analysis of these results indicates that the magnetic field persists long after the drive laser turns off, decaying exponentially with a time constant of ~2 ns. This is consistent with a circuit theory treatment of the interaction, where the drive laser acts as a voltage source while turned on; after the voltage is turned off, the current (and consequently the magnetic field) decays with the \( \frac{L}{R} \) time constant of the target. For the geometry of the 17B targets, this \( \frac{L}{R} \) time is ~2.6 ns.

**Study of Shock Fronts in Low-Density Single- and Multi-Species Systems**

Principal Investigator: R. Hua (UCSD)
Co-investigators: Y. Ping, S. C. Wilks, R. F. Heeter, and J. A. Emig (LLNL); H. Sio (MIT); C. McGuffey, M. Bailly-Grandvaux, and F. N. Beg (UCSD); and G. W. Collins (LLNL, LLE)

This series of shots is designed to study the shock-front structure of low-density single- and multi-species systems in a planar geometry, using proton radiography from a broadband TNSA source and x-ray emission spectroscopy. Unlike most other low-density shock experiments that are mostly carried out in convergent geometry with near-monoenergetic \( ^3\text{He} \) proton radiography, the planar geometry of these experiments enables one to distinguish the shock front from its pusher, and the TNSA protons provide broadband energy measurements.

In this platform, the strong shock of interest is initiated by launching three long-pulse beams from OMEGA EP onto a \( 2\mu \text{m} \) SiO\(_2\) foil (Fig. 152.92). By ablation, the foil pushes into a cylindrical kapton tube that is prefilled with either pure helium or a helium/neon mixture. A short-pulse beam is fired a few nanoseconds after the shock-driving beams, onto a TNSA proton target normal to the axis of the kapton tube. The protons are driven through two windows on the tube into a stack of RCF’s on the other side. Photons emitted from the shock front through another window are recorded by a 1-D spatially resolved soft x-ray spectrometer [called variable-spaced grating (VSG)]. The details of the design and its feasibility have been published.\(^7\)

With pure helium fill, a proton accumulation layer at the shock front, resulting from a self-generated electric field, has been recorded on the RCF. A discussion of the field information inferred from the proton behavior was published recently.\(^{12}\) Shocks in the mixed gas displayed different features. Not one, but two proton accumulation layers are observed. In addition, line emission from the added neon gas is recorded by the VSG, which provides the potential to constrain the temperature and density of the shock front. Further analysis of these results is in progress. Results of these campaigns are being used to further optimize the platform to make more measurements in FY18.
High-Energy-Density Experiments
1. Material Equation of State and Strength Measured Using Diffraction

Measurement of Pb Melt Curve Along the Hugoniot Using In-Situ Diffraction
Principal Investigator: C. E. Wehrenberg
Co-investigators: R. Kraus, F. Coppari, J. M. McNaney, and J. H. Eggert

This series of shots was designed to measure the melt curve of Pb on the Hugoniot using in-situ x-ray diffraction on OMEGA EP’s PXRDIP diagnostic. The physics package consists of a 40-μm epoxy ablator, a 12.5-μm lead foil, and a 100-μm LiF window for VISAR. A single UV beam is used to drive a steady shock through the ablator and into the Pb sample. A separate beam drives a Ge x-ray source target, and the diffracted x rays are collected on image plates mounted on the inside of the PXRDIP box. The solid or liquid state of the shock-compressed lead is determined by the presence of a diffraction signal as either sharp lines (solid) or a single diffuse line (liquid). Figure 152.93 shows a diffraction pattern containing both the diffuse liquid signal from lead under 47 GPa of shock compression and sharp diffraction lines generated by the Pt pinhole.

This campaign is performed in conjunction with similar campaigns on the NIF and the Linac Coherent Light Source.
(LCLS). The epoxy ablator setup, which eliminates reverberations between the ablator and any glue layers, and the Ge x-ray source target were developed on the NIF campaign before being applied on these shots. On this OMEGA EP day, alternating beams to increase the shot rate achieved 11 shots that provided a large data set to map the melt curve on the Hugoniot. This data will be compared with similar experiments performed on LCLS, making it possible to compare shots done with greater x-ray diffraction precision (because of the x-ray free-electron laser) with the greater-precision shock drive available on OMEGA EP. In addition, these melt curve measurements will be used in the design of future NIF shots using shock-ramp combination drives to measure melt curves in off-Hugoniot states.

**Development of a New Platform for Measuring Recrystallization**
Principal Investigator: F. Coppa
Designer: R. Kraus
Co-investigators: C. E. Wehrenberg and J. H. Eggert

This campaign seeks to develop a platform for measuring recrystallization of Pb through shock-ramp compression. By first launching an initial shock to compress the sample along the Hugoniot close to the melting pressure, letting the sample release into the liquid phase, and then recompressing it with ramp compression across the solid–liquid phase boundary, one can measure high-pressure melting lines of materials. The structure of the Pb upon shock, release, and ramp compression is monitored by x-ray diffraction. The onset of melting is identified by the appearance of a diffuse scattering pattern and the disappearance of the Bragg diffraction lines characteristic of the solid. The pressure is monitored by VISAR looking at the interface between the Pb and a LiF window.

Building upon the successful use of Be ablators initiated in FY16, the FY17 series used a slightly modified target design that included a Ge preheat shield (Fig. 152.94). Hydrocode simulations allowed us to design a suitable laser pulse shape to compress the Pb sample along this complicated shock–release–ramp path.

![Example diffraction image showing a diffuse signal generated by liquid Pb under shock compression at 47 GPa and sharp diffraction lines generated by the Pt pinhole.](image1)

![Target schematic and example of a diffraction image showing diffuse (liquid) and sharp (solid) features from the Pb sample. VISAR: velocity interferometer for any reflector; XRD: x-ray diffraction.](image2)
Diffraction data indicate that Pb completely melts above 50 GPa along the Hugoniot. Resolidification into a structure consistent with the bcc lattice has been observed upon subsequent ramp compression to higher pressures. Data analysis is ongoing and will provide valuable information for future OMEGA and NIF campaigns.

**High-Pressure Crystal Structure of Ramp-Compressed Pb and Ta to Resolve Known or Potential High-Pressure Phase Transformations**

Principal Investigator: A. E. Lazicki  

Several previous experiments on OMEGA, OMEGA EP, and the NIF in the past five years have examined the crystal structure of Pb and Ta at high pressure. Pb is of specific interest because its high-pressure/high-temperature phase diagram, including two phase transformations, is well constrained between 0 to 100 GPa from diamond-anvil cell experiments. Therefore, it provides an opportunity to test for the effects of rapid compression and heating, inherent from ramped laser drives, on phase boundaries. Although very high quality data have already been collected in the 100- to 500-GPa pressure regime, previous measurements have failed to reach sufficiently low pressures to cross the two known phase boundaries. Meanwhile, measurements of Ta structure have suggested a possible new high-pressure phase, but contamination of the diffraction pattern by diamond, a material that has always been present in the target packages, has made the determination ambiguous. The goals of this OMEGA EP campaign were to capture the low-pressure phase transformations in Pb and, in the case of Ta, to eliminate all diamond from the target and examine the pressure regime where previously a new phase was suggested. Both target packages used beryllium ablators and LiF windows, with thin (3- to 5-µm) layers of Pb. The shots used a combination of Cu and Ge backlighter sources.

Very high quality diffraction data were collected on all shots and the low-pressure phases of Pb were clearly observed and followed with pressure, constraining the phase boundaries (Fig. 152.95). The diagnostic damage incurred from the explosive ablation of beryllium made it impossible to reach the desired pressure regime for the Ta measurements, but the measurements provided at least a clear lower bound to the transition. These results were presented at an invited talk at the 55th meeting of the European High-Pressure Research Group and are being prepared for publication.

![Figure 152.95](U2274JR)

(a) Raw data and (b) integrated diffraction patterns from three phases of Pb. A Cu Heα x-ray source was used.
Development of an In-Situ Pressure Standard for Diffraction Experiments
Principal Investigator: F. Coppari
Co-investigators: J. H. Eggert and R. Kraus

The goal of this campaign is to develop a new way of determining pressure in diffraction experiments, based on the use of an in-situ pressure gauge. By measuring the diffraction signal of a standard material (whose equation of state is known) compressed together with the sample, one can determine the pressure reached during ramp compression.

Currently, pressure is determined from VISAR measurements of diamond free-surface velocity or particle velocity through a transparent window (such as LiF or MgO). This method is, in some cases, ambiguous (lack of reflectivity or shock formation) or relies on assumptions and equation-of-state models. Cross-checking the VISAR measurement with in-situ pressure determination using the diffraction signal of a standard material will improve the diffraction platform by providing a complementary way of determining the pressure state within the sample, with great impact to the programmatic effort of determining structures and phase transitions at high pressure and temperature. In addition, combining pressure determination from VISAR and from the in-situ gauge can provide information about the temperature of the sample by measuring the calibrant thermal expansion.

Building upon the results obtained in FY16, this FY17 campaign focused on the study of Au, Ta, Pt, and W pressure standards ramp compressed to 2 Mbar on OMEGA. Data were collected on eight successful shots in a half-day, including diffraction patterns of the Au/Ta and Pt/W pairs. All shots provided useful data (Fig. 152.96). Comparison of the pressures obtained from VISAR analysis and from the diffraction patterns will yield information on the accuracy of the VISAR method as well as on the existence of preheating.

Future directions of this work will look at characterizing the pressure standards at higher pressure and implementing this technique into other diffraction experiments. This platform still requires additional development before it can be used routinely in diffraction experiments, but the data collected so far are extremely encouraging and suggest that the use of an in-situ pressure gauge can be a viable path forward in future x-ray diffraction (XRD) measurements on both OMEGA and the NIF.

Development of Simultaneous Diffraction and EXAFS Measurements
Principal Investigator: F. Coppari
Co-investigators: Y. Ping and J. H. Eggert

Being able to measure simultaneous diffraction and an extended x-ray absorption fine-structure (EXAFS) signal in the same shot will be an enormous advance for laser-based materials experiments, providing a simultaneous probe of both the long-range [x-ray diffraction (XRD)] and short-range (EXAFS) order of the material, as well as two complementary probes of the Debye–Waller factor to gain information about the temperature state of the material under investigation. The approach successfully developed in FY16 was to use the PXRDIP diagnostic to measure diffraction and the x-ray source spectrometer to measure EXAFS. The challenge was to find a single suitable backlighter that would generate both a monochromatic (for diffraction) and broadband (for EXAFS) x-ray source. Success was achieved in measuring simultaneous
XRD and EXAFS of Fe by using a dual-foil bremsstrahlung backlighter, where one foil is optimized to generate He\textsubscript{α} radiation for diffraction and the other to generate a continuous and broadband x-ray source for EXAFS (see Fig. 152.97).

The goal for the FY17 campaign was to improve the quality of the EXAFS measurement by using a higher-Z foil to increase the bremsstrahlung radiation, thereby generating a brighter x-ray source to allow EXAFS measurements at a higher energy (Zr K edge at 18 keV). Both Au and Ta were tested as x-ray sources for EXAFS and Fe or Cu as an x-ray source for diffraction. The spectrum generated by Au is indeed brighter, but it was not yet possible to see the EXAFS modulation above the Zr K edge. Poor spectral resolution certainly played a role in the deterioration of the EXAFS data: in this experimental setup, the spectral resolution is limited by the x-ray source size, which is enlarged by the expanding plasma generated by direct laser ablation of the foil. Further improvement could be obtained by using a spectrometer with focusing geometry that would be less sensitive to the effects of the source size.

Texture Diffraction and Recovery of Shock-Compressed Samples for Ex-Situ Study
Principal Investigator: A. Krygier
Co-investigators: C. E. Wehrenberg, H.-S. Park, and D. C. Swift

This shot day continued the previous work of Wehrenberg on highly textured Ta that studied the deformation response of Ta to shock waves. The goal of these shots is to study if effects seen on the initial shock are present during a second shock. In particular, twinning, which is not included in the Livermore strength model for Ta, was observed to play a significant role in deformation over a wide range of shock pressures (30 to 150 GPa). X-ray scattering was measured using the PXRDIP platform in combination with VISAR on OMEGA EP.

These shots included a time series of x-ray diffraction measurements of various pressure histories. In the first configuration, the sample was shocked to the twinning regime, released, and then shocked above twinning. Twinning was observed, as shown by the red circle in Fig. 152.98, in the initial shock but not on release or on the second shock. In the second configuration, the sample was shocked to the twin-
ning threshold, released, and then shocked into the twinning regime. Twinning was not observed on either the release or the second shock, suggesting that dislocations generated by a shock wave moderate the deformation response. In addition, a shock-ramp history was performed, using the initial shock in the twinning regime that is approximately analogous to strength experiments in Ta performed on the NIF. In this regime, a new texture was observed that cannot be explained by twinning. Analysis is ongoing.

2. Material Equation of State Using Other Techniques

**Development of a Platform for Equation-of-State Measurements Using Flyer Plate Impact**
Principal Investigator: F. Coppari

This ongoing campaign is developing a platform to use OMEGA to accelerate flyer plates to a high velocity for absolute equation-of-state (EOS) measurements by symmetric impacts. The concept is to ramp compress a metallic foil through indirect laser ablation across a vacuum gap and observe the flyer impact on a same-material sample, mounted side by side with a transparent LiF window (Fig. 152.99). By measuring the flyer-plate velocity through the transparent window prior to impact, and the resulting shock velocity in the metallic sample using transit time measurements, the principal Hugoniot of the metallic foil can be determined absolutely (e.g., without needing a known pressure reference), enabling one to develop an EOS standard.

Specifically, this campaign tested three flyer-plate materials to check performance and hydrodynamic prediction capabilities. Plates of Mo, Cu, and W were chosen because they can be “easily” ramp compressed to a high pressure and do not exhibit structural solid-phase transitions. Building upon the previous FY16 campaign, a successful half-day of eight shots was fielded in FY17, accelerating flyer samples to different velocities. While Cu and W showed anomalous behavior, such as flyer breakup resulting in the loss of VISAR reflectivity before impact, Mo was successfully accelerated to 14 km/s, corresponding to a shock pressure into the Mo sample of ~1 TPa. As Fig. 152.99 shows for the lower-pressure shot, the velocity ramps up smoothly and VISAR data are obtained up to impact. The impact actually shocked the LiF window into a pressure range where it was no longer transparent, impairing the VISAR signal. Future experiments will look at accelerating Mo flyer plates to different pressures to characterize its Hugoniot EOS in a wide pressure range, using quartz as a transparent window.

**Development of Simultaneous EXAFS and VISAR Measurements**
Principal Investigator: Y. Ping
Co-investigators: F. Coppari and J. H. Eggert

This campaign aims to test mirror shielding for simultaneous VISAR measurements in the presence of the implosion backlighter needed for EXAFS measurements. Both the VISAR and EXAFS measurements must pass through the sample, so an x-ray transparent optical mirror is used to redirect the VISAR probe beam onto the sample. This mirror is, however, vulnerable to blanking. For these FY17 shots a tilted mirror design was implemented to move the mirror out of the line of sight of the x-ray transmission path, so that more shielding of the mirror could be applied. This design extended the mirror life time by about 1 ns, but after that, the mirror was still blanked and the VISAR signal lost. On the other hand, good EXAFS data were collected for the Co K edge for the first time, as shown in Fig. 152.100. Separate shots for VISAR without the implosion also produced good VISAR data to characterize the pressure. The similar experimental design on the NIF has
the mirror positioned behind the target, which survived and produced a VISAR signal in the presence of the implosion in a recent NIF shot.

**Development of a Platform for EXAFS Measurements at the L Absorption Edge of High-Z Materials**

Principal Investigator: F. Coppari  
Co-investigators: Y. Ping and J. H. Eggert

EXAFS measurements in the 7-keV x-ray region are routinely performed on the OMEGA laser during shock and ramp compression. However, because the brightness of the capsule implosion used as an x-ray source decreases at high energy, measuring the EXAFS of materials with an absorption edge higher than 7 keV is very difficult in a single shot. In addition, as one seeks to study higher-Z materials and observe the L absorption edge, the amplitude of the EXAFS signal decreases as well because the cross section for the absorption at the L edge is lower than at the K edge. These issues make these measurements extremely challenging.

In prior years the TaXAFS Campaigns looked at the Ta L3 edge (~10 keV) and were able to obtain a good signal-to-noise ratio by averaging 15 shots or by using a multicrystal spectrometer. The goal of the FY17 campaign is to look at materials whose L edge is closer to 7 keV, where the number of photons generated by the capsule implosion is higher, to see if a good signal can be obtained in a single shot. The shots studied the Ce L3 edge (5.7 keV), but although a nice contrast at the Ce L edge has been measured, no clear EXAFS modulations could be observed in a single shot (Fig. 152.101), probably also a result of limited spectral resolution. This suggests that in order to successfully measure L-edge EXAFS, a spectrometer with focusing geometry is needed to increase the signal because of the higher collection efficiency and reduce the sensitivity to x-ray source size broadening, therefore improving the quality of the data.
Development of Spherically Convergent Equation-of-State Measurements
Principal Investigator: A. E. Lazicki

This campaign is developing a platform for measuring Hugoniot EOS at pressures much higher than can be achieved using a standard planar drive. This platform is intended to collect data in the pressure regime of 100+ Mbar, where currently very little data exist for any material, for the purpose of constraining EOS models.

The FY17 campaigns first used a hohlraum (indirect drive) to launch converging shock waves into solid spheres of CD (deuterated) and CH (normal) plastic. Along the axis of the hohlraum, vanadium He_x backlight 2-D x-ray images of the imploding sphere were collected with a framing camera. On some shots, x-ray Thomson-scattering measurements were also made using a Zn backlighter and a spectrometer at the hohlraum equator [Figs. 152.102(a) and 152.102(b)]. The radiographs yield density and shock velocity that make it possible to calculate the shock state using the Rankine–Hugoniot equations, and the scattering data yield information about temperature and ionization state.

The FY17 shots improved on prior measurements by increasing hohlraum gas fill to eliminate suspected hohlraum blowoff features and using faster-gating cameras to improve spatial resolution. Neutron diagnostics were also fielded to detect neutrons from the hot spot. The new design for the x-ray scattering measurement yielded a high-quality Compton feature and an elastic feature potentially artificially elevated because of Zn plasma leakage. These results, together with the FY16 measurements on CH_2, are being summarized for a publication that will describe the principal Hugoniot of plastic from the initial densities of CH, CH_2, and CD.

In addition, a separate half-day of shots continued development of a platform to achieve hundreds-of-Mbar pressures in a spherically converging shock wave, launched by using direct laser ablation of the sphere. This measurement probed deuterated plastic (CD) using radiography, x-ray Thomson scattering, and neutron yield. Data improved in quality compared to FY16 but indicated some drive asymmetry and preheating effects, requiring further design optimization.

Development of a Conically Convergent Platform for Hugoniot Equation-of-State Measurements in the 100-Mbar to 1-Gbar Pressure Regime
Principal Investigator: A. E. Lazicki

This campaign was designed to develop a platform for measuring Hugoniot EOS of arbitrary (including high-Z) materials at pressures much higher than can be achieved using a standard planar drive. This platform is intended to collect data in the pressure regime of 100+ Mbar, where currently very little data exist for any material, for the purpose of constraining EOS models.

To achieve the desired pressure amplification, converging shock waves are launched into a cone inset in a halfraum. For appropriate cone angles, nonlinear reflections of the shock wave result in the formation of a Mach stem: a planar high-pressure shock that propagates along the axis of the cone. This approach was tested in FY16 and produced promising results but suggested the need for preheat shielding. The FY17 half-day campaign attempted to overcome this need with Au preheat shield layers by experimenting with a porous cone material, from which simulations suggested that increased pressure amplification would be possible. Additionally, these shots tested multiple cone angles and fielded targets with quartz.
windows and Al steps to attempt to quantify the Mach wave’s strength and steadiness. Two results are shown in Fig. 152.103. Transit time calculations indicate that 100+ Mbar shock waves were generated in the CH cones. At peak laser intensities the quartz windows blanked, indicating that the Au preheat shield thicknesses were not sufficient, and the profile of the shock breakout suggested the Mach wave was decaying (unsupported) through the full target thickness. A subset of the shots also attempted to use area-backlit radiography to image the Mach wave formation in the cones themselves, through slits in the halfraums. Geometric issues made it very difficult to interpret the framing-camera images. However, the lessons learned on this shot day were very important in designing a successful test shot on the NIF, completed on 31 August 2017.

![Experimental configuration and raw SOP data showing the breakout time from (a) a stepped Al sample with Au preheat shield layers and (b) a thick quartz sample with a Au preheat shield layer.](image)

Figure 152.103

3. Material Dynamics and Strength

**Copper Rayleigh–Taylor (CuRT) Growth Measurements**

Principal Investigator: J. M. McNaney

Co-investigators: S. Prisbrey, H.-S. Park, C. M. Huntington, and C. E. Wehrenberg

The CuRT Campaign is part of the material strength effort, which is aimed at assessing the strength of various metals at high pressure and high strain rate. The goal of the CuRT platform is to measure Rayleigh–Taylor (RT) growth of samples that behave “classically,” i.e., can be fully modeled using a fluid description. In this series of experiments the intent is to measure RT growth in liquid copper at high pressures. A second goal is to demonstrate the dynamic range of the technique by measuring RT growth in solid copper.

Without the stabilization of strength, classical RT growth is characterized by a growth rate $\gamma = \sqrt{kgA_n}$, where $k$ is the wavelength of the unstable mode, $g$ is the acceleration, and the Atwood number $A_n$ quantifies the magnitude of the density jump at the interface. Acceleration of the sample in the experiment is provided by the stagnation of a releasing shocked plastic “reservoir,” which is directly driven by –2 to 8 kJ of laser energy, depending on the desired material condition. The growth of preimposed ripples is recorded using transmission x-ray radiography of a copper He$_\alpha$ slit source, where the opacity of the sample is calibrated to the ripple amplitude. The pre-shot metrology and measured $\rho$ of the driven sample together yield the growth factor, which is compared to models of RT growth. Diagnostic features such as a gold knife edge on the sample allow one to measure the modulation transfer function and create an opacity look-up table on each shot, resulting in error bars of approximately ±10%.

The March 2017 campaign produced the first results for liquid copper RT. Analysis of the velocimetry (Fig. 152.104) indicated that the copper RT samples were subjected to a shock of $\approx 5$ Mbar, leading to a complete melt of the sample and subsequent RT growth in the liquid phase. The liquid Cu RT growth curve is presented in Fig. 152.105. A second day of liquid copper shots took place on 13 September 2017. The results of those experiments (also in Fig. 152.105) were consistent with the values obtained in the March data. Simulations of the liquid Cu growth are in progress.

**Evaluation of Additively Manufactured Foams for Ramp-Compression Experiments**

Principal Investigator: R. F. Smith

The FY17 “AMFoam” Campaigns continued to evaluate the use of 3-D–printed or additively manufactured foams as surrogates to carbonized resorcinol foams (CRF’s) in ramp-compression target designs. The 3-D–printed foams may be characterized as follows: The $100 \times 100 \times 16\mu m^3$ log pile blocks, composed of individually printed lines, are stitched together to form $1.7\text{mm}$-diam layers. Seven $16\mu m$ layers are then stacked on top of one another to arrive at cylindrical AM foams that are $112 \mu m$ tall (Ref. 37). These foams are glued onto a $25 \mu m$ Br $+ 120 \mu m$ 12% Br/CH assembly.
4. National Security Applications

**SolarCellESD: Solar Cell Electrostatic Discharge Experiments**
Principal Investigator: K. Widmann
Co-investigators: P. Jenkins (NRL); S. Seiler (DTRA); and P. Poole and B. Blue (LLNL)

The overall goal of the SolarCellESD Campaign is to determine experimentally whether prompt x rays can induce failure modes in solar arrays that are not accounted for by simply testing the individual solar cells alone. The solar-cell array is fielded as part of the x-ray Langmuir probe detector (XLPD) cassettes and exposed to x rays from a laser-driven source. The FY17 campaigns added a partial electromagnetic-interference enclosure of the XLPD front end—“partial” because the enclosure has a large rectangular opening providing an unobstructed view of the x-ray source for the solar-cell array. The bias and diagnostic electronics for the solar cells were also improved such that the bias circuit is fully isolated from any of the target chamber components and the bias voltage can be changed manually between shots. This bias voltage allows one to mimic the voltage difference between two adjacent solar cells from different strings in large arrays, which can range from tens to a few hundred volts. Figure 152.107 shows a schematic and a photo of the new and improved XLPD front end.

The FY17 campaign continued using the x-ray source developed in FY15: a small gold halfraum (600-μm diameter, 600-μm long) with a small pinhole (60-μm or 100-μm diameter) at the “closed” side of the halfraum. The targets were driven with three sets of 1-nsec laser pulses to generate x-ray pulses of 3-ns duration. These targets provide two significantly different

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**Figure 152.104**
(a) The active shock breakout (ASBO) data from shot 26423 (September 2017); (b) analyzed results and comparison to March 2017 results.

**Figure 152.105**
Growth factor data for liquid copper from the March and September shot days.

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[Fig. 152.106(a)]. Following the ramp-compression design described in Ref. 30, 15 OMEGA beams with 300 J in 2 ns result in a ramp-compression wave being launched into a stepped-Al/LiF sample [Fig. 152.106(a)]. At a controlled time after this compression begins, the OHRV (2-D VISAR probe) takes a 2-D snapshot of the reflectivity and velocity field with a spatial resolution of ~3 μm (Ref. 38). An example of the intensity field recorded on the 2-D VISAR is shown in Fig. 152.106(c).

These campaigns have varied the structure of the 3-D-printed foam with the goal of optimizing the temporal ramp profile. These experiments have been conducted in support of material strength experiments on the NIF.
Figure 152.106
(a) Schematic diagram describing the experimental setup for the AMFoam-17A Campaign, which combined the ASBO and OHRV (2-D VISAR) diagnostics on a single shot. The goal of these shots was to characterize the temporal and spatial drive associated with additively manufactured (3-D-printed) foam targets. The line VISAR shows the 1-D velocity measure as a function of time. Two-dimensional VISAR provides the 2-D velocity measurement of 3-D–printed foams at a snapshot in time. (b) The ASBO record shows a characteristic ramp/pressure-hold/ramp compression profile. (c) Sample reflectivity at OHRV probe time. For the AMFoam-17B Campaigns, only the ASBO (VISAR) diagnostic was used. Preliminary 2-D VISAR non-fringes image and velocity analysis shows AMFoam exhibits a spatial structure comparable or perhaps smoother than CRF foam.

Figure 152.107
(a) Front end of the x-ray Langmuir probe detector (XLPD) with the solar cell (mini) array, i.e., 2 x 1 cells and (b) photo of the XLPD (side-on view) with the new electromagnetic-interference enclosure and the various bias and diagnostic electronics.
x-ray flux (and fluence) levels depending on the target orientation with respect to the solar cell array (“open side” versus “pinhole side”). Despite the large difference in the measured radiant x-ray power, from 300-GW/sr peak flux for the open side down to a few tens of GW/sr for the pinhole side, the spectral intensity distribution should be very similar. Unfortunately, the low signal level on the Dante x-ray diodes for the low-flux case did not allow us to make a good quantitative comparison of the unfolded spectral intensities. Figure 152.108 shows the x-ray flux measured with Dante for two open-side shots.

With the improved hardware, it was possible to generate and observe, for the first time, x-ray–induced arcing just between the solar cells without creating discharge between the solar cells and any of the other components in the target chamber. These shots also demonstrated the feasibility of mitigating the x-ray–induced arcing by altering the bias voltage for the solar-cell array. The next goal for the this campaign is a detailed parameter scan to determine the threshold values for the x-ray fluence and associated bias voltage levels, respectively, for x-ray–induced arcing in solar-cell arrays.

**Plasma Instability Control to Generate a High-Energy Bremsstrahlung X-Ray Source**

Principal Investigator: P. L. Poole

In FY17 a campaign began to develop a high-fluence x-ray source in the 30- to 100-keV range by deliberately stimulating and optimizing plasma instabilities. High-fluence sources at lower energies are currently used for materials effects studies in extreme environments, but stronger sources are needed for >30-keV x rays. This project aims to enhance laser conversion to plasma instabilities such as stimulated Raman scattering (SRS), which accelerate electrons in plasma waves that will convert to high-energy x rays via bremsstrahlung in the high-Z hohlraum wall.

The SRS-Xray-17A half-day centered on optimizing plasma conditions within a hohlraum for high SRS gain by reversing ideas on how to mitigate these effects for fusion conditions (SSD bandwidth off, phase plate changes, etc.). Two types of targets were fielded to test SRS generation: one with 6-mg/cm³ (0.2 n_{e}) SiO_{2} foam fill, the other with 10 μm of parylene-N (CH) inner lining. The CH-lined hohlraums were illuminated with an 80-ps picket prepulse arriving a variable time before the main 1-ns, 450-J/beam laser drive.

Primary results were promising: the best targets exhibited 10× more x-ray yield in the >50-keV regime than typically seen on direct-drive shots. The SRS streak data are being compared to the various x-ray diagnostics fielded to obtain a detailed picture of instability gain and directionality and its impact on x-ray fluence. While the CH prepulsed targets exhibited factors-of-6 more direct SRS backscatter than the no-prepulse CH and foam targets, these latter configurations had the best x-ray yield (Fig. 152.109). The increased indirect backscatter on the no-prepulse CH target provides an avenue of further simulation and future experimental study to investigate the directionality of SRS backscatter within the hohlraum and its impact on the ultimate x-ray yield. Target design simulations are underway, using these valuable experimental results to plan the FY18 OMEGA and NIF campaigns. These promising initial results are a valuable stepping stone toward a new x-ray source that will represent a large capability increase for national security applications and related materials under extreme conditions studies, with the additional benefit of broadening the understanding of plasma instability control for fusion and other applications.

5. Plasma Properties

**Investigation of Orthogonal Plasma Flows in the Presence of Background Magnetic Fields**

Principal Investigator: B. B. Pollock
Co-investigators: T. Johnson, G. F. Swadling, J.S. Ross, and H.-S. Park

The DebrisPlasma-17A shots continued the Magnetized Collisionless Shock–Weapons Effect Campaign from previous
The goal for this series of experiments is to quantify the interaction of orthogonal plasma flows, with and without a background magnetic field. The field is supplied by LLE’s magneto-inertial fusion electrical discharge system (MIFEDS), which delivers 4 to 8 T at the interaction region of the experiments, depending upon the specific geometry of the MIFEDS coils (Fig. 152.110). The orthogonal plasma flows originate from two separate foil targets, one of which is mounted to the MIFEDS structure. The foil material composition, laser drive, spacing, and time of flight to the interaction region can be varied on each shot.

The interaction region of the two flows is simultaneously probed with 2\omega Thomson scattering and protons from the implosion of a D\textsuperscript3He-filled capsule. Initial measurements of the electron density and temperature from the Thomson scattering do not indicate a strong dependence on the strength or direction of the background magnetic field. The proton deflectometry data do show structural differences in the measured proton distribution with and without the field, but further modeling and simulations are needed to quantify these differences.

This campaign was designed to measure the thermal transport through gold layers as well the emitted M-shell gold spectra from a well-characterized and uniform plasma for comparison to atomic kinetic models. The buried layer target geometry used for this experiment is capable of generating plasmas with an electron temperature of \sim{}\textasciitilde 2 \text{ keV} at electron densities of \sim{}10\textsuperscript{21} electrons per cubic centimeter. These are within the range of conditions found inside gold hohlraums used on experiments on the NIF, providing a stable laboratory setting for radiation transport and atomic kinetic studies of hohlraum plasmas.

Planar, buried layer targets composed of Ti, Mn, and Au were illuminated evenly on both sides (Fig. 152.111) to heat the sample. The sample was buried between two 5-\textmu{}m-thick layers of Be serving as an inertial tamp to slow the expansion of the sample. Time-resolved 2-D images of the target’s x-ray emission, viewed both face-on and side-on, were recorded using time-gated pinhole cameras. The K-shell spectra from the Ti and Mn were used to determine the electron temperature of the plasma. The time-resolved spectra were recorded using a
crystal spectrometer coupled to a framing camera, as well as a crystal spectrometer coupled to an x-ray streak camera. Two additional time-resolved crystal spectrometers were used to record the full range of the Au M-shell emission. All of the framing cameras used, for imaging as well as spectroscopy, were co-timed so the plasma conditions at the time of the measured Au M-shell emission could be established from synchronous K-shell and imaging data.

During the campaign two different sample thicknesses were used to measure the thermal transport through Au. Two different pulse shapes were also used to assess which was most efficient for coupling laser energy into the buried layer target. A complete set of data from all six precisely co-timed diagnostics was recorded for both target types, using both pulse shapes during this August campaign, at temperatures ~2 keV. Data analysis is underway.

6. Hydrodynamics

**Experiments on the Rayleigh–Taylor Instability in the Highly Nonlinear Regime**

Principal Investigators: C. C. Kuranz (University of Michigan) and C. M. Huntington (LLNL)

Co-investigators: L. Elgin, G. Malamud, S. R. Klein, R. P. Drake, and D. Shvarts (University of Michigan) and T. Handy and M. R. Trantham (LLNL)

These experiments observe the evolution of the single-mode Rayleigh–Taylor instability (RTI) in low- and high-Atwood-number regimes at late scaled times (Fig. 152.112). Models predict two growth phases of the RTI: exponential growth, followed by a nonlinear stage reaching a terminal velocity. For low-Atwood number systems, numerical simulations show an additional growth phase in the late nonlinear stage, characterized by reacceleration. There are, however, claims that this reacceleration may be an artifact of the simulations and may not reflect the evolution of classical RTI. Prior experimental studies of RTI growth have not created the conditions necessary to observe the late nonlinear stage, which requires large aspect ratios of the spike and bubble amplitudes to the perturbation wavelength (1 \( \leq h_{sb} / \lambda \leq 3 \)) (Ref. 40).

The first two experiments in this new campaign were conducted in FY17. A laser-driven blast wave accelerates an RT-unstable interface in a shock tube. X-ray radiographs along dual orthogonal axes capture the evolution of RTI. Late scaled times are achieved with small-wavelength (\( \lambda = 40-\mu m \)) seed perturbations at the interface. PAI (polyamide-imide) plastic (1.4 g/cm\(^3\)) is used as the heavy fluid. The lighter fluid consists of CRF, with pre-shock densities of 0.05 g/cm\(^3\) (high Atwood) or 0.4 g/cm\(^3\) (low Atwood). The first shot day demonstrated x-ray–backlit imaging capable of resolving...
the small-wavelength RT spikes and bubbles. But the plastic shock tube could not support the higher internal pressure of the dense, low-Atwood targets. A new target design was developed and fielded for the second shot day. Improvements included Be walls, which can be thicker because of the high x-ray transmission of Be, and a larger tube diameter, which delays the effects of transverse waves. The new design extended the time scale for observations of RTI growth in low-Atwood targets from 30 ns to >40 ns. However, the ablators did not meet specifications, compromising the physics of the experiment. The data are being analyzed to extract as much information as possible, and the team is working with the ablator manufacturer to ensure that the parts for the FY18 experiments meet all specifications.

This work is funded by LLNL under subcontract B614207.

**Proton Heating of Copper Foam on OMEGA EP**

Principal Investigator: J. Benstead

This LLNL–AWE campaign studies the heating of a cylindrical puck of copper foam irradiated by a short-pulse–generated proton beam. This shot day was an extension of previous shot days in 2014 and 2016 and featured a refined target and diagnostic design. The two major aims of the experiment were to measure the temperature distribution through the target and to quantify the extent of expansion of the rear surface.

The experimental setup is shown in Fig. 152.113. A gold foil was irradiated with the OMEGA EP sidelighter (SL) beam delivering 200 J over 3 ps. The SL produced a beam of protons and ions that were used to heat a copper-foam puck positioned 0.5 mm away. An aluminum foil was placed between the gold foil and the copper puck to improve heating by filtering out heavier ions and low-energy protons, which nonuniformly heat the target. The subsequent sample expansion was imaged with an x-ray radiography system. This used a nickel area backlighter, irradiated with three long-pulse beams, coupled to an x-ray framing camera (XRFC) that imaged the backlit target. The backlighter (BL) beams were delayed with respect to the SL beam in order to observe the heated and expanded target at different times. The streaked optical pyrometry (SOP) diagnostic was fielded orthogonally to the heating axis with its imaging slit oriented such that the temperature through the central section of the disk could be measured front to rear over the first 5 ns of heating (see Fig. 152.114); an RCF stack measured the proton/ion beam spectrum on each shot (see Fig. 152.115).

In total, seven shots were fired with data acquired on the XRFC, RCF, and SOP diagnostics. Full data analysis is still in progress, but preliminary results indicate that the degree of heating achieved was as desired. Slightly unusual features present on the SOP data have been attributed to the reduced target size relative to previous shot days, causing unexpected interactions with the SL pulse.
Figure 152.115
RCF data showing the attenuation of the proton and ion fields on one shot in the final four pieces of the film stack. The film pieces move progressively farther away from the target beginning at the bottom left image, then bottom right, then top left, and finally top right.

Development of Radiography-and-VISAR Platform for Hydrodynamics Measurements
Principal Investigators: M. Rubery (AWE) and D. A. Martinez (LLNL)
Co-investigators: G. Glendinning (LLNL); and S. McAlpin, J. Benstead, and W. Garbett (AWE)

As part of the LLNL/AWE Carisbrook Campaign, one and a half shot days of experiments were performed on the OMEGA Laser System during FY17. The platform consists of a halfraum and shock tube package driven by 15 × 500-J beams from the OMEGA H7 axis [Fig. 152.116(a)]. The objectives for these shot days were to diagnose the evolution of a hohlraum-driven interface using simultaneous point-projection x-ray radiography and VISAR, a configuration that was successfully demonstrated during FY16. A secondary objective for FY17 was to demonstrate the use of a reduced-mass backlighter [3-mm Ta disk versus 4-mm Ta square, Fig. 152.116(b)]. If successful, the new configuration should sufficiently reduce the amount of vaporized metal generated during the experiment to allow future campaigns to use the OMEGA EP short-pulse beam with debris shields removed (higher energy).

To generate the 7.8-keV He-like point-projection x-ray source, a 600-μm-sq nickel microdot was driven to 2 × 10^{15} W/cm^{2} using 4 × 400-J, 1-ns backlighter beams [Fig. 152.116(b)]. The x-ray emission is projected through a 20-μm Ta pinhole plate aligned along the shock tube and toward the TIM-6 cranked snout axis of θ = 123.1 and φ = 172.76. Images were recorded on film using a gated XRFC.

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### Development of Radiography-and-VISAR Platform for Hydrodynamics Measurements

<table>
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<tr>
<th>Beams</th>
<th># CPP’s</th>
<th>CPP size</th>
<th>Pulse</th>
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<td>15</td>
<td>IDI 300</td>
<td>1-ns square</td>
<td>500 J/beam</td>
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<tr>
<td>4× BL</td>
<td>Removed</td>
<td>280 μm</td>
<td>1-ns square</td>
<td>400 J/beam</td>
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Figure 152.116
(a) Radiography + VISAR configuration overview and (b) VISRAD isometric view of the experiment. TPS: target positioner system.
A quartz window, aluminum flash coating, and light shield cone were added to the rear of the target, allowing us to make VISAR measurements along the TIM-5 H14 axis.

The VISAR diagnostic performed well during all shots from Ca-17A and 17B. Figure 152.117 shows good-quality streak-camera images from both legs of the VISAR diagnostic on shot 86473. The SOP diagnostic, which uses the same optical relay as the VISAR, also produced good quality data on both shot days.

Unfortunately, during the first shot day, no radiography images were recorded. The source of this failure was found to be in the assembly of the backlighter. The mounting location of the nickel microdot was measured >100 μm away from the design position; this deviation is sufficiently large to move the nickel x-ray emission out of the XRFC field of view. In response to this, the backlighter design was modified to reduce sensitivity to assembly and target misalignment. The micro-dot was increased from 400 μm square to 600 μm square, and the backlighter beam spots increased from 200-μm to 280-μm diameter. The backlighter energy was also increased from 200 J/beam to 400 J/beam to maintain intensity. On the second half-day of shots these modifications were found to be 100% successful.

In addition, the new reduced-mass backlighter was successfully fielded on the second day and produced a radiograph with no observable drop in image quality, opening up the possibility of a future joint shot day, without the limitations on pulse energy introduced by parabola debris shields.

**Development of Gamma-Ray Sources for MeV Radiography**

Principal Investigator: F. Albert
Co-investigators: N. Lemos and J. Shaw (LLE); and D. A. Martinez and V. A. Smalyuk (LLNL)

This series of shots was designed to develop gamma-ray sources intended for a future MeV radiography capability on NIF’s Advanced Radiography Capability (ARC) short-pulse laser. Megavolt radiography on the NIF will serve a number of applications, such as double-shell implosions and imaging of dense objects.

The first FY17 campaign alternatively focused the backlighter and sidelighter short pulses (10 ps, 900 J) onto tantalum targets coated with 10 μm of plastic to produce hot electrons and subsequent gamma-ray emission from bremsstrahlung radiation (Fig. 152.118). The electron spectrum was measured with an electron positron proton spectrometer (EPPS) (along the short-pulse laser axis and also at 90°), and the gamma-ray spectrum with HERIE (a high-energy radiography imager) using a tantalum step-wedge filter pack. This diagnostic was used to retrieve an emitted photon spectrum of the form \( f(E) = A \exp(E/E_T) \), where \( E_T \) is the spectrum temperature, comprised between 0.5 and 1 MeV for this experiment. About \( 10^{12} \) to \( 10^{13} \) photons/eV/steradian were detected with this process. Targets included both 1-mm-thick, 1-mm-diam, as well as 500-μm-thick, 25-μm-diam, tantalum pucks. The source size was measured by imaging an 800-μm-diam tungsten sphere onto HERIE and was found to be around a few 100 μm for the thick targets. Further shots are required to determine the effect of the target geometry on the source size.
The second campaign (MeV-RadioEP-17B) aimed to look at an alternative electron production scheme (laser-wakefield acceleration in underdense plasmas) for future development of MeV photon sources on ARC. During this shot day, the 3- to 10-ps OMEGA EP short-pulse beams (alternately sidelighter and backlighter) were focused at intensities $10^{18}$ W/cm$^2$ onto a 3-mm plastic gas tube filled with helium at atmospheric pressure. The gas tube was closed with 1-$\mu$m-thick mylar windows, which were blown down using an OMEGA EP long-pulse beam timed 5 to 10 ns before the short pulse. Plasma density at the entrance of the gas tube was monitored with the 4$\omega$ probe diagnostic. The main diagnostic, EPPS, measured accelerated electron energies up to 50 MeV.

Analysis of these campaigns is ongoing, and the results of the two shot days will be used to design efficient gamma-ray sources for LLNL programs.

**Measurements of Instability Growth and Shell Trajectory Relevant to NIF Double-Shell Designs**

Principal Investigator: Y. Ping

Co-investigators: V. A. Smalyuk, P. A. Amendt, R. Tipton, J. Pino, O. L. Landen, F. Graziani

The goal of this campaign is to measure instability growth rate and shell trajectory in planar geometry under conditions relevant to the double-shell design on the NIF. The target is a halfraum with an attached physics package consisting of an ablator, CRF foam, and a Cu inner shell. For the instability growth measurements, ripples with 30-$\mu$m and 60-$\mu$m periods and 0.5-$\mu$m amplitude were imprinted on the Cu inner surface, and face-on gated x-ray radiography was employed to measure the ripple growth over time. For the shell-trajectory measurements, side-on x-ray radiography with a streak camera was employed. The shot day was very successful with excellent data in both configurations, as shown in Fig. 152.119. Data analysis shows reasonable agreement with simulations on the ripple growth, yet the observed preheat of the Cu shell was underpredicted in the modeling. These results are being organized for publications and will also be used for target designs in FY18.

**ACKNOWLEDGMENT**

This work was performed under the auspices of the U.S. Department of Energy by Lawrence Livermore National Laboratory under contract DE-AC52-07NA27344.
FY17 LANL Report on Omega Facility Experiments

In FY17, Los Alamos National Laboratory carried out 17 shot days on the OMEGA and OMEGA EP Laser Systems in the areas of high-energy-density (HED) science and inertial confinement fusion (ICF). In HED our focus areas were on radiation flow, hydrodynamic turbulent mix and burn, warm-dense-matter equations of state, and coupled Kelvin–Helmholtz (KH)/Richtmyer–Meshkov (RM) instability growth. For ICF our campaigns focused on the Priority Research Directions (PRD’s) of implosion phase mix and stagnation and burn, specifically as they pertain to laser direct drive (LDD). We also had several shot days focused on transport properties in the kinetic regime. We continue to develop advanced diagnostics such as neutron imaging, gamma reaction history, and gas Cherenkov detectors. The following reports summarize our campaigns, their motivation, and our main results from this year.

**BeBoron**

Shots were taken to measure the ablation rate of polyethylene to be evaluated for use as an ablator material for imploding capsules. Polyethylene was used at the last minute to replace Be:B flats that did not survive transportation, nor polishing. Polyethylene thin flats positioned at the end of the half-hohlraum and driven by the hohlraum radiation were used. This technique is similar to that used previously.\(^{41,42}\) The shock breakout data gave us good measurements at 175 and 204 eV, with an ablation rate between 3 to 4 mg/cm\(^2\)/ns (Fig. 152.120).

Because of the poor reflectivity of the targets, only three shots gave data. Shot 84428 at 100-eV drive gave good VISAR data with a velocity reaching 30 km/s, with the data being consistent with a reflecting shock in the quartz (Figs. 152.121 and 152.122).

**COAX**

The goal for COAX in FY17 was to complete platform development to prepare the way for addressing radiation-transport...
physics questions in supersonic → subsonic fronts on OMEGA and the NIF. To accomplish this, COAX had three shot days in FY17. This platform uses a halfraum to launch a radiation front down a cylindrical target featuring a doped aerogel foam coaxially contained in an undoped aerogel and a Be sleeve. This target is point-projection backlit by a V foil and pinhole target to radiograph the density spike that occurs when the radiation front cools and hydrodynamic expansion exceeds the rate of radiation flow. A 600-μm CH capsule with 1.5-atm Kr fill is used to backlight the target earlier in time to collect the temperature of the supersonic front from the ionization balance of dopant revealed through K-shell absorption spectroscopy of 1s–2p and 1s–3p Ti or Sc. Dante is used to measure the halfraum temperature through the laser entrance hole (LEH). As the impedance offered by the foam to the radiation front is increased, the slope of the front increases. This provides a window into radiation-transport physics in the supersonic to subsonic regime, which is important for astrophysical objects such as supernovae.

In November 2016, COAX collected useful radiography data of the subsonic front. The spectroscopic backlighter still needed development work to achieve high-resolution Ti and Sc K-shell spectra. We tested a method for evaluating the temperature sensitivity of the platform: putting an 8-μm-thick Cu ring or “top hat” around the lower half of the doped foam. In Fig. 152.123 we observed the difference between two shots in which the doped foam’s density was equivalent, but one target contained a top hat and the other did not. A simple examination of the data was shown at the 2016 APS DPP Meeting (Fig. 152.124).

**Figure 152.123**
Comparison of COAX radiography from targets (a) with and (b) without the copper “top hat.”

**Figure 152.124**
Data presented at APS DPP 2016 showing emission from the subsonic density spike superimposed on the spectral data.
We observed emission from the subsonic density spike superimposed on the spectral data for the first time.

In May 2017 we collected high-resolution (~1000 resolving power) Ti and Sc spectroscopic data (Fig. 152.125). We tested a new design for the radiography backlighter, which unfortunately did not meet specification. The quality of the radiography was reduced as a result. We drove some targets with 600-ps and some with 1-ns square laser pulses to evaluate whether increased coating time would improve the backlighter brightness. While that did occur, the physics target was much cooler than is typical because of the reduced drive energy. We repeated the observation of the subsonic front in the spectroscopic data with this higher-resolution, brighter spectra (produced as a result of improvements to the beam pointing for the capsule backlighter). A comparison between data from May and synthetic spectra produced with OPLIB in PrismSPECT was included in a presentation at the 2017 Anomalous Absorption Conference (Fig. 152.125).

In August 2017, we made a platform design change to move the physics target closer to target chamber center (TCC). This was partly to move the hohlraum back into the field of view for Dante and partly to improve backlighter flux through the physics target. We improved our design for the radiography V-foil backlighter and moved that target closer to TCC as well, to take advantage of the maximum drive range for the beams that this design allowed. We collected high-quality radiography and spectroscopy data with the new design (Fig. 152.126). We collected data at a sequence of times to study the transition between the supersonic and the subsonic flow over time. The times shown in Fig. 152.126 are the difference in time between the start of hohlraum drive and the time at which data were actually collected. At this year’s APS DPP Meeting we will discuss early analysis results for this data, such as shown in Fig. 152.127. To address radiation-transport physics questions precisely enough to constrain radiation–hydrodynamics models, we need $T_e$ uncertainty <5 eV and spatial uncertainty (from radiography) ~20 μm, which early analysis of our results...
(Fig. 152.127) implies we have achieved. Provided the rest of the analysis of the data from COAX-17B and COAX-17C is as promising, we will be starting the process for applying for NIF shots in 2018.

**CylDRT**

The first CylDRT/CylStalk shot day was scheduled on 29 August 2017. The main objective for this shot day was to measure the deceleration-phase Rayleigh–Taylor (RT) growth of the sine wave perturbations on the inner side of an aluminum marker layer in cylindrical geometry using face-on radiography.

We are reviving the cylindrical problem after more than ten years on OMEGA. The cylinder is 250 μm longer than what was previously used because of the requirement to use SG5 phase plates. The Target Fabrication Group had to re-establish their capabilities including an outside vendor (AlumiPlate) for the marker layer coating. The goal is to measure the deceleration RT growth of the \( m = 10 \) mode sine wave inner perturbation on the aluminum marker layer.

Figure 152.128 shows the axial and transverse views of the cylindrical target. The cylinder was 2.5-mm long and had a 986-μm outer diameter. The inner diameter of the cylinder was 860 μm. In the middle and inner side of the cylinder was a 500-μm-long aluminum marker band, and the inner aluminum layer had sine wave perturbations mode 10 with 3-μm amplitude. Figure 152.129 shows the experimental geometry.
The cylinder was driven by 40 beams with a 1-ns square pulse. An iron or nickel backlighter was attached to the TIM-4 end of the cylinder. The x-ray framing camera was on TIM-6. We used 12× magnification with a 15-μm pinhole onto a four-strip framing camera. We also used a TIM-2/TIM-3 side imaging of the target to see the uniformity of the implosion.

Figure 152.130 shows the experimental image from a smooth target that had no sine wave perturbation for reference. Figure 152.131 shows the experimental image for a target with mode-10 sine wave. The preliminary analysis of the experimental results shows deceleration phase RT growth agreeing with hydro calculations.
The Los Alamos Double-Shell team completed two planar double-shell (DSPlanar) experiments at the Omega Laser Facility in FY17. This campaign of experiments is part of the larger overall double-shell effort at Los Alamos, which is intended to test double-shell–relevant physics and materials in a planar geometry for ease of simulation and diagnosis. Specifically, these experiments are intended to validate our ability to predict hydro coupling, instability growth, with and without tamper mitigation, and the impact of target-fabrication artifacts such as fill tubes or joints in a double-shell target in the simplified planar geometry.

The first FY17 experiments examined momentum transfer during the ablator “inner shell” collision to benchmark RAGE simulations for similar systems. In these experiments, an indirectly driven ablator propagated down a shock tube [Fig. 152.132(a)] to impact a SiO\textsubscript{2} layer serving as an inner-shell surrogate. VISAR measurements [Fig. 152.132(b)] of the shock propagation, for pre- through post-ablator/layer impact, showed good breakout timing and average velocity agreement with pre-shot simulations.

The second FY17 shot day focused on studies of joint perturbation growth in support of NIF target-fabrication efforts.
These experiments again used indirectly driven ablators, but this time the ablator had a joint down the center. We varied the width of the joint gap as well as the geometry of the joint seam [Fig. 152.132(d)] for comparison to RAGE simulations. Early analysis suggests that RAGE simulations bound the observed joint growth rate and can serve to calculate an upper limit on the joint perturbation growth in a double-shell target.

HKMix

HKMix is a continuing LANL campaign to study mix physics in implosions between the fuel and shell, with the overarching goal of benchmarking models used in simulations of ICF and HED systems. These shots used 9- or 15-μm-thick, 860-μm-diam plastic shells with a thin deuterated (CD) layer in the shell and a HT gas fill. Using the LANL Cherenkov-based gamma ray detectors, both the HT (core) and DT (mix) burn histories can be measured. The timing difference between mix and core burn is a novel constraint on modeling these implosions; good data were acquired, with analysis and interpretation ongoing.

These campaigns also used much thinner deuterated layers (0.15 μm) than past work, which enables higher resolution of the effect mix depth. Past work with 15-μm-thick shells could be explained using turbulence-based mix models, and these were used to predict the mix (DT) yield versus CD layer recession depth for these shots, shown by the green curves in Fig. 152.133. The data, however, show a dramatic 100% decrease in mix yield as the CD layer is recessed by only 0.3 μm. This trend cannot be explained by a turbulence-based model, but good agreement is found with a diffusion-based model (black curve). At a 1-μm recession depth, the mix yield is observed to increase, indicating another “inversion” mix mechanism, potentially caused by a jet from the target support stalk.
In ICF experiments, interspecies ion separation is considered to be a candidate mechanism for yield degradation compared to radiation–hydrodynamic code predictions. This species separation could be driven by multispecies diffusive processes (e.g., thermodiffusion), as suggested by simulations employing recently implemented first-principles–based multi-ion–species transport models. The same physics also appears to drive ablator/gas mix in ongoing experiments focused on studying that problem.

Detailed analyses of x-ray–imaging spectroscopy data obtained from the earlier IonSepMMI-15A Campaign (in 2015) provided the first direct experimental evidence of interspecies ion separation in ICF experiments.\textsuperscript{46,47}

To obtain the earlier results, we assumed 1-D spherical symmetry in the analysis procedure. In the recent IonSepMMI-17A Campaign (August 2017), we recorded x-ray–imaging spectroscopy data along three different lines of sight, which will allow us to analyze the data without needing to assume spherical symmetry and to perform 2-D/3-D reconstructions of spatial profiles of ion densities to infer ion species separation in the implosion core. We conducted 12 target shots with a high return of x-ray–imaging spectroscopy data on four target types of varying fill pressure and dopant (Ar) concentration. We also collected time-evolution data of both x-ray and DD-neutron reaction histories in collaboration with MIT, which fielded their PXTD diagnostic for this purpose. We anticipate improved observations of stronger and weaker interspecies ion separation as a function of target type. We are now processing and analyzing the x-ray spectroscopy data (Figs. 152.134 and 152.135). Further processing/analysis is being conducted to infer the spatial profile of argon concentration versus time, which is the smoking gun for whether species separation is occurring.

Results from these campaigns will add to our experimental database for validating first-principles models of multi-ion–species transport and diffusion that have been implemented in LANL ASC (advanced simulation and computing) codes. The latter will allow us to better quantitatively assess the impact of species separation that are initially mixed, as well as the mix of species that are initially separated (e.g., ablator/fuel) in ICF implosion performance.
Marble VC

The Marble Void Collapse Campaign was developed to address (1) fundamental issues relevant to understanding of Marble implosions on the NIF and (2) simulation capability of macro-pore engineered foams. This year, the Marble team refined a laser-driven shock-tube platform and tested two topics. First, shock propagation through macro-pore engineered foams was investigated to examine if the pore size affects shock speed. Three types of macro-pore engineered foams (<1, 50, and 90 μm in diameter) were used in shock-tube experiments driven by the OMEGA laser. X-ray radiographic data indicate that shock speed through macro-pore engineered foams depends strongly on foam density, less on pore size [Fig. 152.136(a)]. Data were successfully used to validate LANL simulation capability. Second, we designed a single foam-filled void (250 μm in diameter) [Fig. 152.136(b)] and shocked it from two opposing directions, aiming to increase turbulence at the spherical boundary and as a result induce magnetic fields. X-ray radiographic data show that while the first shock compressed a spherical foam-void without much turbulence, the time-delayed second shock seems to increase a turbulent motion. D³He proton radiographic data (supported by the MIT group) were successfully obtained and are being analyzed.

MShock

The LANL MShock Campaign is studying the feedthrough of the RM instability in a thin layer, analogous to previous gas-curtain experiments. RM is relevant to mix in an ICF capsule where the ablative drive on capsule imperfections gives rise to RM and secondary shocks re-shock the linearly growing RM instability. It is known from fluids that such a re-shock can drive RM to turbulence. The MShock platform utilizes a beryllium shock tube analogous to the previous Shear Campaign. A thin high-density layer ~10× denser than surrounding foam is located a short distance from the first drive ablator. Two opposing laser drives with a 3-ns time delay directly drive ablators on the opposite side of the shock tube. This allows for a growth period between the initial excitation of the RM instability and the re-shock. X-ray radiography is used to capture the evolving layer; mix-width measurements are compared with the LANL Besnard–Harlow–Rauenzahn (BHR) turbulence model.

FY17 was the first year of the MShock Campaign. The first shot day was focused on platform development and proof of concept for the target design. Data were collected to calibrate shock timings with simulations, and a new central doping technique applied to the high-density layer was verified (Fig. 152.137). In addition, the experiment showed no distortions of the mix layer, which had been previously observed in the OMEGA Re-shock Campaign. The second shot day focused on capturing time sequences for three surface perturbations.
and collecting photometric data for dopants in the high-density layer. Preheat characterizations on this day showed that preheat effects are strong enough to affect surface profiles. Sufficient data were collected for one of the surface modes and nearly all data were collected for another. The results from both of these days will aid in preparing future experiments for both OMEGA and the NIF.

Figure 152.137
Data from a series of shots in the MShock Campaign.

**Oblique Shock**

The LANL Oblique Shock Campaign on OMEGA EP had two shot days in FY17 with a total of 14 shots. This campaign is designed to study the interplay between RM and KH hydrodynamic instabilities. It was conceived as a collaborative effort between LANL and the University of Michigan. The platform will also allow us to provide experimental input for testing and validation studies for turbulent transitional models like the LANL’s modal model, which will provide input for initial conditions for full turbulence models like BHR.

The first day was dedicated to testing (1) a 30-ns extended drive as an alternative to the typical 10-ns drive, and (2) the platform in a regime that more closely mirrors the University of Michigan’s hydro experiments on tilted interfaces (light to heavy). The Oblique Shock Campaign has been looking at shocks driven across an inclined (heavy-to-light) surface into a low-density foam (Fig. 152.138). Instead of three beams driving the target at one time, the three beams were stacked in time with 10 ns each, making a 30-ns train. In this configuration, the intensity driving the shock is reduced by a factor of 3, but the shock is supported for almost the entire experiment.

The second shot day was dedicated to testing a new multimode surface to study the effects of mode coupling on the growth of the interface. The 10-ns drive focused three beams...
on the ablator to drive a strong shock into the tube as in previous experiments; the setup was the same as Fig. 152.138. Figure 152.140 shows the imposed interface for studying mode coupling. The idea is for the mid modes to couple to the low modes, which would allow us to track the mode growth. The shot day showed very promising results; however, it also showed the limitations of our diagnostics and alignment procedures (Fig. 152.141). These are being enhanced for the next shot day to improve our resolution and contrast using a new Mn-Heα quartz asphere for the spherical crystal imager (SCI), which will allow us to use long-pulse (0.5- to 1-ns) beams to illuminate the backlighter targets; this in turn will allow us to produce far more x rays to view the layer and increase our resolution by using film instead of image plates. In the future, we will take advantage of the new SCI magnification (15×) to further increase resolution.

**WFEOS**

The goal of the WFEOS-17A shot day was to study the equation of state (EOS) of wetted-foam material. The LANL ICF Program is conducting liquid layer implosions on the NIF, where the liquid is wicked into a supporting foam shell on the inner surface of the capsule. The implosion dynamics and performance are sensitive to the EOS of the mixed foam/DT material because it sets the implosion adiabat, but the EOS of such mixtures has not been measured. The WFEOS-17A Campaign used the planar cryogenic capability on OMEGA with the active shock breakout (ASBO)/streaked optical pyrometer (SOP) diagnostic to measure the shock propagation through a liquid-D₂ wetted foam. Good data were acquired on several shots. An example interferometry streak from VISAR on shot 85712 is shown in Fig. 152.142. In addition to the shock release
from a high-density carbon pusher into the wetted-foam material, we observe a change in shock velocity at the wetted-foam interface with pure D$_2$, indicating a noticeable difference in the EOS of the two materials.

**FY17 SNL Report on Omega Facility Experiments**

**MagLIFEP and MagLIFSNL**  

The MagLIFEP and MagLIFSNL Campaigns at LLE operated by Sandia National Laboratories conducted a total of four shot days in FY17 (one on OMEGA and three on OMEGA EP) aimed at characterizing the laser heating of underdense plasmas (D$_2$, Ar) at parameters that are relevant to the magnetized liner inertial fusion (MagLIF) ICF scheme.$^{51,52}$ MagLIF combines fuel preheat, magnetization, and pulsed-power implosion to significantly relax the implosion velocity and $\rho R$ required for self-heating. Effective fuel preheat requires coupling several kilojoules of laser energy into the 10-mm-long, underdense (typically $n_e/n_c < 0.1$) fusion fuel without introducing significant mix. Barriers to achieving this include the presence of laser–plasma instabilities (LPI’s) as laser energy is coupled to the initially cold fuel, and the presence of a thin, polyimide laser entrance hole (LEH) foil that the laser must pass through and that can be a significant perturbation.

Experiments, having different goals, were performed on the OMEGA and OMEGA EP lasers. The objectives of the OMEGA EP experiments were to develop a spectrometer capable of viewing Ne K-shell emission ($h\nu = 920$ to $1100$ eV), and to continue to investigate the effects of pulse shaping and LEH foil thickness on energy coupling. Capturing Ne spectra required developing a new spectrometer based on the spatially (in 1-D) and temporally resolved multipurpose spectrometer (MSPEC) design that used a KAP or RbAP with a maximized Bragg angle giving an observable energy range of 891 to 1773 eV (Ref. 53). In addition, modifications to the targets were required that allowed soft x rays to escape while still accommodating high-pressure (up to 10 atm) gas fills. This was achieved by machining up to five slots in the sides of the 115-μm-thick CH tube target and covering the slots with a 2-μm-thick polyimide foil, as shown in Fig. 152.143(a). Capturing Ne spectrum was challenging but was achieved as shown in Fig. 152.143(b). The instrument should allow one to observe cooling of the plasma after the laser has turned off, potentially facilitating the study of magnetic field’s effect on the electron thermal conduction and cooling process.

The objective of the OMEGA laser experiments was to compare energy coupling into underdense D$_2$ plasmas with $2\omega$
and $3\omega$ beams, with and without smoothing by spectral dispersion (SSD), and for different beam intensities. This shot series was the first Sandia-led effort to investigate preheat on this laser, and, as such, scaled-down versions of the OMEGA EP targets were required because of the reduced energies per beam available ($\sim 450 \text{ J max compared to } >3 \text{ kJ on OMEGA EP}$). Of particular interest in this series was the effect that changing the laser wavelength, intensity, and smoothing had on-beam propagation and on-stimulated Brillouin scattering (SBS) and stimulated Raman scattering (SRS) levels. The results suggest that for similar intensities and similar values of $n_e/n_c$, the laser wavelength has a significant effect on the beam propagation, as measured by x-ray framing camera (XRFC) imaging [Fig. 152.144(a)], and on LPI levels, as inferred by hard x-ray signal levels [Fig. 152.144(b)], while SSD has little impact. The results have implications for the future of MagLIF laser preheat, which currently uses a $2\omega$ laser and is susceptible to significant LPI at relatively low intensities. The data suggest that a $3\omega$ laser could make preheat possible at higher intensities for given plasma parameters while minimizing LPI, ensuring that greater energy coupling could be possible over a given propagation distance.

ACKNOWLEDGMENT

FY17 NRL Report on Omega Facility Experiments
During FY17, an NRL/LLE collaboration on laser imprint led to three successful shot days on OMEGA EP. The experiments showed that the application of a prepulse that pre-expands and lifts off the coating prior to the arrival of the main laser pulse gives an order-of-magnitude reduction of laser imprint, as expected on the basis of the original experiments on the Nike laser and understanding the mechanism of the imprint suppression (Fig. 152.145). Further experiments demonstrated imprint reduction with prepulse times compatible with pulse durations available for implosions on the Omega 60-beam laser. Moreover, we were able to utilize thinner coatings, minimizing the risk of fuel preheat and increasing the chances that they will be compatible with low-adiabat, thin-ablator shell implosions on OMEGA.
Multibeam effects in LPI at the megajoule scale were first considered through the cross-beam energy transfer between two beams that was used to adjust the symmetry of irradiation. Recent experiments performed in near-vacuum hohlraums have revealed another multibeam effect, namely the electromagnetic seeding of the sidescattering for some beams by the backscattering or transmission of additional beams. The issue of collective scattering instabilities has recently become an emerging field with several reported experimental studies. In this context, two experiments have been performed on OMEGA to investigate collective SRS and collective SBS. The SRS of two beams has been evidenced in an inhomogeneous plasma produced in an open planar geometry where the significant amplification of the Raman light at large angles from the density gradient has been observed for the first time. The collective Brillouin amplification of shared IAW’s has been observed in indirect-drive experiments using a rugby-ball–shaped hohlraum. In both types of experiments, the flexibility of the Omega Laser Facility and its large battery of diagnostics have played a critical role toward the physical understanding of the collective mechanisms.

The first experiment was performed to investigate the collective SRS produced by two beams sharing a common electromagnetic daughter wave in an inhomogeneous plasma. The targets were 7-mg/cm² C12H16O8 foams with a diameter of 2.5 mm and a length of 950 μm aligned along the H3–H18 axis of the OMEGA target chamber. The laser beams were focused by f = 6.7 lenses through elliptical phase plates, producing spot sizes with a 200 × 300-μm diameter [full width at half maximum (FWHM)]. These beams were fired at an energy level of 400 J in a 1-ns square pulse. Twelve beams were used, incident at 60° from the foam axis, making a six-beam cone on each side of the target, as is illustrated for the H18 side in Fig. 152.146(a). After ~0.5 ns, the regions of foam ionized by the different beams of the same cone began to superimpose.

Figure 152.146(b) shows a typical measurement with the near-backscatter imager (NBI) of the time-integrated angular distribution of the light scattered in the Raman wavelength range (450 to 900 nm) around the midplane of Beamlines 45 and 50 with angles between 20° and 60° from the target normal. The SRS signal was maximum in the bisecting plane of beams 45 and 50 at angles between 42° and 54° from the target normal, close to the full-aperture backscatter station of Beamline 25 (FABS25). Figure 152.146(c) shows the time-resolved spectrum of the light scattered in this direction collected in an aperture of ±4° in FABS25 for the same shot. The SRS signal started at ~0.5 ns, as soon as the beams started overlapping in the foam plasma. It lasted until the end of the laser pulses with an almost constant wavelength of λSRS ~ 600 nm, corresponding to the electron density (~0.17 n₀) in the region of beam overlap.

This observation in vacuum of the SRS scattered light at ~48° from the target normal corresponds to SRS light produced at ~80° from the density gradient in the plasma region of interaction. This optimum angle results from the increase of
giving a hot-electron temperature of 20 keV compared to an expected temperature of 25 keV.

The collective nature of the coupling and the amplification at large angles from the density gradient are found to increase the global SRS losses and to produce light scattered in novel directions outside the planes of incidence of the two beams. The mechanism evidenced in this experiment can occur in both direct- and indirect-drive experiments between any pair of beams. Such a collective amplification of a common scattered-light wave has been proposed to explain the large amount of SRS measured in the NIF indirect-drive experiments. It results in an underestimation of the scattered-light losses because this red-shifted Raman light experiences significant refraction and absorption before exiting the plasma and in the acceleration of hot electrons, which can preheat the capsule.

The second experiment was performed to investigate collective Brillouin scattering in indirect-drive experiments. The interaction was studied at the laser entrance hole (LEH) of rugby-ball–shaped hohlraums filled with 1 atm of methane gas. The revolution axis of the hohlraum was aligned along the P5–P8 axis of the OMEGA target chamber. Twenty laser beams were incident on each side of the target distributed along three cones [see Fig. 152.147(a)]: five beams, at 21° from the axis, were pointed at 500 nm from the window outside the hohlraum; similarly, five beams at 42° and ten beams at 59° were pointed at the LEH. The beams were focused by $f = 6.7$ lenses through random-phase plates, producing focal spots with a diameter of 300-μm FWHM. The laser pulse, with a total duration of 2.5 ns, was made of a prepulse of intensity $\approx 8 \times 10^{13}$ W/cm² per beam followed by a main pulse for $t = 1.8$ to $t = 2.3$ ns at an intensity per beam of $5 \times 10^{14}$ W/cm². The plasma conditions were characterized thanks to thermal Thomson-scattering (TS) measurements performed at the LEH and at 300 μm from the LEH outside the hohlraum as illustrated in Fig. 152.147(a).

The scattered light was collected in the backward direction of one beam of the 59° cone (Beamline 30) and one beam of the 42° cone (Beamline 25) and analyzed in time and wavelength by full-aperture backscatter stations (FABS25 and FABS25). Figure 152.147(b) shows a typical spectrum measured in FABS25. For the time interval (1.8 ns to 2.4 ns), when the laser intensity was close to its maximum, two contributions were detected in FABS25. The first signal, starting at $t = 1.8$ ns, was caused by Brillouin backscattering of Beamline 30 developing in the plasma inside the hohlraum. The backscattering was identified in the single-beam linear gain calculations performed by post-processing the hydrodynamics simulations. It simply

![Figure 152.146](a) Angular distribution of the six beams incident on the H18 side and scheme of the near-backscatter imaging (NBI) diffuser plate. (b) Typical image recorded with the NBI diagnostic for a 7-mg/cm³ foam. (c) Time-resolved spectrum of the stimulated Raman scattering (SRS) light collected in the direction of FABS25 (full-aperture backscatter station) in the same shot as for (b).
the Thomson-scattering measurements for the light scattered by Beamline 66 off the collective IAW’s generated by the ten beams at 59° [IAWa in Fig. 152.147(c)]. The observation of a sidescattering signal peaked in wavelength was a strong indication of multibeam effects. In a different experiment, using a different pointing of the ten beams of the 59° cone to improve the uniformity of irradiation on the hohlraum wall, significant spectral broadening of this second signal was observed as expected from the widening of the crossing volume of the ten beams of the 59° cone.

The absolute energy measurements performed in the FABS made it possible to evaluate the energy losses caused by the collective SBS instability of the 59° cone. To do so, we first considered the geometry of the scattering, taking into account the aperture of the beams and the diagnostic. The geometry of the SBS of Beamline 66 in the collective instability that drove the IAW along the hohlraum axis (IAWa) is shown in Fig. 152.147(c), assuming straight-line propagation of the light in front of the crossing-beam region at LEH. Its analysis showed that only a small fraction (<1/15) of the scattered light was collected in the FABS. The signal associated with the collective instability of the ten beams peaked at +3% of the laser power per beam, but we estimate that the scattering losses may be 10× higher than those directly measured in the FABS. Therefore, collective Brillouin scattering can result in high scattering amplitude and significantly impair the laser–target coupling in indirect-drive experiments.

REFERENCES

Publications and Conference Presentations

Publications


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Conference Presentations


The following presentations were made at the Industrial Associates Fall Meeting 2016, Rochester, NY, 9–12 October 2016:


The following presentations were made at Frontiers in Optics, Rochester, NY, 17–21 October 2016:


J. M. Schoen, “History of the Center for Optics Manufacturing” (invited).


The following presentations were made at the 37th Tritium Focus Group Meeting, Rochester, NY, 25–27 October 2016:


W. T. Shmayda, “Properties of DT Ice in Cryotargets.”


The following presentations were made at the 58th Annual Meeting of the APS Division of Plasma Physics, San Jose, CA, 31 October–4 November 2016:


S. Bucht, D. Haberberger, J. Bromage, and D. H. Froula, “Transforming the Idler to Seed Raman Amplification.”

D. Cao, P. W. McKenty, J. P. Knauer, and D. R. Harding, “Investigation of Acquired Fuel Motion Caused by Ice Roughness in OMEGA Cryogenic Experiments.”


D. Clarkson, R. Ume, R. Sheets, S. P. Regan, T. C. Sangster, S. Padalino, and J. McLean, “Bulk Etch Rate and Swell Rate of CR-39.”


K. Cook, M. Coats, M. Yuly, S. Padalino, T. C. Sangster, and S. P. Regan, “Measurement of the 6He Decay Produced by the 9Be(n,α)6He Reaction.”


L. Hao, R. Yan, J. Li, and C. Ren, “Development of a New Fluid Code to Study Laser-Plasma Instabilities.”


J. Li, R. Yan, and C. Ren, “Density-Modulation-Induced Absolute Laser-Plasma Instabilities in Inertial Confinement Fusion.”


A. Pak, “Shock-Wave Acceleration of Protons on OMEGA EP.”


The following presentations were made at the 40th IEEE EDS Activities in Western New York Conference, Rochester, NY, 4 November 2016:


The following presentations were made at the Rochester Academy of Science 43rd Annual Fall Session, Rochester, NY, 12 November 2016:


The following presentations were made at the 2016 International Workshop on Radiative Properties of Hot Dense Matter, Santa Barbara, CA, 5–9 December 2016:


The following presentations were made at the NIF and JLF User Group Meeting, Livermore, CA, 6–8 February 2017:


The following presentations were made at the IAEC–NNSA Meeting on Hydrodynamic Instabilities in HED Systems, Livermore, CA, 8–10 February 2017:

R. Betti, “Deceleration Phase Hydrodynamic Instabilities, Pressure Degradation from Low to High (Mid) Modes.”


The following presentations were made at the 22nd Target Fabrication Meeting, Las Vegas, NV, 12–16 March 2017:


N. D. Viza and D. R. Harding, “Performance of Different ‘Lab-on-Chip’ Geometries for Making Double Emulsions for Polystyrene Shells.”

M. D. Wittman, M. J. Bonino, C. Fella, and D. R. Harding, “Effect of Tritium-Induced Damage to Plastic Targets from High-Density D-T Permeation.”


The following presentations were made at the 13th Direct Drive and Fast Ignition Workshop, Salamanca, Spain, 22–24 March 2017:


The following presentations were made at the Ninth Omega Laser Facility Users Group Workshop, Rochester, NY, 26–28 April 2017:


L. H. Xiao, R. S. Craxton, D. Barnak, and J. Davies, “Simulations of Laser-Driven Magnetized Liner Inertial Fusion.”


The following presentations were made at CLEO 2017, San Jose, CA, 14–19 May 2017:

S.-W. Bahk, C. Dorrer, and J. Bromage, “Two-Dimensional Characterization of Spatiotemporal Coupling of Ultrashort Pulses Based on Chromatic Diversity.”

C. Dorrer and J. Hassett, “High-Accuracy, Model-Based Laser Near-Field Beam Shaping.”


The following presentations were made at the Sixth International Conference on High Energy Density Physics, Shirahama, Japan, 5–9 June 2017:

A. B. Sefkow, “Adventures in ICF and HEDP with Magnetic Fields.”


The following presentations were made at the 47th Annual Anomalous Absorption Conference, Florence, OR, 11–16 June 2017:

S. Bucht, D. Haberberger, J. Bromage, and D. H. Froula, “Transforming the Idler to Seed Raman Amplification.”


A. B. Sefkow, “Adventures in ICF and HEDP with Magnetic Fields.”

A. B. Sefkow, J. M. Koning, M. R. Gomez, S. B. Hansen, K. Cochrane, C. Thoma, D. R. Welch, and M. M. Marinak,
“Unprecedented Stability in Z-Pinch Implosions Due to Magnetic Fields and Plasma Physics.”


The following presentations were made at the 16th International Superconductive Electronics Conference, Sorrento, Italy, 12–16 June 2017:


The following presentations were made at the 20th Conference on Shock Compression of Condensed Matter, St. Louis, MO, 9–14 July 2017:

Celliers, “Hugoniot Measurements of Silicon Shock Compressed to 25 Mbar.”


The following presentations were made at the 20th International Conference on Electron Dynamics in Semiconductors, Optoelectronics, and Nanostructures, Buffalo, NY, 17–21 July 2017:


The following presentations were made at High Energy Density Science Summer School, La Jolla, CA, 30 July–11 August 2017:


G. W. Collins, “Physics of Matter at Extreme Pressure.”

Y. H. Ding, “A First-Principles Equation-of-State Table of Beryllium for High-Energy-Density Plasma Simulations.”


A. Hansen, “OMEGA Supersonic Gas-Jet Target System Characterization.”


M. Stoeckl and A. Kozlov, “Dependence of Readout Fade Rate on X-Ray Energy for BaFBr$_{0.85}$I$_{0.15}$:Eu Image Plates.”


The following presentations were made at Liquid Crystals XXI, San Diego, CA, 6–10 August 2017:

Liquid Crystal Devices. II. Transition-State Modeling in Azo-benzene and Spiropyran Oligomers.”


The following presentations were made at the 10th International Conference on Inertial Fusion Sciences and Applications, Saint Malo, France, 11–15 September 2017:


The following presentations were made at the 49th Annual Symposium on Optical Materials for High Power Lasers, Boulder, CO, 24–27 September 2017:


The following presentations were made at the 11th International Laser Operations Workshop, Rochester, NY, 26–28 September 2017:


L. J. Waxer, C. Dorrer, E. M. Hill, A. Kalb, and W. A. Bittle, “Development and Implementation of a Single-Shot Diagnostic for Characterizing 0.5- to 250-ps Pulses on OMEGA EP.”

