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About the Cover:

The cover depicts the phase diagram of magnesium oxide (MgO) and the newly measured melting curve (solid black). Various theoretical predictions for the phase diagram and melting curve are given with dashed–dotted lines. Core–mantle boundary conditions of Saturn and of 1-, 7.5-, and 15-Earth-mass super-Earths are indicated. The principal Hugoniot (blue curve) defines the states that are accessible with a single shock wave. With single-shock experiments, the melt curve can only be explored up to the pressure where it crosses the principal Hugoniot—600 GPa in MgO. A different technique was necessary to probe melting at higher pressures. A double-shock self-impedance–matching technique was used to measure the melt curve of MgO to 2000 GPa (20 Mbar); this is the highest pressure to which any material's melt curve has been probed experimentally.

On the cover figure, solid gray circles represent the first shock B1 states in this work; pressure was measured and temperature was taken from previous work on the principal Hugoniot of MgO. Red



open and solid circles are the second shock states; both pressure and temperature were measured. The solid red circles indicate points that are interpreted to lie on the melt curve of MgO because of a lack of observed temperature increase over a wide pressure range due to the latent heat of melting. We find that at 1950 GPa, the measured melting temperature of MgO is 17,600 K; this is 17% lower than recent theoretical predictions (purple dashed–dotted curve). This double-shock technique, depicted in the figure above, will lead to new advances in probing phase-transition behavior in transparent materials to multi-terapascal conditions.

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

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LLE Review Quarterly Report



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In Brief

This volume of LLE Review 166 covers the period from January–March 2021. Articles appearing in this volume are the principal summarized results for long-form research articles. Readers seeking a more-detailed account of research activities are invited to seek out the primary materials appearing in print, detailed in the publications and presentations section at the end of this volume.

Highlights of research presented in this volume include:

- T. W. Overton reviews the many successes both in advancing inertial confinement fusion research and training the next generation of scientists that have been enabled by the close collaboration between LLE and General Atomics (p. 67).
- W. Y. Wang and R. S. Craxton propose pentagonal prism hohlraum experiments on OMEGA as a test bed for high-symmetry hohlraum experiments on future laser facilities (p. 76).
- K. M. Woo and R. Betti develop an analytic model for the impact of 3-D ρR asymmetries on the generalized ignition criterion that allows the degradation of the Lawson criterion to be inferred from ion-temperature asymmetry measurements (p. 81).
- J. L. Shaw *et al.* use OMEGA EP as the driver for a laser-plasma accelerator that generated a relativistic electron beam with charge exceeding 700 nC and laser-to-electron conversion efficiencies up to 11% (p. 83).
- K. L. Nguyen *et al.* investigate the nonlinear saturation of cross-beam energy transfer (CBET) using VPIC simulations. Trapping induced modification of the ion-velocity distribution was found to detune the CBET resonance and limit the gain (p. 86).
- J. L. Peebles *et al.* use the magneto-inertial fusion electrical discharge system (MIFEDS) to collimate a relativistic chargeneutral electron–positron beam generated using OMEGA EP (p. 90). The technique has the potential to generate pair plasmas that can be used to simulate astrophysical phenomena.
- S. Jiang *et al.* demonstrate enhanced positron production using microstructured targets in high-intensity laser-plasma interaction experiments on OMEGA EP (p. 93).
- D. H. Froula *et al.* provide an overview of the use of Thomson scattering as a spatially and temporally resolved diagnostic of plasma conditions and the associated velocity distribution functions (p. 95).
- L. E. Hansen *et al.* report on equation-of-state measurements of CO₂ up to 800 GPa using laser-driven diamond-anvil-cell targets (p. 101).
- M. C. Marshall *et al.* investigate the metastability of the liquid-to-ice VII phase transition under rapid compression and find that the liquid phase can persist at pressures up to 4× higher than the onset of metastability (p. 104).
- M. D. Bergkoetter *et al.* develop a method for laser wavefront phase retrieval in the presence of chromatic aberrations (p. 107). Forward fitting is used to retrieve the wavefront phase at the pupil plane based on a model where both monochromatic and chromatic aberrations are modeled using expansions over Zernike polynomials.
- I. A. Begishev *et al.* demonstrate high-efficiency (37%) optical parametric chirped-pulse–amplification on the Multi-Terawatt Laser System using large-aperture DKDP crystals (p. 111).

- C. Dorrer *et al.* develop a novel sum-frequency generation technique for broadband frequency conversion (p. 114). They demonstrated the scheme by combining the broadband 1ω output of an optical parametric amplifier with narrowband 2ω light, resulting in broadband 3ω with ~10 THz of bandwidth.
- V. V. Ivanov *et al.* describe magnetized plasma experiments performed on the Zebra pulsed-power machine (p. 117). Megagauss magnetic fields were used to significantly modify the plasma expansion from an interacting laser pulse.
- J. Puth et al. summarize operations of the Omega Laser Facility during the second quarter of FY21 (p. 121).

Russell K. Follett Editor

LLE and General Atomics: A Partnership for the Future

T. W. Overton

General Atomics

The Three-Decade Relationship has been a High-Impact Influence on Inertial Fusion Research

For a collaboration that is often measured in microns, it is a bit ironic that the Laboratory for Laser Energetics (LLE) and General Atomics (GA) are more than 2500 miles apart. LLE was established in 1970 as a center for the investigation of the interaction between matter and intense laser radiation (Fig. 1). The development of a series of high-powered, neodymium-glass laser systems at LLE (DELTA, ZETA, and the 24-beam OMEGA) eventually led to the 60-beam OMEGA Laser System. The 60-beam OMEGA laser has been operational since 1995 and is one of the primary research tools for inertial confinement fusion (ICF) and high-energy-density (HED) physics research in the U.S.



Figure 1 Exterior of the Laboratory for Laser Energetics at the University of Rochester.

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OMEGA is maintained and operated by LLE for the Department of Energy's (DOE's) National Nuclear Security Administration (NNSA). OMEGA can focus up to 30,000 J of 351-nm laser energy onto a target that measures less than 1 mm diam in approximately one billionth of a second. In addition to the 60-beam OMEGA, the 4-beam OMEGA EP laser has also been a main research tool for the community since 2008.

OMEGA EP consists of four beamlines similar to those at the National Ignition Facility (NIF). Two of those beams can be compressed for short-pulse, petawatt-class operation. Both of these facilities allow scientists to explore physics conditions at extremely high pressures and temperatures—including fusion, the process that powers the Sun. Approximately 60% of all experiments conducted at the Omega Laser Facility are led by researchers from outside LLE.

Since 1991, nearly all of the LLE target capsules, which hold the material that is compressed by high-powered laser pulses, have been manufactured by GA in San Diego. Dr. Mike Campbell, the director of LLE, is frank about the close interdependence

of LLE and GA. "LLE is the main facility for direct-drive fusion research, where the laser directly impinges on the fusion capsule. So, in addition to the laser and diagnostics that make up the facility, you need to have targets to shoot. All those capsules are made for us at General Atomics, and the progress that we can make in the fusion program is directly dependent on the characteristics of the targets that GA delivers."

"It's been a three-decade, high-impact relationship," said Mike Farrell, vice president of inertial fusion at GA. "Whether it's been engineered systems, targets, diagnostic instrumentation, or other activities in support of the science that researchers are conducting at LLE, GA's contributions to enabling the physics research performed at LLE have been a constant and significant enabling factor."

GA's involvement in fusion research reaches back to its founding in the 1950s. Although its global operations now extend into technologies as diverse as aerospace and biotech, fusion remains a core focus of its research and development activities (Fig. 2). For decades, GA has worked with DOE and its predecessor agency, the Atomic Energy Commission, on a wide variety of fusion energy initiatives. Many GA employees have worked at DOE laboratories and facilities like LLE—including Dr. Campbell, who headed up GA's target fabrication operations from 2000 to 2007.



Figure 2 The General Atomics campus in San Diego (photo Ronalyn Conception/GA).

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Target Development

The LLE–GA collaboration began in the early 1990s, when DOE awarded GA the contract to fabricate targets for the U.S. ICF program, which supports much of the research at LLE (Fig. 3). One of the first major projects to emerge from LLE's work with GA was an upgraded Cryogenic Target Handling System. Direct-drive fusion targets on OMEGA need to be filled with fusion fuel, which is a mixture of the hydrogen isotopes deuterium and tritium (DT). In order to increase the amount of fuel each target can hold and to start with the highest density, the DT fuel is frozen at cryogenic temperatures. These tiny capsules are then placed into the target chamber.

Dr. David Harding, group leader for target fabrication at LLE, joined the team during development of the new system in 1995. "On the old system, the researchers would fill the targets with a very small amount of tritium at room temperature so that the capsule wouldn't burst, and put it into the cryostat at 19 K to create a very rudimentary DT ice layer," he explained.

This was sufficient for early experiments in the late 1980s, but researchers wanted to shoot targets with significantly more fuel. The goal of the new system was to load very thick (up to ~100- μ m) DT layers into very thin-walled (less than 5- μ m) polymer capsules (Fig. 4). The experiments also required that the ice wall be uniformly thick with a smooth surface. As might be expected, creating targets with these attributes was a significant challenge.



Figure 3 The (a) Target Bay and (b) Laser Bay at LLE during a laser shot.



Figure 4 One of the cryogenic targets used on OMEGA, fabricated by GA.

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"The capsules are filled with fuel by diffusion, which is placing the capsules under more 1000 atm of pressure in a vessel filled with the fuel gas, so that it diffuses across the thin capsule walls," Dr. Harding said. "After you've pressurized the capsule, you have to freeze the fuel to keep it from bursting once the target is removed from the pressure vessel."

The challenge was doing this in a way that the capsules could be filled, frozen, placed into the target chamber, and shot before the ice thawed and the capsules burst. "The job was to design a cryostat with a pressure vessel inside it, so you could ramp the pressure up, cool the targets, and do it in a very controlled manner," Dr. Harding said.

On top of that, the targets still needed to get to the chamber where they could be shot by the lasers. The team designed a smaller, mobile cryostat that would take the target from the larger cryostat and place it into the target chamber. To maintain the cryogenic temperatures, the target would be covered by a copper shroud during transport.

That might sound simple, but the process needed to operate in an extremely precise manner. The target had to be placed within 5 μ m of the center of the target chamber, and the cryogenic protective shroud had to be retracted very rapidly, less than 100 μ s prior to the shot.

The challenging demands of this system led to some initial difficulties in development. GA and LLE began working together on the process in 1995, and GA designed all the major components with assistance from LLE. The parts were then shipped to Rochester and assembled, integrated, and tested by LLE staff. "It was a complete collaboration between LLE scientists and engineers and GA's personnel," Dr. Harding said. "It was a great success, especially given that it was completed in such a short time."

The first target using the new system was shot on 14 July 2000—just five years after the project began. The new system was a major improvement over the old one on the 24-beam OMEGA laser. It was able to process and shoot four targets per day in exactly the conditions the researchers needed (Fig. 5). "There were no compromises on the physics quality of the targets," Dr. Harding said.



Figure 5

Operators from LLE's Cryogenic and Tritium Facility are shown performing routine maintenance and performance checks on moving cryostat transport carts (MCTC's) in the Cart Maintenance Room. Steven Verbridge and Chad Fella (foreground) are preparing to raise the moving cryostat (MC) from MCTC#2 in order to perform service. Sean Adams and Michael Coffey (background) are employing a coordinate measuring machine on an (elevated) MC on MCTC#7 to verify proper alignment of the removable shroud. Operators are wearing disposable jackets and gloves given the trace tritium surface contamination on internal equipment.

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The system was so successful that elements of the tritium-handling approach were used to inform the design of the cryogenic target system on the NIF at Lawrence Livermore National Laboratory (LLNL), which GA also developed. LLE shared its experiences with the new system with GA, which incorporated it into the work on the NIF. "Many of the features in our cryogenic system fed directly into the NIF system," Dr. Harding said.

Another innovation that flowed from the LLE–GA partnership related to the composition of the polymer target capsules. Researchers at LLNL first demonstrated a process for making the polymer shells known as glow-discharge polymerization (GDP). However, LLNL had little interest in pursuing the process at the time. LLE staff asked GA to improve and develop the technology for use on OMEGA. "We used the targets and told GA how they performed and what changes they needed to make," Dr. Harding said. "They used results from OMEGA implosions to further define the target specifications. That technology is now the baseline target fabrication methodology."

Working together, GA and LLE scientists have continued to refine methods for producing the polymer capsules and making them much smoother and free from defects. This is important because any defect or imperfection in the capsule serves as a "noise" source from which hydrodynamic instabilities can grow and adversely affect the quality of the implosion experiments. At LLE's request, GA developed a technique for producing polystyrene shells using a solvent-based microencapsulation method, resulting in capsules that are significantly smoother than GDP targets (Fig. 6). "GA was able to control the uniformity and thickness of the shells and make them many, many orders of magnitude smoother than the GDP targets," Dr. Harding said.

Diagnostics

Another key element of LLE's work are the diagnostic systems that record data from the laser shots. After all, the experiments are of little use unless scientists are able to analyze exactly what happened. Here too, LLE and GA have worked together



Figure 6

The improvement in smoothness between (a) 1-mm-diam GDP shells and (b) polystyrene shells, both made by GA for experiments on OMEGA. The shell wall is 0.01 mm thick. The notable difference is the absence of features on the surface of the polystyrene shell, which is important for high-performing implosions.

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to bring significant innovations to the field. "In these implosion experiments, we need to know the shape of the central hot-spot plasma formed in a laser-direct-drive implosion," said Dr. Sean Regan, LLE's Experimental Division Director. "To do that, we need some very specialized instrumentation."

The ideal implosion is spherically symmetric. "We don't get that," Dr. Regan said, "because we don't position the target accurately enough, or the laser is stronger on one side than the other, or some other factor. So, we need to understand the causal relationships between the implosion inputs and the outputs."

Several years ago, through the National Diagnostics Working Group, LLE joined a productive collaboration between GA, LLNL, Sandia National Laboratories (SNL), and Kentech Instruments that involved the development of a single-line-of-sight (SLOS) camera that can capture multiple images with a shutter speed of about 25 ps (Fig. 7). GA has delivered two versions of the SLOS instrument so far. The SLOS-CBI (crystal backlighter imager) was delivered to LLNL in 2017 and is in use on the NIF. The SLOS-TRXI (time-resolved x-ray imager) was delivered to LLE and is now operating as part of OMEGA (Fig. 8).



Figure 7

Terry Hilsabeck and Kyle Engelhorn of General Atomics with the SLOS-TRXI instrument during assembly (photo: Eugene Kowaluk).

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SLOS-TRXI is a primary diagnostic for the DT cryogenic implosion campaigns on OMEGA (Fig. 8). It images the hot-spot plasma emission with 25-ps temporal resolution and 10- μ m spatial resolution using a pinhole camera and a time-dilation tube. To put that speed in perspective, 25 ps is the time it takes a beam of light to travel almost half of an inch. As impressive as this is, the researchers at the Omega Laser Facility are not satisfied with the status quo. "Because the hot spot is about 40 μ m," Dr. Regan said, "a 10- μ m resolution doesn't give you a very crisp image. We want to do better. We also want to diagnose the low-mode structure of the hot-spot plasma."



Figure 8 The SLOS-TRIXI instrument being installed on OMEGA.

The team is currently designing what might be termed "SLOS-TRXI 2.0" that will add a third line of sight to the instrument, allowing for a full 3-D low-mode reconstruction of the hot-spot plasma. Under this same collaboration, a hot-spot x-ray imager with 20-ps temporal resolution and $5-\mu m$ spatial resolution is being developed for OMEGA. This imager will allow scientists to determine the overall shape of the hot spot and how it deviates from a perfectly spherical shape. This can also be compared to data from the other OMEGA diagnostics.

"Using nuclear spectroscopy measurements recorded along three quasi-orthogonal diagnostic lines of sight, we can infer the hot-spot flow velocity," Dr. Regan said. "If there's a significant flow in the hot spot caused by an asymmetric implosion, it will doppler-shift the DT fusion neutron spectrum, and you'll see the mean energy of the spectrum shifted up or down. Combining the hot-spot flow measurements with the SLOS-TRXI x-ray images of the hot spot allows us to model the structure of the hot spot. This will ultimately help us understand how the hot-spot formation is affected by multidimensional effects on the implosion."

The working group is drawing on lessons learned from SLOS-TRXI to develop the third line-of-sight instrument. LLE is responsible for project oversight and design and the x-ray pinhole camera, while GA is developing the drift tube. SNL, LLNL, Kentech, and Sydor Technologies are developing other elements. A conceptual design has been developed, and discussions about the construction have recently started. "It's a fantastic collaboration between LLE, GA, the national labs, and these other innovative and highly talented private firms," Dr. Regan said.

National Laser Users' Facility

Another key element of the collaboration has occurred through the DOE's National Laser Users' Facility (NLUF) program, which provides beam-time access at the Omega Laser Facility for scientists in both academia and private industry to conduct basic research and train graduate students. These experiments explore a wide range of HED science topic areas such as plasma physics, laboratory astrophysics, high-pressure materials, magnetized HED plasmas, nuclear science, and novel diagnostic development. NLUF is part of the Joint Program in High-Energy-Density Laboratory Plasmas, which is sponsored jointly by the NNSA Office of Research, Development, Test, and Evaluation and the DOE Office of Fusion Energy Sciences.

DOE funds the operation of NLUF, making it possible for researchers, including students, to conduct experiments without a direct facility charge. In addition, DOE provides research funds directly to these users for experiments. To broaden the science scope and grow the user community, NLUF will become a facility access program starting in 2022 with no restrictions on the source of the research funds users may have.

"The benefits of the NLUF Program are really tremendous," said Dr. Mingsheng Wei, an alumna of NLUF and GA who now serves as the NLUF Manager at LLE. "NLUF has compiled a strong record of excellence in HED and frontier science research

and trained over 200 Ph.D. graduate students and postdoctoral researchers." More than 60 graduate students from 18 institutions (excluding UR) are currently conducting thesis research using OMEGA, primarily through the NLUF.

GA fabricates a wide variety of targets and components and performs metrology and target assembly to support over 350 shots at the Omega Laser Facility each year for NLUF users. "There are a lot of specialized targets and materials that we aren't able to make in our lab because we don't have the machinery or the expertise here," said Dr. Carolyn Kuranz, associate professor of nuclear engineering and radiological sciences at the University of Michigan. "We depend on GA for that. For example, there are hydrodynamic instability experiments that we wouldn't be able to do or characterize without those targets and materials. And we've been able to do some really cool and exciting work because of it."

After earning her Ph.D., Dr. Wei worked as a project scientist performing HED physics research at the University of California San Diego (UCSD) including NLUF-supported research on OMEGA. She joined GA in 2010, where she continued her research in HED physics, leading several NLUF projects with experiments on OMEGA and supporting target development, before moving to LLE in 2018. Kuranz also participated in the NLUF Program and now supervises graduate students doing their own research through it.

Professor Farhat Beg at UCSD has been involved in the NLUF for more than ten years. The UCSD campus is directly across the street from GA's main campus in the Torrey Pines neighborhood of San Diego. "NLUF makes it possible for us to do stateof-the-art science that we simply cannot do on any other facility," Prof. Beg said. "The program has given us an outstanding platform and opportunities for our students and post-docs to carry out top-quality work on OMEGA. GA has been instrumental in providing us with the complex targets we need to do this."

A 2016 article¹ in *Nature Physics* co-authored by Dr. Wei and Prof. Beg serves as one example of what the NLUF has helped achieve. A team of researchers from UCSD, GA, LLE, LLNL, and several other institutions (Fig. 9) conducted a series of experiments on OMEGA using copper-doped plastic shell targets from GA to demonstrate a significant improvement in energy coupling of high-intensity, laser-produced relativistic electrons in integrated cone-in-shell fast-ignition experiments. The lead graduate student on the project, Leonard C. Jarrott from UCSD, is now a staff scientist at LLNL.



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Figure 9

This NLUF-supported team of scientists from UCSD, GA, LLE, and LLNL demonstrated a significant improvement in energy coupling of high-intensity, laser-produced relativistic electrons in integrated cone-in-shell fast-ignition experiments on OMEGA. Front row (left to right): Chris McGuffey (UCSD), Chad Mileham (LLE), and Wolfgang Theobald (LLE); middle row: Farhat Beg (UCSD), Gennady Giksel (then at LLE, currently with University of Michigan), and Mingsheng Wei (then at General Atomics, currently with LLE); back row: Leonard Charlie Jarrott (then a graduate student at UCSD, currently a Staff Scientist at LLNL), Toshinori Yabuuchi (then at UCSD, currently at SACLA, Japan), Richard Stephens (then at GA, currently retired), and Hiroshi Sawada (then at UCSD, currently at University of Nevada, Reno). (Photo: Farhat Beg.)

Professor Beg praised the record of the NLUF in preparing students for careers in HED physics and particularly the work at DOE national labs. He noted that most of his graduate students have gone on to work at LLNL, LLE, and GA. One of his former post-doc researchers, Dr. Christine Krauland, did her doctoral work at LLE with NLUF funding while at the University

of Michigan, and now works in GA's Inertial Fusion Division studying an alternative ICF scheme called "shock ignition," also under an award from NLUF (Fig. 10). "I have only good things to say about NLUF and LLE," Dr. Krauland said. "It's a critical talent pipeline for the ICF labs and the HED community as a whole."



Figure 10 Dr. Christine Krauland (photo: Christine Krauland).

Dr. Johan Frenje, assistant head of the High Energy Density Physics Division at MIT's Plasma Science and Fusion Center, had similar praise for the NLUF program. "It's everything for us," he said of the opportunity to work on OMEGA through the collaboration with GA. Frenje has guided about 15 to 20 doctoral students through the program in the 20 years he's been at MIT. "Of those," he said, "80% to 90% have gone on to work at the national labs."

"GA's participation in the NLUF Program is absolutely essential," Dr. Frenje said. "There's no question that we can't do this otherwise. We rely on GA to help us push the boundaries of science. Our requirements for targets have gotten tougher and tougher, but GA has delivered."

"We are equal partners, and it is a very collaborative process," Prof. Beg said of GA. "They provide opportunities for us to do really challenging science. They work to make the targets we need and have devoted substantial R&D money to make these projects possible."

Dr. Wei found that her experience working at GA to design targets for experiments on OMEGA gave her additional perspective on what the program has accomplished: "Working on the target side, it was really rewarding being able to bridge the gap between the physics and design and the actual experiments."

Toward the Future

"We do 10 to 15 experiments a day," Dr. Campbell said, "and maybe 3 or 4 cryogenic experiments. I'd like to do a hundred a day. That means we have to learn how to fabricate targets much more rapidly, and higher quality, and that's a real challenge. To be more relevant to NNSA and others, we need more-complicated targets made from multiple layers-multi-shell rather than single shell. That's something that GA is working on."

"Target design offers the greatest level of flexibility of all of the experimental parameters available to academic and laboratory researchers and physicists," M. Farrell said. "Enabling today's experiments while advancing the state of the art in targets for future designs is a passion for the engineers and scientists at GA. Transitioning from a "single shot" to "repetition-rated shot" capability is a grand challenge the GA and LLE collaboration is looking forward to taking on."

As much as the collaboration has accomplished, Dr. Campbell believes there is still much more to be done, given sufficient support. "I think the country does not devote enough R&D toward target fabrication," he said, "and I would like to see that grow. We have these three active facilities: Omega, the NIF, and the Z Machine at Sandia. We need to not only supply the targets for the ongoing programs, but also supply the targets for the future. That's critical to maintaining U.S. leadership in this field and ensuring we can reap the benefits of all this effort."

1. L. C. Jarrott et al., Nat. Phys. 12, 499 (2016).

A Proposal for Pentagonal Prism Spherical Hohlraum Experiments on OMEGA

W. Y. Wang and R. S. Craxton

Laboratory for Laser Energetics, University of Rochester

An important requirement for achieving ignition and gain through inertial confinement fusion is obtaining high levels of drive uniformity on a spherical capsule.^{1–3} In the indirect-drive approach,² the fuel capsule is placed inside a case, known as a hohl-raum, which is made of a high-*Z* material (typically gold) and converts the laser energy into an x-ray radiation field that provides a smooth drive on the capsule surface. On the National Ignition Facility (NIF),⁴ the laser beams enter through two laser entrance holes (LEH's) on the axis of a cylindrical hohlraum.² For future laser systems, however, spherical hohlraums have attracted recent interest as a means of achieving better uniformity.^{5–7} The work presented here explores the use of a spherical hohlraum, known as a pentagonal prism (PEPR) hohlraum, that will allow spherical hohlraums to be tested before future large-scale laser systems are constructed. The PEPR hohlraum has seven LEH's and is well suited to the OMEGA geometry. Proposed experiments on OMEGA are predicted to produce highly uniform compressions of the capsule.

The first spherical hohlraum to be proposed was the tetrahedral hohlraum,^{8–10} shown in Fig. 1(a), with four LEH's located at the vertices of a tetrahedron. Tetrahedral hohlraum experiments were performed on OMEGA, producing highly uniform capsule compressions^{11,12} consistent with the radiation drive on the capsule having less than 1% nonuniformity. Recently, octahedral hohlraums [Fig. 1(b)] were proposed as a more-uniform alternative to cylindrical and tetrahedral hohlraums, with flux nonuniformity as low as 0.1% (Refs. 5–7). The octahedral hohlraum has six LEH's corresponding to the centers of the faces of a cube or the vertices of an octahedron.



Figure 1 (a) Tetrahedral, (b) octahedral, and (c) pentagonal prism (PEPR) hohlraums. Laser entrance holes on the far side of the hohlraums are shown in outline.

Although the 60-beam OMEGA laser is geometrically unsuitable for driving octahedral hohlraums (as is also true of the NIF), the PEPR hohlraum [Fig. 1(c)] is well matched to the symmetry of the OMEGA target chamber, whose beam configuration has fivefold symmetry about the vertical axis. The LEH's of the PEPR hohlraum are based on the faces of a pentagonal prism, with five LEH's around the equator and one on each pole. This configuration was first suggested by Farmer *et al.*¹³

The PEPR hohlraum design presented here has dimensions taken from Ref. 10: the hohlraum diameter is 2800 μ m, the capsule diameter is 550 μ m, and the LEH diameter is 700 μ m. Five beams enter each of the two polar LEH's and ten beams enter each of the equatorial LEH's. Figure 2(a) shows which beams enter each LEH. The angle of incidence θ_i relative to the LEH normal ranges from 21.4° (for beams passing through the polar LEH's) to 69.7°. For comparison, θ_i ranges from 23.2° to 58.8° for the OMEGA tetrahedral hohlraum¹⁰ and from 21.2° to 52.4° for the NIF. The ray paths of the beams passing through the polar LEH's are shown in Fig. 2(b). They are focused inside the hohlraum to maximize the clearance from the capsule, as was done for the earlier tetrahedral hohlraum experiments on OMEGA. There are problems associated with the use of small angles of incidence on large systems such as the NIF, including laser–plasma instabilities along the large propagation distances and absorption in the hohlraum plasma. The proposed octahedral hohlraum avoids low values of θ_i because all beams enter in the optimal range of 50° to 60°.



Figure 2

(a) LEH assignments used in *LORE* simulations for a PEPR hohlraum on OMEGA. Colors of beam ports represent LEH assignments. The small colored circles indicate the LEH locations. (b) Ray paths of beams entering through the LEH's (6 and 7) on the z axis. These beams are focused inside the hohlraum to maximize the clearance from the capsule.

The PEPR hohlraum is analyzed using a new view-factor code *LORE*,¹⁴ which follows the physics model used in the code *BUTTERCUP*.⁹ *LORE* traces beam paths starting from the target chamber port. Each beam is divided into multiple rays, each traveling through the best-focus point of the beam. *LORE* finds the intersection of each ray with the hohlraum wall and includes an *ad hoc* model of how much energy is deposited at that point and how much is reflected to the next intersection point. Typically, all the energy is deposited at the first intersection since the hohlraum wall is strongly absorbing. Figure 3(a) shows contours of deposited energy. One can see 60 distinct laser spots, spread fairly uniformly over the hohlraum wall. As recognized in Ref. 9, this is desirable for capsule uniformity. The beam spots are all clear of the LEH's.

After tracing all the beams, *LORE* determines a spatially independent background radiation temperature T_r by assuming a Planckian radiation field in the hohlraum. T_r is calculated by balancing the power entering the radiation field (the absorbed laser power multiplied by the laser-to-radiation conversion efficiency) with the power lost to the hohlraum wall, the capsule, and the LEH's. Of particular importance is the loss to the wall, equal to $\sigma T_r^4 (1-\alpha_w) A_w$, where σ is the Stefan–Boltzmann constant, α_w is the wall albedo, and A_w is the wall area. Early in time, the albedo is low and most of the radiation incident on the wall goes into heating the wall. Later in time, the heated wall re-radiates most of its incident energy into the hohlraum and the albedo approaches unity.

Next, *LORE* calculates the emitted radiation flux at every point on the wall as the sum of the re-radiated portion of the incoming radiation ($\alpha_w \sigma T_r^4$) and the portion of the absorbed laser flux that is converted to radiation. The emitted radiation flux I_e is parametrized in terms of the effective radiation temperature T_e , defined at each point on the wall such that ($I_e = \sigma T_e^4$).

To determine the radiation uniformity on the capsule, *LORE* scans over multiple points on the capsule. For each point, *LORE* integrates the radiation flux I_e over all viewing directions. These integrals typically involve scanning over 60,000 points on the capsule and, for each point, looking along ~100,000 directions. A contour plot of the flux variations on the capsule is given in Fig. 3(b) for an albedo α_w of 0.85. The nonuniformity level is very low at 0.6% rms.

Figure 3(b) shows that, while the variations in drive on the capsule are very small, the strongest drive occurs at the poles. This is because at late times (high albedos) the heated hohlraum wall provides the dominant contribution to the drive. As a result of the two LEH's on the poles being spaced farther from other LEH's than the five equatorial LEH's, the poles of the capsule receive more drive. Conversely, at early times (low albedos), the laser-heated spots provide the dominant contribution to the drive. As a result of these spots being more "clumped" around the equator than the poles [Fig. 3(a)], the equatorial region of the capsule receives slightly more drive (but the nonuniformity is still small at 1.1% rms).



Figure 3

(a) Contour plot of deposited laser energy per unit area on the wall of the PEPR hohlraum. The LEH's are indicated in red. (b) Contour plot of radiation flux variations on the capsule for an albedo of 0.85.

The design was optimized to provide good uniformity at all albedos, i.e., good time-dependent uniformity. This was accomplished by adjusting the aim points of the beams within their LEH's to shift the laser-heated spots closer to the poles. Figure 4(a) shows the dependence of the nonuniformity on albedo for the optimized case, varying from 1.1% at low albedo to 0.6% at high albedo. To achieve this level of time-dependent uniformity requires that both contributions to the drive (laser spots and wall) produce good uniformity since they each dominate at a different time. This is hard to accomplish with cylindrical hohlraums, for which "beam phasing" (different pulse shapes in different sets of beams²) is typically required to provide the best balance between the two contributions. While this can produce a low time-averaged nonuniformity, time-dependent nonuniformity can limit the attainable fuel convergence. In the PEPR hohlraum, as in tetrahedral and octahedral hohlraums, all beams can be given the same laser temporal pulse shape.

Figure 4(a) includes, for comparison, an unoptimized PEPR hohlraum design, in which all beams are aimed through the centers of the LEH's. The 6% nonuniformity at low albedo results from the beam spots clumping closer to the equator than shown in Fig. 3(a). The nonuniformity declines at higher albedo as the wall contribution increases. Also shown in Fig. 4(a) is a prediction for optimized tetrahedral hohlraums with the same dimensions. At values of albedo below 0.5, the tetrahedral hohlraum provides better uniformity than the optimized PEPR hohlraum because the locations of deposited laser energy on the tetrahedral hohlraum are more evenly spread out. At albedos greater than 0.5, however, the nonuniformity is lower for the PEPR hohlraum.

A critical role in hohlraum design is played by the case-to-capsule ratio, i.e., the hohlraum radius divided by the capsule radius. It has long been recognized that a large ratio provides better uniformity at the expense of a lower radiation temperature.² This



Figure 4

(a) The rms nonuniformity on the capsule for optimized and unoptimized PEPR hohlraums and an optimized tetrahedral hohlraum as a function of albedo. (b) The rms nonuniformity as a function of the ratio of hohlraum radius to capsule radius for the optimized PEPR hohlraum (with the capsule and LEH radii held fixed), illustrating the tradeoff between uniformity and radiation temperature. The albedo here is 0.8.

tradeoff is shown in Fig. 4(b) for the PEPR hohlraum, where the hohlraum radius is varied with the capsule and LEH radii held fixed. The point at a ratio of 5.09 corresponds to the design used in this article, with $T_r = 195$ eV. This can be increased to 215 eV at a ratio of 3.5 at the expense of a greater nonuniformity of 2%. These values of T_r are limited by the 18-TW laser power assumed here (approximately the peak power used for the OMEGA tetrahedral hohlraum experiments). A NIF-scale PEPR design predicts $T_r = 293$ eV at a ratio of 3.5 with a nonuniformity of 1.23% (Ref. 14).

An ignition-scale laser system irradiating an octahedral hohlraum would have lower nonuniformity than a PEPR hohlraum for a given case-to-capsule ratio because of the better geometrical symmetry. It would also benefit from more-favorable beam paths because of the elimination of small angles of incidence. While not a candidate for an ignition system, the PEPR hohlraum has the advantage that it can be used on an existing facility, offering a platform for the performance of a variety of experiments. It can be used to demonstrate high-quality spherical implosions using minimal tuning compared with the NIF. Beam phasing is not required: all beams use the same temporal pulse shape. The ratio of hohlraum-to-capsule radius may be adjusted to explore the trade-off between capsule uniformity and background radiation temperature. In addition, the anticipated ease with which near-symmetric implosions can be generated offers a platform for the examination of hot-spot physics and the development of improved diagnostics.

While the PEPR hohlraum promises to drive implosions that are substantially symmetric and 1-D, the geometry is inherently 3-D, requiring 3-D simulations for detailed hohlraum design. Many years ago, the difficulty of carrying out 3-D simulations may have favored the selection of cylindrical hohlraums, which, in spite of their uniformity issues, are well suited to 2-D modeling. The PEPR platform on OMEGA can provide a useful test bed for 3-D modeling.

Further information on this work can be found in Ref. 14.

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Impact of Areal-Density Asymmetries on the Loss of Confinement and Ignition Threshold in Inertial Confinement Fusion Capsules

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In inertial confinement fusion implosion experiments, variations in the shell areal density ρR reduce the shell's inertia to confine the core pressure. Distorted capsules with large areal-density modulations decompress faster than uniform capsules in the disassembly phase. In this summary, a simple 3-D analytic hot-spot model is derived to include the effects of low-mode arealdensity modulations in the ignition criterion. The generalized 3-D ignition criterion for low modes is shown to depend on both the harmonic mean and the arithmetic mean of the areal density. The "thin spots" in the shell are shown to dominate the loss of confinement as reflected by the harmonic mean definition of areal densities.

Effects of low-mode asymmetries on confinement can be described by a highly idealized ignition model. In this analysis, the 3-D hot-spot energetics in the disassembly phase are studied by a simple form of the 3-D hot-spot energy equation [Eq. (9) in Ref. 1]

$$\frac{\mathrm{d}}{\mathrm{d}t}\ln(PV^{\gamma}) = P/S_{\mathrm{f}}.\tag{1}$$

Here $S_f \equiv 24 T_i^2 / (\langle \sigma v \rangle_{DT} E_{\alpha})$ is approximated with a constant, valid for DT fusion reactivities $\langle \sigma v \rangle_{DT}$ within the ion-temperature range 6 < $T_i < 20$ keV; $S_f \simeq 7.24$ atm s is the minimum ignition threshold;^{2,3} $\gamma = 5/3$ is the ratio of specific heat for an ideal gas; and $E_{\alpha} = 3.5$ MeV is the DT fusion alpha particle's initial kinetic energy; heat conduction loss and radiation loss are neglected. The time evolution of the hot-spot volume *V* is approximated by a second-order expansion of *V* in time *t* after the capsule is compressed to the minimum hot-spot volume. For igniting capsules, the hot-spot pressure *P* grows rapidly as the rate of alpha heating exceeds the rate of PdV work in the disassembly phase. As a result, the generalized 3-D ignition threshold is obtained, $\chi_{3-D} \equiv P_s \tau_c / S_f \rightarrow 1$ as the hot-spot volume V_s . The impact of areal-density asymmetries on the loss of confinement is hidden in the second time derivative of the hot-spot volume. It can be shown that the confinement time depends on the harmonic mean (HM) definition of areal densities:

$$\tau_{\rm c} \approx \sqrt{\frac{V_{\rm s}}{P_{\rm s}A_{\rm s}}} \langle \rho R \rangle_{\rm HM},\tag{2}$$

where A_s is the hot-spot surface area at stagnation. In the harmonic mean definition, the areal densities of thinner regions are weighted more than that of thicker regions. As a result, in the presence of low modes, the hot spot disassembles faster, leading to a shorter confinement time τ_c with respect to 1-D since $\langle \rho R \rangle_{HM} < \rho R_{1-D}$. The harmonic mean is the only way to account for "thin spots" (areas of ultralow areal density) in the shell that dominate the loss of confinement. If localized, thin spots do not significantly contribute to the arithmetic mean of the areal density. Measuring only the arithmetic mean would overestimate the confinement and therefore the Lawson parameter in the presence of severe asymmetries, leading to areas of ultralow areal density.

As shown in Fig. 1(a), the no- α Lawson criterion $\chi_{no\alpha}^{3-D} = P_s \tau_c / S_f$ is shown to capture the onset of ignition for low modes $\ell = 1$ to 2 when $\chi_{no\alpha}^{3-D} \approx 1$. The confinement time τ_c is calculated from the curvature of the temporal hot-spot volume at stagnation as the hot spot is compressed to the minimum volume. The yield amplification is shown to be well approximated by the fitting

formula $Y_{\text{amp}}^{3-\text{D}} \approx (1 - \chi_{\text{no}\alpha}^{3-\text{D}}/0.96)^{-1.14}$. In Fig. 1(b), the harmonic mean ρR of mode $\ell = 1$ is shown to degrade faster with the ion-temperature ratio $T_{\text{min}}/T_{\text{max}}$ than that of the arithmetic-mean ρR . This result implies that the 1-D values of areal densities can be inferred from the measured ion-temperature measurement asymmetry. Consequently, the degradation of the 3-D Lawson criterion $\chi_{\text{no}\alpha}^{3-\text{D}}$ for $\ell = 1$ asymmetries can be assessed as a unique function of the ion-temperature measurement asymmetry.



Figure 1

(a) Yield amplification as a function of the no- α Lawson criterion $\chi^{3-D}_{no\alpha}$ using the confinement time τ_c defined by the second time derivative of the hot-spot volume at the time when the capsule is compressed to the minimum volume. (b) The degradation of harmonic-mean and arithmetic-mean areal densities against the ion-temperature measurement asymmetry.

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Microcoulomb (0.7±0.4/0.2-µC) Laser-Plamsa Accelerator on OMEGA EP

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Laser-plasma accelerators (LPA's) driven by short-pulse, kilojoule-class lasers provide a path to producing compact sources of high-charge, high-energy electron beams for conversion into x-ray and positron sources. Here, we report on the first LPA driven by a short-pulse, kJ-class laser (OMEGA EP) connected to a multikilojoule high-energy-density science (HEDS) driver (OMEGA).

Experiments were performed on the OMEGA EP Laser System.¹ The laser was run with a central wavelength λ of 1054 nm at best compression (pulse duration of 700±100 fs). To improve the quality of the focal spot and increase the Rayleigh length, the focusing geometry of the short-pulse laser beams was converted from its nominal f/2 geometry by using spatially filtered apodizers² located at the injection plane before amplification in the Nd:glass beamline to control the beam diameter and generate an f/5, f/6, f/8, or f/10 geometry. At focus, the R_{80} spot size of the laser (i.e., radius that contains 80% of the total energy) was between 11.5 and 19.9 μ m. The apodized laser energy varied from 10 to 115 J, which produced on-target peak normalized vector potentials [$\alpha_0 \cong 8.6 \times 10^{-10} \sqrt{I_0 (W/cm^2)} \lambda(\mu m)$], where I_0 is the vacuum intensity, between 1.8 and 6.7. The apodized laser pulse was focused 500 μ m inside a Mach 5 gas jet with nozzle diameters varying between 2 and 10 mm as shown in Fig. 1(c). The gas was 100% He, and the resultant plasma densities in the plateau ranged from 1.5 × 10¹⁸ to 4.5 × 10¹⁹ cm⁻³, depending on nozzle diameter and backing pressure.

Figure 2(a) shows that the total charge in the electron beams scales approximately linearly with a_0 . The data shown are for a 6-mm-diam nozzle operating at a plasma density of 5×10^{18} cm⁻³, but plasma densities of 1, 2, and 3×10^{19} cm⁻³ showed the same trend. This trend was also seen for 4-mm-diam nozzles operating at 1×10^{19} cm⁻³ and 10-mm-diam nozzles at densities of 0.2, 0.5, 1, and 3.5×10^{19} cm⁻³. The charge in the electron beams was calculated using the method described in Ref. 3.

Figure 2(b) shows that the charge in the electron beam scales approximately linearly with plasma density until a density of 1×10^{19} cm⁻³. The two data sets shown each have a different a_0 value; the rate of increase of charge with plasma density is steeper for the higher a_0 value. The highest-charge electron beam measured in this experiment, which had a charge of $707\pm429/224$ nC, was produced at an a_0 of 6.6 and a plasma density of 7.5×10^{18} cm⁻³. Using an electron energy of 17.9 MeV, which is the weighted average electron energy of the representative electron spectrum from this experiment [Fig. 1(d)], this charge corresponds to a conversion efficiency from laser energy to electron energy of 11%. The details of this calculation can be found in Ref. 3. Of that total energy, 30%, 50%, and 90% is contained in electrons with energies below 18.5 MeV, 25.6 MeV, and 85.1 MeV, respectively. Figure 2(c) shows that when the charge scaling was extended to higher plasma densities, the maximum charge produced plateaus with density. A similar trend was seen for data taken on a 6-mm-diam nozzle for both a_0 values of 5 and 6.



Figure 1

[(a),(b)] Examples of a target spot at the focal plane for the standard f/2 focus and the f/6 apodized focus, respectively. The peak fluence per energy for (a) and (b) is 7.9 and 11.8×10^5 cm⁻², respectively. (c) Relative layout of the laser, target, and diagnostics. (d) Electron spectrum from an $a_0 = 5.1$ laser shot propagating through a plasma density of 5.4×10^{18} cm⁻³ generated by a 6-mm-diam nozzle. The shaded region marks the detection limit of the electron–positron–proton spectrometer (EPPS). OAP: off-axis parabola.



Figure 2

(a) Electron-beam charge versus a_0 for a 6-mm-diam nozzle operating at a plasma density of 5×10^{18} cm⁻³. Electron-beam charge as a function of plasma density (b) up to $\sim 1 \times 10^{19}$ cm⁻³ for $a_0 \sim 3$ (magenta circles) and $a_0 \propto 6$ (blue squares) for a 6-mm-diam nozzle and (c) over the entire sampled plasma density range for $a_0 \sim 5$ and a 10-mm-diam nozzle. The dashed lines are added to guide the eye.

A microcoulomb-class, high-conversion-efficiency laser-plasma accelerator was demonstrated, providing the first laser-plasma accelerator driven by a short-pulse, kJ-class laser (OMEGA EP) connected to a multi-kJ HEDS driver (OMEGA). The produced electron beams have maximum energies that exceed 200 MeV, divergences as low as 32 mrad, record-setting charges that exceed 700 nC, and laser-to-electron conversion efficiencies up to 11%. The total charge in the electron beam is found to scale with both a_0 and plasma density. Based on these empirical scalings, higher-charge electron beams may be possible using laser systems that can deliver a_0 values larger than the maximum a_0 of 6.7 produced in this configuration while still maintaining longer f numbers and near-Gaussian, single-moded laser spots on target.

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Cross-Beam Energy Transfer Saturation by Ion-Trapping–Induced Detuning

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The performance of direct-drive inertial confinement fusion (ICF) implosions relies critically on the coupling of laser energy to the target plasma. Cross-beam energy transfer (CBET), the resonant exchange of energy between intersecting laser beams mediated by ponderomotively driven ion-acoustic waves (IAW's), inhibits this coupling by scattering light into unwanted directions. The variety of beam intersection angles and varying plasma conditions in an implosion results in IAW's with a range of phase velocities. Here, we show that CBET saturates through a resonance detuning that depends on the IAW phase velocity and that results from trapping-induced modifications to the ion distribution functions. For smaller phase velocities, the modifications to the distribution functions can rapidly thermalize in the presence of mid-Z ions, leading to a blue shift in the resonant frequency. For larger phase velocities, the modifications can persist, leading to a red shift in the resonant frequency. Ultimately, these results may reveal pathways toward CBET mitigation and inform reduced models for radiation-hydrodynamic codes to improve their predictive capability.

Laser-driven ICF experiments are subject to numerous nonlinear couplings between the electromagnetic waves and plasma waves. Among these couplings, CBET, the resonant exchange of energy between intersecting laser pulses mediated by ponderomotively driven IAW's, has emerged as particularly troublesome.¹ CBET inhibits the performance of both direct- and indirect-drive implosions by scattering light into unwanted directions.^{2–4} In direct drive, this reduces the coupling of laser energy to the capsule, while in indirect drive, it can spoil the symmetry of the x-ray illumination. Both approaches have achieved some success in mitigating CBET by using independent wavelength shifts on the beams to detune the interaction.^{5–8} More-extensive mitigation, however, requires pulses with a much larger bandwidth⁹—a technology in active development at LLE and the Naval Research Laboratory.^{10–12}

Comprehensive, predictive models of CBET can guide both ongoing and future mitigation strategies and help define the expanded ICF design space that these strategies afford. Current integrated models of ICF implosions, using radiation-hydrodynamic simulations, typically implement simple linear models of CBET due to the computational expense associated with more-complete models. While more-sophisticated models have been developed,^{13–16} common approximations include ray optics (i.e., speckle and diffractive effects are ignored) and a steady-state plasma response, while neglecting nonlinear processes.¹⁷ This is in spite of a mounting body of work pointing to the importance of processes such as ion trapping, stochastic heating, two-ion decay, nonlinear sound waves, and IAW breakup.^{18–23} Perhaps the most-convincing indication comes from a number of experiments that have observed nonlinear saturation.^{24–26} The most recent of these experiments, performed on the OMEGA laser, demonstrated that, at high intensities, a drop in the power transferred from a pump to seed pulse was accompanied by an ~7× increase in the ion temperature.²⁶

Motivated by this observation, this work provides a detailed description of the underlying physics responsible for CBET saturation for conditions relevant to these experiments. Specifically, we show that depending on the phase velocity of the IAW (v_p) , CBET can saturate through two types of resonance detuning, both of which result from trapping-induced modifications to the ion distribution function. For "small" IAW phase velocities, the modifications to the distribution function rapidly thermalize in the presence of mid-*Z* ions, blue shifting the resonant IAW frequency. For "large" IAW phase velocities, the modifications to the

distribution function persist for a longer time and red shift the resonant frequency. These results, obtained using the collisional particle-in-cell code VPIC,²⁷ avoid many of the pitfalls associated with the reduced models used in radiation-hydrodynamic codes and provide insight into the evolution and feedback of CBET with the coronal plasma.

Figure 1 demonstrates that, in both the small and large v_p cases, CBET evolves through three stages: an initial linear stage (<5 ps), an early saturation stage (~5 to 20 ps), and a final late-time saturation stage (>20 ps). During each of these stages, the gain, i.e., $G = \ln(P^{\text{out}}/P^{\text{in}})$, where P^{in} and P^{out} are the probe input and output powers, tracks the electrostatic energy. The initial stage corresponds to transient growth of the IAW as the interaction attempts to evolve toward a linear, steady state.



Figure 1

Evolution of the CBET gain and electrostatic energy for the (a) "small" and (b) "large" phase velocity cases. The CBET gain generally tracks the electrostatic energy. Due to an interplay of IAW transverse breakup and reduced Landau damping from ion trapping, the gain saturates at a value lower than the linear, steady-state gain.

Before this state can be reached, however, the interaction becomes nonlinear and the IAW undergoes transverse breakup.²⁸ The IAW initially exhibits coherent phase fronts, but after 20 ps, the fronts have broken up into smaller transverse structures. Due to its observed correlation with ion trapping, the breakup likely results from the trapped particle modulational instability (TPMI).^{28–31} In the TPMI, the nonlinear frequency shift from ion trapping with inhomogeneity in the IAW amplitude creates variations in the phase velocity across the phase fronts. If a section of the phase front advances or retards by more than $\sim \pi/2$ with respect to adjacent sections, the front breaks. At this point, the wave amplitude crashes and the energy is transferred to the ions. The local dissipation of the wave prevents additional trapping and changes to the phase velocity. In fact, the rapid drop in electrostatic energy after ~10 ps results from initially trapped ions carrying away the energy of the now-broken IAW.

The rapid decay of the electrostatic energy (Fig. 1) is followed by a slow drop in the gain and marks the beginning of the latetime saturation stage. During this stage, the small and large v_p interactions exhibit strikingly different behaviors. Foremost, the gain drops substantially for small v_p and only modestly for large v_p . While both trends have their origin in ion-trapping-induced detuning, the cause of this detuning depends on the role of each ion species in collisional energy transfer and thermalization.

After ~50 ps, the plasma conditions and gain evolve slowly enough that CBET occurs in a quasi-steady state (Fig. 1). In this quasi-steady state, the kinetic coupling coefficient Im[Γ], calculated using the electron and ion distribution function and averaged over the interaction region, provides the response and resonant behavior of the plasma. For small v_p , the thermalization of the H and N ions causes a gradual blue shift in the resonant IAW frequency [Fig. 2(a)]. With the fixed wavelength shift of the probe beam, the gradual blue shift in resonant frequency causes the gain to drop. In addition, the increased damping from the modified distribution function has broadened the resonance peak and lowered its maximum. For large v_p all of the collision rates are generally lower than in the small- v_p case because of the larger phase velocity. This allows the trapping-induced modifications to the distribution function to persist for a much longer time. Consistent with the trapped ion frequency shift,^{32,33} these modifications cause a red shift in the resonant frequency [Fig. 2(b)].



Figure 2

The Im[Γ], which determines the gain in the steady-state approximation, calculated using the electron- and ion-velocity distribution functions from VPIC at t = 0 and 80 ps. Initially, the IAW is driven on resonance for both the (a) small- (b) and large- v_p cases (intersection of the dashed black lines and the peak of the solid curve). Later, (a) the increase in ion temperatures and flow blueshift the peak in the small- v_p case, while (b) the persistent tails of the distribution function due to trapping redshift the peak in the large- v_p case. Both detunings cause the gain to drop.

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Magnetically Collimated Relativistic Charge-Neutral Electron–Positron Beams from High-Power Lasers

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Relativistic electron–positron pair plasmas are important objects for study in fundamental plasma physics; they have unique properties resulting from mass symmetry and are used to explain observations of multiple astrophysical phenomena, such as gamma-ray bursts, black holes, and active galactic nuclei.^{1–7} The creation of dense, relativistic pair plasmas in the laboratory has remained elusive due to the short positron lifetime and difficulty in producing plasmas of high-enough density.⁸ To create a useful pair plasma for laboratory astrophysics, three conditions must be satisfied: (1) high-energy (MeV) particles must be confined longer than the plasma time scale of interest; (2) the plasma dimensions must be significantly greater than the Debye length; and (3) the plasma must be charge neutral. To date, there are limited experiments that produce high-temperature pair plasmas and high-flux pair jets required to simulate astrophysical phenomena.^{9,10}

One method to generate energetic positron–electron pairs relevant to laboratory astrophysics experiments is to use ultrashort high-intensity lasers. Previous experiments have produced jets of positrons and electrons with MeV energies in a small volume (<3 mm³), with particle densities of about 10^{15} cm⁻³ and 10^{13} cm⁻³ at the source for electrons and positrons, respectively.^{11,12} These positrons are generated through the Bethe–Heitler process, whereby bremsstrahlung γ rays from high-energy electrons decay into electrons and positrons. So far, of the three plasma conditions, the confinement condition (2) has been achieved experimentally using an imposed magnetic field,¹² and the neutrality condition (3) has been achieved in another experiment using a very thick, high-*Z* target to generate equal numbers of electrons and positrons with >5-MeV energy.¹³ No previous experiment has produced a high-temperature, neutral pair-plasma beam; many experiments have produced jets with electron density exceeding the positron density by severalfold but have not confined or collimated them for long periods.

Magnetic focusing has provided a promising path toward satisfying the conditions of charge neutrality and MeV confinement. Indeed, previous experiments demonstrated the use of an externally imposed magnetic field to collimate electrons and positrons for measurement.^{14,15} These experiments measured nearly 70× as many positrons on magnetized shots as unmagnetized shots due to collimation by the magnetic field. Ratios of electrons to positrons ranged from 10 to 3 for various shots during the experiment, but did not reach charge neutrality. This improvement suggests that externally imposed magnetic fields have great potential to collimate fully neutral pair plasmas. Here we report the measurement of a neutral, collimated electron–positron beam by utilizing recent upgrades to the pulsed-power system known as MIFEDS (magneto-inertial fusion electrical discharge system).¹⁶ This represents a significant step toward generating charge-neutral electron–positron pair-plasma jets in the laboratory.

Five magnetized shots and one unmagnetized reference shot were carried out in the experiment on the OMEGA EP laser, which used 10-ps pulses with $\lambda = 1054$ -nm light, focused on a 500- μ m-diam, 20- μ m-thick gold disk. The laser energy was 900±20 J (intensity = 9 × 10¹⁸ W/cm²). It should be noted, however, that MIFEDS-related debris and copper deposition on laser optics caused an estimated loss of energy and intensity of 15% to 20% after the first magnetized shot, leading to on-target energies in the range of 770±40 J with an intensity of 7±1 × 10¹⁸ W/cm². At these conditions, the positron yield can be estimated from Myatt *et al.*¹⁷ Assuming a laser energy of about 800 J and a laser-to-electron conversion efficiency of 30%,¹⁸ the expected positron yield is about

 7.5×10^{10} in total, with a density of approximately 3×10^{10} positrons/sr since without collimation, the positrons diverge with an angle of 1 to 2 sr (Ref. 16). The combination of positron number and divergence results in a measurement that is barely above the detection threshold of the spectrometer.

In order to magnetize shots, field-generating coils were set up in a magnetic mirror configuration with wire loops 14 mm apart with an inner diameter of 10 mm. These produced a field of up to 13 T at the ends of the mirror and 5 T in the center near the source. For two of the five shots, the MIFEDS charge voltage was tuned to examine a lower field of approximately 4 T and 9 T at the center and edges, respectively. The energy distributions of electrons and positrons were measured with an electron–positron–proton spectrometer (EPPS), which was placed along the primary magnetic-field axis. The magnetic field axis and spectrometer collection region were aligned perpendicular to the laser and target axis.

Initial particle-tracing simulations, carried out with the multiphysics code COMSOL, show that when the field is imposed, particles with energy less than 2.5 MeV are completely confined radially. For higher-energy particles, those with a small initial pitch angle relative to the mirror field axis will be well focused and collimated by the magnetic field due to slight deflections. Particles of roughly 13 MeV from a point source 8 mm away from the coil would be collimated, given the coil radius and peak magnetic-field profile used on the experiment. For particles with energy higher or lower than this optimum energy, particles will be overfocused prior to reaching the spectrometer or not focused at all.

These simulation predictions were borne out in the experiment, which consisted of five magnetized shots (three at high field and two at low field) and one unmagnetized reference shot. In the no-field case, almost no positrons were measured above noise level (as expected) and electrons followed a single temperature spectrum across all energies of interest. When the field was applied, positrons of equal energy and number to the electrons were measured. The positron signal peaked at 13 MeV (Fig. 1), which matches the estimated best focus from the initial simulations. Two shots were also performed using a weaker, 9-T field, which decreased the peak energy of focused particles to 10 MeV, while still maintaining an even ratio between positrons and electrons. These collimated beams of pair particles were maintained over the 50-cm distance between the experiment and the spectrometer.



Figure 1

Electron and positron spectra measured using the field axis EPPS. With no field (dotted curves), only electrons characterized by a single slope temperature were measured along this direction. When the 13-T collimating field was applied, nearly identical quantities of electrons and positrons peaked at around 13-MeV energy were measured.

While these measurements of charge-neutral, focused electron–positron beams represent a step forward toward a true pair plasma, the densities are still insufficient for many astrophysics purposes. To effectively create scaled astrophysical shocks within the pair plasma, densities of at least 10^{13} cm⁻³ and energies of over 10 MeV are required. While our jet contains energies of 13 MeV and is collimated over a long distance (>50 cm), the density is considerably lower—roughly 2×10^8 cm⁻³ based on the assumption that the particles are accelerated throughout the laser pulse. To increase density toward meeting the confined pair-plasma goals, future work will concentrate on increasing the density of the pair-plasma beam. Thicker gold or microstructured targets can be used in

the future (similar to those used in other pair-plasma–generating experiments¹⁴), which will substantially increase the conversion to positrons. Further upgrades to MIFEDS will also be applied in order to increase the energy range for confined or collimated particles.

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Enhancing Positron Production Using Front-Surface Target Structures

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Electron–positron pair plasmas generated by laser–solid interactions offer a wide range of potential applications in different fields, including astrophysics, material science, biology, etc. We report a target design that produced a substantial gain in relativistic electron–positron pair production using high-intensity lasers and targets with large-scale microstructures on their surface. Compared to an unstructured target, a selected Si microwire array target yielded a near-100% increase in the laser-to-positron conversion efficiency and produced a 10-MeV increase in the average emitted positron energy under nominally the same experimental conditions.

The experiment was performed on the OMEGA EP Laser System. A schematic diagram of the experimental setup is shown in Fig. 1(a). The target was irradiated using a short-pulse laser with a wavelength of 1.053 μ m, an energy of 500 J, and a pulse length of approximately 700 fs. The peak intensity was estimated to be 4.5×10^{20} W/cm². Figures 1(b) and 1(c) show scanning electron microscope (SEM) images of two different target structures used in the experiment. Structure 1 was optimized through particle-in-cell (PIC) simulations of the hot-electron temperature prior to the experiment. For comparison, structure 2 showed detrimental effects on electron energies in simulations. The structures were made of Si microwires and were embedded in a thin plastic layer and then glued to a 1-mm-thick Au backing layer. The high-energy electrons generated and guided by the surface structures transport through the thick Au layer and induce pair production. The positron and electron spectra were measured by a spectrometer on the back side of the target along the laser direction (which was also the target normal direction).



Figure 1

(a) Schematic diagram of the experimental setup; (b) SEM image of the pre-optimized target structure 1; (c) SEM image of the unoptimized structure 2. PDMS: polydimethylsiloxane.

The measured spectra for both types of structured targets as well as a flat unstructured target are shown in Fig. 2. Target structure 1 generated about 50% more positrons than the regular flat target, and its laser-to-positron conversion efficiency increased by \sim 97%. The spectrum peak also shifted from \sim 50 MeV for the flat target to \sim 60 MeV for structure 1. Structure 2 showed fewer as well as much lower energy positrons, in accordance with expectations since the length and spacing of the microwires encumber the laser focusing. The electron spectrum from structure 2 also showed the same trend, in agreement with the positron measurements. The electron spectra from the flat unstructured target and from structure 1 were mutually similar, however, with both having an electron temperature of about 21 MeV.



Experimentally measured spectra for (a) positrons and (b) electrons. Different colors indicate the results from different targets under the same laser conditions.

We have performed multiple simulations to model the entire process and explain the observed phenomena. We adopted a two-stage approach since the laser–plasma interaction was modeled with a 2-D Cartesian PIC simulation and the electron transport and pair production process was modeled with a 2-D cylindrical simulation. The simulations successfully reproduced the experimental results. Compared to a flat target, structure 1 generated more high-energy (tens to hundreds of eV) electrons. Two acceleration mechanisms are responsible for these electrons, including the loop-injected direct acceleration,¹ which is associated with any targets having moderate scale-length pre-plasma, and the structure-guided direct laser acceleration, which occurs only with the structured target.² We have also found strong Weibel instabilities near the critical-density surface for both target types, which largely widens the electron divergence, explaining why the electron spectra measured at 0° look similar for both the flat and the structure 1 target. The energy of positrons is largely determined by the sheath field on the back side of the target. The simulations suggested that the integrated sheath voltage for structure 1 is about 10 MV higher than for the flat target, which is consistent with the measured energy difference between their positron peaks.

In summary, front-surface target structures have been shown experimentally to substantially enhance the positron yield and energy. The follow-up simulations explain the entire process of how the laser–plasma interaction that is manipulated by the target structure affects the yield and energy of positrons. The agreement between the simulated and experimental spectra indicates the possibility of further target optimization using two-stage PIC simulations.

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Optical Thomson Scattering

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Introduction

Thomson scattering provides a direct observation of electron motion in a plasma by encoding their velocities on the frequency spectrum of the scattered light. By propagating a beam of photons (ω_0 , \mathbf{k}_0) through a plasma and isolating the Thomson-scattering volume collected into a spectrometer (Fig. 1), a spatially resolved measurement of the plasma conditions can be determined from the scattered frequency spectrum (ω_s , \mathbf{k}_s) (Ref. 1). The scattered-power spectrum observed by the detector is given by

$$\frac{\mathrm{d}P_{\mathrm{s}}}{\mathrm{d}\omega_{\mathrm{s}}} = \frac{P_{\mathrm{i}}r_{0}^{2}L\mathrm{d}\Omega}{2\pi} \Big(1 + \frac{2\omega}{\omega_{0}}\Big)n_{\mathrm{e}}S(\mathbf{k},\omega),\tag{1}$$

where $r_0^2 = 7.95 \times 10^{-26}$ cm is the classical electron radius, *L* is the length of the scattering volume along the probe beam, $\mathbf{k} = \mathbf{k}_s - \mathbf{k}_0$, $\omega = \omega_s - \omega_0$, d Ω is the solid angle of the collected scattered photons, and P_i is the average incident laser power. The density fluctuations of the plasma around its average density dictate the primary shape of the scattered spectrum through the dynamic structure factor. For a collisionless plasma with no magnetic fields affecting the motion of the particles, the dynamic structure factor is

$$S(\mathbf{k},\omega) = \frac{2\pi}{k} \left| 1 - \frac{\chi_{e}}{\varepsilon} \right|^{2} f_{e}\left(\frac{\omega}{k}\right) + \sum_{j} \frac{2\pi}{k} \frac{Z_{j}^{2} n_{j}}{n_{e}} \left| \frac{\chi_{e}}{\varepsilon} \right|^{2} f_{j}\left(\frac{\omega}{k}\right), \tag{2}$$

where f_e and f_j are the normalized 1-D electron and ion-velocity distribution functions, respectively, projected along the scattering vector (**k**), Z_j is the average charge of the *j*th ion species, $n_e = \sum_j n_j Z_j$, and n_j is the density of *j*th ion species. The longitudinal dielectric function is

$$\mathcal{E} = 1 + \chi_e + \sum_j \chi_j,\tag{3}$$

where the kinetic plasma susceptibilities are given by

$$\chi_{\rm e}(\mathbf{k},\omega) = \frac{4\pi e^2 n_{\rm e}}{m_{\rm e}k^2} \int_{-\infty}^{\infty} \mathrm{d}\mathbf{v} \frac{\mathbf{k} \cdot \partial f_{\rm e}/\partial \mathbf{v}}{\omega - \mathbf{k} \cdot \mathbf{v}},\tag{4}$$

$$\chi_{j}(\mathbf{k},\omega) = \frac{4\pi Z_{j}^{2} e^{2} n_{j}}{m_{j} k^{2}} \int_{-\infty}^{\infty} \mathbf{d} \mathbf{v} \frac{\mathbf{k} \cdot \partial f_{i} / \partial \mathbf{v}}{\omega - \mathbf{k} \cdot \mathbf{v}},$$
(5)



Figure 1

A typical Thomson-scattering system is shown where the probe laser beam propagates through a distributed phase plate (DPP) before being focused to an area (A) at the Thomson-scattering volume. An aperture stop is imaged into the plasma to define the Thomson-scattering volume along the propagation of the probe beam (L).

The scattering spectrum can be used to measure the electron distribution function, which is most evident in the high-frequency noncollective Thomson-scattering regime. Here the collective motion of the electrons is heavily damped and the power scattered at a particular frequency is proportional to the number of electrons with a velocity that Doppler shifts the frequency of the probe laser to the measured frequency [Fig. 2(a)]. In this strong damping regime, where the scattering parameter $\alpha \equiv 1/k\lambda_{De} \ll 1$ ($\lambda_{De}^2 = \kappa T_e/4\pi e^2 n_e$ is the electron Debye length), Eq. (1) is reduced to light that is scattered from an ensemble of uncorrelated electrons,

$$\frac{\mathrm{d}P_{\mathrm{s}}}{\mathrm{d}\omega_{\mathrm{s}}} = \frac{P_{\mathrm{i}}r_{0}^{2}L\mathrm{d}\Omega}{2\pi} \Big(1 + \frac{2\omega}{\omega_{0}}\Big) n_{\mathrm{e}}f_{\mathrm{e}}(\omega/k). \tag{6}$$

From here it is evident that the noncollective spectrum provides a direct measurement of the electron distribution function, but in practice, the small scattering cross section of the electron and the small number of electrons at high velocities lead to low signal-to-noise, typically limiting this technique to measuring electrons in the bulk of the distribution function.

It is also possible to measure the electron distribution function in the regime where the high-frequency scattering spectrum is governed by the collective electron motion introduced by weaker damping of the density fluctuations.^{2,3} In this collective regime, the thermal particle motion drives a rich spectrum of fluctuations, which when probed, can present themselves in the scattering spectrum as peaks shifted around the incident frequency of the laser (Fig. 2). As charged particles propagate through the plasma, they leave electrostatic fluctuations in their wake. The amplitude of each fluctuation is determined by the balance of its damping rate with the rate at which it is driven by the plasma particles. The high-frequency electron plasma wave fluctuations start to play an important role in the scattering spectrum when $\alpha \sim 2$, but when the fluctuations are more weakly damped ($\alpha \ge 2$), the resonant features have separated clearly from the noncollective scattering spectrum [Fig. 2(a)]. For the low-frequency fluctuations [Fig. 2(b)], there are similar regimes, but related to the ion motion. The transition between the collective and noncollective regime in a collisionless plasma is governed by the ion Landau damping, and the collective low-frequency regimes occur for the scattering parameter $\alpha > (ZT_e/3T_i - 1)^{-1/2}$. For $ZT_e/3T_i < 3$, the fluctuations are heavily damped by the ions.

The frequency of these resonant peaks can be found approximately by solving the dispersion relation ($\varepsilon = 0$) for the natural modes of the plasma by finding the real parts of the roots of the dielectric function [Eq. (3)], which is where one can see the power of collective Thomson scattering in determining the plasma conditions. Assuming Maxwellian electron distribution functions and weakly



Figure 2

(a) High-frequency spectrum calculated from Eq. (1) in the heavily damped noncollective regime, $\alpha = 0.25$ (red dotted curve); mildly damped collective regime, $\alpha = 2.0$ (black dashed curve); and weakly damped collective regime, $\alpha = 4.0$ (blue solid curve). The temperature was maintained at $T_e = 100$ eV and the density was scaled $n_e = 1 \times 10^{17}$ cm⁻³ (red), $n_e = 6 \times 10^{18}$ cm⁻³ (black), $n_e = 2.5 \times 10^{19}$ cm⁻³ (blue). The low-frequency spectrum has been suppressed. (b) The low-frequency spectrum calculated from Eq. (1) in the heavily damped noncollective regime, $ZT_e/T_i = 0.5$ (red); mildly damped collective regime $ZT_e/T_i = 3.5$ (black); and weakly damped collective regime $ZT_e/T_i = 10$ (blue). The scattering parameters $\alpha = 2$ and $T_e/T_i = 10$ were held constant. For all calculations, the angle between incident and scattered light was held constant ($\theta = 90^\circ$).

damped fluctuations, the real part of the dispersion relation for the ion-acoustic waves simplifies to $\omega_{iaw} \simeq k \sqrt{[(Z\kappa T_e + 3\kappa T_i)/m_i]}$, in the low-frequency spectrum, and the real part of the high-frequency part of the dispersion relation, corresponding to the electron plasma waves simplifies to $\omega_{epw}^2 = \omega_{pe}^2 + 3\kappa T_e k^2/m_e$, where $\omega_{pe}^2 = 4\pi n_e e^2/m_e$ is the electron plasma frequency. Thomsonscattered light from these collective electron motions generates constructive interference at the detector plane, and the frequency of this interference can be directly related to the plasma conditions through the plasma dispersion relations; note that measuring the difference between the frequency of the laser and the peak features in the spectrum ($\omega = \omega_s - \omega_0 = \Delta \omega$) is a measure of the plasma conditions through the associated dispersion relations ($\Delta \omega/\omega_0 \approx \Delta \lambda/\lambda_0$).

Collective Thomson scattering is a powerful diagnostic regime used to overcome background radiation because of the need to resolve only the frequencies of the spectral peaks. This is in contrast with the noncollective regime, where the shape of the scattering spectrum is used to infer the plasma conditions, therefore challenging one to understand the background radiation spectrum and the wavelength sensitivity of the diagnostic. In practice, modern collective Thomson-scattering systems can resolve the complete spectrum, providing detailed measurements of the electron distribution functions,⁴ electron temperatures, ion temperatures, ^{5,6} plasma flow velocities, and electron densities.^{7,8}

Laser Beam Propagation

The small Thomson-scattering cross section is one of the most-challenging aspects of Thomson scattering. Integrating Eq. (1) over frequency provides the total power scattered, $P_s/P_1 \simeq (8\pi/3)n_er_0^2Ld\Omega \sim 10^{-12}$ for typical parameters ($n_e = 10^{19}$ cm⁻³, $L = 50 \ \mu$ m, $d\Omega = 10^{-4}$). To overcome this small cross section, lasers are used to deliver sufficient power to the Thomson-scattering volume, but the laser power must be balanced against laser–plasma instabilities that can prevent the laser beam from reaching the Thomson-scattering volume. One of the most-limiting instabilities is ponderomotively driven self-focusing. For a laser beam with a Gaussian spatial profile, the self-focusing power threshold is $P_c = 3 \times 10^7 [T_e (\text{keV})/n_e/n_c]$, where $n_c = m_e \omega_0^2/4\pi e^2$ is the critical density for the probe laser.

By limiting the power of the laser to the critical power for self-focusing, the maximum power scattered is given by

$$P_{\rm s}^{\rm max}(W) \simeq 4 \times 10^{-29} \omega_0^2 T_{\rm e}({\rm eV}) L({\rm cm}) \mathrm{d}\Omega. \tag{7}$$

To demonstrate how restrictive this condition is on the parameter space accessible by Thomson scattering, the signal-to-noise can be calculated by assuming Poison statistics, $S/N \approx \sqrt{P_s^{\max} \Delta t/\hbar \omega_0}$, where \hbar is Planck's constant.⁹ For typical conditions $(T_e = 100 \text{ eV}, L = 10^{-2} \text{ cm}, d\Omega = 10^{-4}, \omega_0 = 3.8 \times 10^{15} \text{ l/s})$, spread evenly over 100 resolution units in an ideal system suggests $S/N \sim 10$. From here, it is evident that Thomson scattering requires high electron temperatures, long integration times (Δt) , large Thomson-scattering volumes along the axis of the probe beam (L), or large solid angle collection optics $(d\Omega)$ to increase the signal-to-noise, but each of these parameters has significant constraints within the experimental design.

Intuitively one would expect higher laser powers or higher densities to improve the signal-to-noise, but once the laser power has reached the critical power for self-focusing, the beam will not propagate well to the Thomson-scattering volume. Increasing the density does not help because the increased signal that results from the higher density is directly compensated by the need to reduce the laser power to remain below the critical power for self-focusing. One way to overcome self-focusing, typically at the cost of increasing the Thomson-scattering volume, is to use a DPP.¹⁰ A DPP introduces spatial phase modulation across the laser beam prior to the focusing lens. This phase increases the diameter of the laser spot by distributing the laser power into many speckles, which increases the self-focusing threshold by a factor of ~100 (Ref. 9).

Thomson Scattering from a Maxwellian Plasma

Figure 2 shows the high-frequency and low-frequency parts of the Thomson-scattering spectrum calculated using Eq. (1) assuming Maxwellian ion and electron distribution functions. To measure these spectra, a typical Thomson-scattering instrument uses two spectrometers to independently resolve the high-frequency and low-frequency regimes.^{11,12} The high-frequency spectrum requires lower dispersion to spread the $\Delta\lambda/\lambda_0 \sim 0.1$ spectrum over a detector with approximately 200 resolution units. This can be achieved with a 1/3-m spectrometer with a 150-grooves/mm grating. Resolving the low-frequency spectrum requires a high-dispersion system that can resolve the separation between the ion-acoustic peaks $\Delta\Lambda/\lambda_0 \sim 10^{-3}$ over at least 20 resolution units. This can be achieved with a 1-m spectrometer with a 2400-grooves/mm grating. Often the spectrometers are coupled to optical streak cameras to measure the evolution of the plasma conditions. In these systems, the temporal resolution is determined by the pulse-front tilt introduced by the spectrometers, which is typically of the order of 100 ps (Ref. 13). By trading unrealized spectral resolution for improved time resolution, the temporal resolution can be optimized to the Heisenberg limit.^{14,15}

1. High-Frequency Fluctuations-Electron Plasma Waves

Figure 3 shows the sensitivity of the high-frequency spectrum to the plasma conditions in three different scattering regimes. In the weakly damped regime, the scattering features are very narrow and the sensitivity of the frequency of their peaks provides an accurate measure of the electron density. In this regime, the width of these features is typically dominated by instrument



Figure 3

The sensitivity of the spectrum shown in Fig. 2 to (a) the electron density in the weakly damped regime and electron temperature in (b) the strongly damped regime and (c) the mildly damped regime. The parameters were varied around the central value (black) by +10% (blue) and -10% (red).

broadening and density gradients within the Thomson-scattering volume.¹⁶ By reducing the scattering parameter such that the waves are mildly damped, their width can be increased significantly beyond typical broadening due to gradients and the shape becomes an accurate measurement of the electron temperature, while the peak location remains a measure of the electron density. Further increasing the damping results in a noncollective spectrum where the shape of the scattering spectrum represents the electron distribution function.

2. Low-Frequency Fluctuations—Ion-Acoustic Waves

Figure 4(a) shows the sensitivity of the low-frequency spectrum in the collective regime to the product ZT_e . In this weakly damped regime, the scattering features are very narrow and the sensitivity of their peak location in frequency provides an accurate measure of ZT_e , provided $ZT_e \gg 3T_i$. When this condition is not met, it is convenient to work in the mildly collective regime, where the shape of the ion-acoustic peaks can be resolved, providing a measure of the ion temperature [Fig. 4(b)]. Another technique that is often used to measure the ion temperature in low-Z plasmas is to introduce a small fraction of higher-Z atoms.^{5,6} When the ratio of atomic number to the average ionization (A/Z) is sufficiently different between the two species, additional low-frequency modes are resolvable in the scattering spectrum [Fig. 4(c)].¹⁷ From the relative amplitudes of these two modes, an accurate measure of the ion temperature can be obtained.^{8,18}



Figure 4

Low-frequency spectrum (a) for a single species nitrogen plasma, where $ZT_e = 630 \text{ eV}$ (red), $ZT_e = 700 \text{ eV}$ (black), and $ZT_e = 770 \text{ eV}$ (blue), where $T_i = 20 \text{ eV}$. (b) In the mildly damped regime, the width of the ion feature can be used to measure the ion temperature: $T_i = 18 \text{ eV}$ (red), $T_i = 20 \text{ eV}$ (black), $T_i = 22 \text{ eV}$ (blue), and $ZT_e = 700 \text{ eV}$. (c) Introducing 5% nitrogen (Z = 7) to a hydrogen (Z = 1) plasma provides two low-frequency modes, and their relative amplitudes provide an accurate measure of the ion temperature: $T_e/T_i = 5 \text{ (red)}, T_e/T_i = 3.3 \text{ (black)}, T_e/T_i = 2.5 \text{ (blue)}; and <math>T_e = 100 \text{ eV}$ was held constant. For all calculations, $\alpha = 2$.

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Melting of Magnesium Oxide to 2 TPa Using Double-Shock Compression

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Magnesium oxide (MgO, periclase) is an end-member of the (Mg, Fe)O magnesiowustite mineral, a major constituent of the Earth's lower mantle.^{1,2} It is likely present in the deep interiors of gas giants such as Jupiter and Saturn and in rocky extra-solar planets known as super-Earths.^{3,4} As an abundant component in planets, the physical properties of MgO can influence planetary structure and evolution. The B2 phase (CsCl type) of MgO is expected to be abundant in the mantles of super-Earths and in the rocky cores of gas giants due to the dissociation of MgSiO₃-perovskite.⁴ The melting of MgO could therefore be an important driver of thermal and chemical exchange in the mantles and the core–mantle boundary regions of these planets.^{5,6} Quantifying the melting behavior of MgO to the high pressures and temperatures of planetary interiors is therefore relevant to investigating a number of topical issues in planetary science.

The melt curve of MgO has been studied up to 40 GPa using laser- and resistance-heated diamond-anvil cells,^{7–10} and up to 550 GPa on the principal Hugoniot with decaying-shock experiments.^{11,12} Single shock waves can be used to study melting of a material to the pressure at which the principal Hugoniot crosses the melt curve; however, different experimental techniques are necessary to probe melting at higher pressures. In this work, we apply the double-shock self-impedance-matching technique^{13,14} to measure the melt curve of MgO to 2 TPa—the highest pressure to which any material's melt curve has been studied experimentally.

These experiments were performed on the OMEGA EP Laser System.¹⁵ The targets consisted of a 20- μ m-thick CH polystyrene ablator, a 50- μ m-thick quartz pusher, and a 100- or 200- μ m-thick single-crystal (100) MgO sample. All pieces were laterally 3-mm squares. The target components were held together with 1 to 3 μ m of low-viscosity epoxy. The quartz pusher produced steady shocks in the MgO sample and served as a temperature/reflectivity reference.^{16,17} Two successive shock waves were launched into the sample with a dual laser pulse through ablation of the CH. The first shock was produced with 400 J in a single laser beam with a 6- or 4-ns flattop pulse (0.067 TW or 0.1 TW); the second shock was produced with a net 1500 to 6400 J in one to three beams with a 2-ns flattop pulse (0.75 to 3.2 TW). Distributed phase plates were used to create a spatially uniform irradiance profile with a 95% encircled energy spot diameter of 1100 μ m. The time-resolved diagnostics included a streaked optical pyrometer (SOP)¹⁸ and a dual-channel line-imaging VISAR (velocity interferometer system for any reflector).¹⁹

The measured first (black open circles) and second (red open and solid circles) shock pressure and temperature results are plotted in Fig. 1. At a phase boundary, a material's Hugoniot is often marked by a plateau or reversal in temperature with increasing pressure as thermal energy contributes to a phase transition.^{20,21} This behavior has been observed in shock experiments on diamond,²² SiO₂ (Ref. 17), and the principal Hugoniot of MgO.^{11,12} Results of the second shock show a temperature increase of only 3000 K from 1.2 to 2 TPa; above this pressure, temperature rises rapidly. The three central second-shock data points (solid



Figure 1

The phase diagram of MgO. Black open circles represent the firstshock B1 states in the present work. Red open and solid circles are the second-shock states; both pressure and temperature are measured. The three central second-shock states (solid red circles) are interpreted to be on the melting curve of MgO due to a lack of heating across a large increase in pressure. Melting data from previous experiments are plotted with small solid circles (pink,⁷ purple,⁹ brown,¹⁰ green,¹¹ and blue¹²) and B1-B2 transition data are plotted with ×'s (green,¹¹ and orange²³). Dotted–dashed curves are previously predicted phase boundaries (green,¹³ red,²⁴ orange,²⁵ and purple²⁶). The solid blue curve is a prediction for the principal Hugoniot, ^{26–28} and the solid red curve (interpolated with dashed red) is a prediction for the second shock Hugoniot.²⁶⁻²⁹ The core-mantle boundary conditions are plotted for Saturn³⁰ and 1-, 7.5-, and 15-Earth-mass $(M_{\rm F})$ super-Earths.³¹ The solid black curve is the Simon–Glatzel fit [Eq. (1)] to the melting data in this work and lower-pressure anvil cell melt data,^{8,10} with gray shading representing the uncertainty in the fit parameters.

red circles) are interpreted to lie on the melt curve of MgO because they demonstrate a lack of heating across a large increase in shock pressure, which is attributed to the latent heat of MgO melting. These experiments did not determine the structure of solid MgO, and no structural data exist at these pressures. It is assumed that the MgO melts from B2 in these experiments because no other solid phases are predicted above the B1–B2 transition.

To capture the shape of the high-pressure melt curve, we performed a fit to our data by combining select lower-pressure anvil cell melting data^{8,10} with a Simon–Glatzel equation of the form

$$T_{\rm m}[K] = 3098 \left[\frac{P_{\rm m}({\rm GPa})}{a} + 1 \right]^{1/b},\tag{1}$$

where $T_{\rm m}$ and $P_{\rm m}$ are the temperature and pressure of the melt curve and 3098 K is the melting temperature of MgO at atmospheric pressure.⁸ This empirical relation has been used to describe the melting behavior of other oxides including SiO₂ (Ref. 15) and MgSiO₃ (Ref. 32). The best-fit parameters are given by a = 9.15 (2.23) GPa and b = 3.14 (0.19) with a covariance of -0.39, determined from a nonlinear least squares analysis. A previously published melting curve of MgO (Ref. 15) based on extrapolation of anvil-cell and decaying-shock melting data^{8,9,11} overestimates the melting temperature at 1950 GPa by 27%. This simple fit was chosen based on the discrepancy in the melting temperature of MgO on the principal Hugoniot.

The melt curve in Eq. (1) is plotted in Fig. 1 (solid black) and shows strong agreement with recent density functional theory²⁶ (dashed–dotted purple curve) up to 650 GPa before the curves diverge. Reference 20 overestimates the measured melting temperature at 1950 GPa by 17%. The highest-pressure second-shock equation-of-state point in this work is in the liquid regime of the 173-GPa secondary Hugoniot of MgO and shows general agreement with first-principles equation-of-state simulations of secondary Hugoniots from similar initial shock conditions;^{27–29} the slope of the secondary Hugoniot defined by the two highest-pressure second-shock points in this work does appear steeper than theoretical predictions. The discrepancy between experiment and theory on the melt curve could originate from the complex elastic and plastic responses of MgO during the shock/re-shock and phase transformation processes, which have not been considered in the first-principles calculations. This calls for larger-scale nonequilibrium simulations and crystallographic diagnostics to better understand problems as such. The low-pressure second-shock data in this work demonstrate that the double-shock technique is a valuable method for probing the behavior of MgO in the solid phase at the temperatures and pressures directly relevant to the core–mantle boundary of gas giants similar in size and composition to Saturn³⁰ and super-Earths in the 7.5- to 15-Earth-mass range.³¹

In summary, laser-driven double-shock compression is a valuable method for probing the behavior of MgO in the solid phase at extreme conditions. The present work uses this technique to extend the melting curve of MgO up to 2 TPa and 20,000 K, the highest pressures and temperatures to which any material's melt curve has been probed experimentally. These measurements allowed us to explore the state of the deep interiors of Saturn-sized gas giants and super-Earths. This technique can be used to further quantify the melting behavior of other planetary materials to further investigate the diversity of planetary structures. Additionally, the technique presented in this work will lead to new advances in probing phase transitions of transparent materials up to TPa pressures and significantly advance warm dense matter physics.

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Metastability of Liquid Water Freezing into Ice VII Under Dynamic Compression

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Dynamic compression is often used to create the nonequilibrium conditions needed to study metastability and kinetic effects in materials as they undergo phase transitions.^{1,2} In particular, the pressure-induced phase transformation of liquid water solidifying into ice VII has been the focus of many experimental and theoretical works.^{3–9} Under rapid submillisecond compression, the liquid phase can persist metastably well into pressure–temperature conditions where ice VII is the stable phase.^{3–5} Previous experimental studies found that liquid water can remain metastable to at least 7 GPa— ~5 GPa higher than expected based on the equilibrium phase diagram—before homogeneously freezing into ice VII when quasi-isentropically (ramp) compressed over hundreds of nanoseconds.^{3–5} This work ramp compresses liquid water over the highest compression rates to date (up to ~3 GPa/ns) to further investigate its metastability limit.

Water was ramp compressed into the ice VII phase in experiments at the Omega Laser Facility.^{10,11} The liquid–ice VII phase transition in a thin water layer, sandwiched between a baseplate and a sapphire or quartz window, was diagnosed using a velocity interferometer system for any reflector (VISAR).¹² Since ice VII is ~5% more dense than liquid water at the phase transition conditions, the volume of the thin water layer abruptly decreases during the phase transition (~1-ns duration), which alleviates pressure on the water/window interface despite the continuously increasing pressure drive.^{3–5} The VISAR records a corresponding dip in the water/window interface velocity, which we interpret as the liquid freezing into ice VII (Fig. 1).



Figure 1

Interface velocities and corresponding pressures (applicable to all curves) from experiment shot 29419 and the post-shot simulations. An inset of the target components relevant to the experimental measurements and simulations is shown, where "B" is the sapphire baseplate, "W" is the water, "Witness" is the sapphire witness, and "Window" is the sapphire window. The VISAR probes a reflective Al coating at the baseplate/witness and water/window interfaces to measure their velocities. A dip in the water/window interface pressure, resulting from the liquid water freezing into the ~5%-more-dense ice VII phase, is observed near 24 ns and 7.5 GPa in the experiment and simulation using the liquid/ice VII equation of state and classical nucleation theory–based kinetics model.

Water was compressed at rates spanning from 0.2 to 3 GPa/ns over 15 experiments, where the loading rate was varied by changing the laser intensity, the baseplate thickness, and the window material (e.g., the lower impedance of quartz compared to sapphire leads to shallower ramp compression profiles) (Fig. 2). We find that the liquid–ice VII freezing pressure, defined as the pressure in the liquid at the peak velocity before the dip, for water compressed on the principal isentrope increases with compression rate to at least ~8 GPa [Fig. 2(a)]. We observed freezing at pressures as high as ~9 GPa; however, additional heating of <8 K above the principal isentrope cannot be ruled out, which could further raise the freezing pressure.⁴ These results indicate that liquid water can exist to at least ~3.5× higher pressure than the onset of metastability (2.2 GPa) (Ref. 9) and that the metastability limit is at least ~11% higher than previously reported.^{3–5} Agreement between data at 0.1 to 0.3 GPa/ns in Fig. 2(a) from this work (Omega), Dolan *et al.* (Z),³ and Nissen *et al.* (Thor)⁴ (all room temperature), obtained using different target component materials, suggests that ice VII is nucleated homogeneously in the bulk and not heterogeneously at the various window or baseplate surfaces.



Figure 2

(a) Liquid–ice VII freezing pressure versus compression rate [(defined from 2.2 GPa ([onset of metastability)] to the freezing pressure]) and (b) pressure histories of the water/window interface for all shots ordered by decreasing compression rate and shifted in time for clarity. In the legend of (a), S and Q denote sapphire and quartz windows, respectively, and ΔT_0 is the initial temperature increase above the principal isentrope. Asterisks in (b) mark the pressure relaxation interpreted as freezing.

Our experimental results can be reproduced in hydrodynamic simulations (*ARES*) using a kinetics model (SAMSA)¹³ that, remarkably at these extreme conditions, is fundamentally based on classical nucleation theory (CNT).⁸ The baseplate/water/ window portions of the target were simulated using a pressure input on the front baseplate surface that were determined from the shot-specific sapphire "witness" measurements adjacent to the water layer. The same pressure relaxation at the water/window interface observed in the experiment is also observed in the simulation using the CNT-based kinetics model⁸ and a multiphase equation of state (EOS) for the liquid and ice VII phases¹⁴ (Fig. 1). This pressure relaxation is concurrent with the onset and completion of freezing in the simulations. The "null case" of no phase transition, represented by using only the liquid EOS, does not show the dip in the water/window interface pressure, suggesting that the dip observed in the experiment is indeed the result of freezing and not wave reverberations within the target.

The experiments reported here are at the frontier of using experimental ultrafast science to explore metastability and kinetics associated with phase transitions. It is remarkable that recent theoretical and numerical advances provide a detailed understanding of the observed phenomena, while relying on the fundamentally simple picture of homogeneous nucleation using CNT. This could have implications for our general understanding of phase transformations at extreme conditions.

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Measurement of Chromatic Aberrations Using Phase Retrieval

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Phase retrieval is a computational method for estimating the phase of an electromagnetic field based on measurements of the intensity in one or more planes. For wavefront-sensing applications, the phase of interest is in the pupil plane of an optical system, and typically the aperture and image-plane intensity from a point source [the point-spread function (PSF)] are known.¹ The iterative process of retrieving the phase involves forming an initial estimate of the wavefront in the pupil plane, simulating a propagation of that field to the image plane (which typically involves a Fourier transform), and comparing the resulting intensity distribution with the measured intensity via an error metric. The wavefront estimate is then modified to improve agreement in the image plane.

Among the various applications of this general approach is the improvement of diagnostic tools for the OMEGA EP laser. In this case, phase retrieval complements measurements provided by a conventional Shack–Hartmann wavefront sensor (SHWFS) with an estimate of non-common-path error and differential piston between regions of the segmented beam.² A proposed way to improve this system is to also estimate chromatic aberrations in the system. These can arise in a chirped-pulse–amplification laser such as OMEGA EP in the form of (1) residual angular dispersion from the pulse stretcher and compressor and (2) axial dispersion (longitudinal chromatic aberration) due to transmissive elements, both of which can lead to a significant reduction in the focused intensity.

We use a modal approach to modeling chromatic effects. Angular dispersion is modeled as a wavefront tilt that varies linearly with wavelength and longitudinal chromatic aberration as a defocus that varies linearly with wavelength. By forcing the spectrally varying components of the wavefronts to fit this model, we can mitigate uniqueness problems that would arise if they were allowed to vary independently. In our model, we assume that (1) there is an initial estimate of the monochromatic wavefront W_0 measured by a SHWFS and (2) there is also a non-common-path error between the SHWFS and the focusing optics, which must be estimated in terms of a monochromatic wavefront W_M , in addition to the chromatic aberrations $(\lambda - \lambda_r) W_c$, relative to a reference wavelength λ_r , so that

$$W(\xi,\eta) = W_0(\xi,\eta) + W_M(\xi,\eta) + (\lambda - \lambda_r)W_c(\xi,\eta), \tag{1}$$

where $W_{\mathbf{M}}(\xi,\eta) \equiv \sum_{n=0}^{N} a_n Z_n(\xi,\eta)$ and $W_{\mathbf{c}}(\xi,\eta) \equiv \sum_{n=0}^{N} c_n Z_n(\xi,\eta)$.

Having mentioned above "differential piston between regions of the segmented beam,"² this model does not capture the segment-to-segment errors that can exist from a tiled grating. The a_n and c_n coefficients would usually vary from segment to segment (including piston tip and tilt).

In Eq. (1), ξ and η are pupil-plane coordinates, λ is the wavelength of a single spectral component, and Z_n is the *n*th Zernike polynomial. The phase-retrieval algorithm jointly estimates the unknown monochromatic wavefront along with the chromatic

aberrations in terms of the sets of coefficients a_n and c_n , respectively. The exit pupil amplitude is assumed to be a known aperture with transmittance $P(\xi, \eta)$, and each spectral component is given a scalar amplitude weight, based on the known power spectrum. Each spectral component is propagated from the pupil plane to the image plane separately, and the total PSF in the image plane is the incoherent sum of the spectral components. The error metric is the normalized sum-squared difference between the simulated and measured PSF's.

As a simple test case, we simulated PSF's with angular dispersion of 1.2 waves peak-to-valley (p-v) of tip/tilt and axial dispersion of 1 wave of defocus, both across an 8-nm bandwidth. Besides the global minimum, three local minima also appear, where one or both of the signs of the dispersion parameters are reversed. When a known defocus is added, two of the local minima disappear and the error metric appears as in Fig. 1(a).



Figure 1

The twin image problem is illustrated. (a) Error metric as a function of the two dispersion variables, with a known defocus included. The global minimum (correct values) is in the center of the plot at coordinates (1.2, 1). (b) Illustration of how defocus does not eliminate the joint axial–angular twin.

We can explain the appearance of these minima and their relation with defocus by analogy to the well-known twin image problem in monochromatic phase retrieval, which arises because the absolute value of the Fourier transform of any complex signal f(u) is the same as that of its complex conjugate flipped about the origin f(-u). Defocus cannot eliminate the joint axial-angular twin, when the signs of both the angular and axial dispersion are reversed simultaneously. This effect can be visualized with the ray trace in Fig. 1(b), which shows axial dispersion causing each spectral component to focus at a different horizontal position and angular dispersion causing the focal points to spread vertically.

However, the PSF produced by the twin is subtly different due to the scaling effect of the wave number $2\pi/\lambda$ in the complex exponential and any asymmetry in the spectrum. If this results in a significantly higher error metric for the twin than the true solution, a sufficient strategy for avoiding this minimum is to check the error metric against a threshold value after the optimizer converges; if it is too high, reverse the signs of the dispersion coefficients and perform another round of local nonlinear optimization to confirm whether that minimizes the error metric.

To test this strategy, the algorithm was run through a series of trials with a variety of simulated true reference wavefronts and various starting guesses for the dispersion parameters. In all cases, the true chromatic aberrations were 1.2 waves p–v of angular dispersion and 1 wave p–v of axial dispersion across an 8-nm bandwidth (or equivalently stated as 0.15 waves/nm and 0.125 waves/nm, respectively). The spectral weights for intensity were representative of measured spectra in fully amplified shots on OMEGA EP [shown in Fig. 2(a)]. Each of the randomized true reference wavefronts was the sum of a known part W_0 with a magnitude 0.4 waves rms and an unknown part W_M [simulating non-common-path error, an example of which is shown in Fig. 2(b)] that had to be retrieved and had a magnitude of 0.11 waves rms. Known defocus in the amount of 1.5 waves was also added.



Figure 2

Monte Carlo simulation results: (a) Intensity spectrum used in Monte Carlo simulations; (b) representative example of random monochromatic wavefront aberration (defocus removed); and (c) resulting phase-retrieval error metric versus error of the reconstruction, where the error is computed over the full bandwidth. Successful cases in (c) have a low estimation error, and failed cases are identifiable by a large error metric in the PSF fit. Cases above the vertical line are considered failed cases.

Both the unknown part of the monochromatic reference wavefront and the dispersion terms were retrieved. A bootstrapping process was used in which the monochromatic reference wavefront was optimized alone while the dispersion terms were left at zero. Once that converged, a second stage was carried out, allowing both the monochromatic and dispersion terms to vary. In combination with the twin phase check described earlier, this strategy was successful in all but two of the 1600 tested cases. In the failed cases, the normalized root-mean-squared error between the estimated and true PSF's was about 10× greater than the error metric for successful cases, as shown by the outliers in Fig. 2(c). Therefore, if one of these failures occurs in practice, it can be recognized by the large residual error metric and recovered from by performing another optimization run from a different random starting point.

To test the proposed chromatic aberration retrieval algorithm in a real-world scenario, we applied the method to a small-scale laboratory setup. This test bed was previously built to generate and measure chromatic aberrations with a method that utilizes a 2-D grating to simultaneously disperse spectral components and provide focus diversity.³ For consistency with Ref. 3, we represent angular and axial dispersion in terms of pulse-front delay (PFD), and radial group delay (RGD), respectively.

The conceptual layout of the experimental setup is shown in Fig. 3(a). A Superlum SLD-52 superluminescent light-emitting diode (SLED) served as a broadband light source, having a bandwidth of approximately 100 nm and spectral distribution of intensity shown in Fig. 3(b). Given the large bandwidth of the source, significant chromatic dispersion can be introduced with a pair of lenses (L1 and L2), which together provided $\gamma = 10.0$ fs of RGD. A fused-silica wedge (W) with a 1° apex angle imparted $\beta = 2.5$ fs of PFD. Images of the PSF were captured by a camera, which was translated along the optical axis to provide focus diversity.

The focal sweep was performed first with a Semrock narrowband filter (F) with a FWHM bandwidth of less than 7.2 nm and a central wavelength of 1.03 μ m to provide images that were effectively monochromatic; in the second sweep the narrowband filter was replaced by a neutral-density filter. From the spectrally filtered images, we measured the monochromatic contribution to the wavefront without the risk of confounding effects from dispersion. The nonfiltered set of PSF images exhibited a very subtle blurring effect due to dispersion and the increased bandwidth. Because these effects are subtle, the dominant monochromatic wavefront was first estimated in terms of Zernike coefficients; then the dispersion was retrieved in terms of tip, tilt, and focus coefficients that vary linearly with wavelength. Lastly, all of the variables were jointly optimized to produce the final estimate.

The final pupil wavefront estimate, shown in Fig. 3(c), differed by only 16.9-nm rms from the monochromatic result. The final dispersion estimate consisted of 4.68 fs of PFD and 9.13 fs of RGD, which differ by 6.4% and 8.7% from the expected values of



Figure 3

Laboratory experiment setup and results: (a) Experimental system layout (reprinted from Ref. 3); (b) intensity spectrum of the SLED; and (c) final wavefront estimate (full bandwidth). A: apodizer.

5.0 fs and 10.0 fs, respectively. Compared to the results in Ref. 3, where the PFD estimate was within 0.5 fs of expectations and the RGD within 0.1 fs, these results for PFD are similar, while the RGD is less accurate.

In summary, we have developed a simulation model and optimization process for the joint estimation of linear chromatic aberrations in addition to monochromatic aberrations using a measured broadband PSF together with a known aperture, spectrum, and initial wavefront estimate. We found a bootstrapping strategy that first estimated the monochromatic wavefront correction followed by optimization of the chromatic parameters to be highly successful. A test of this approach in a laboratory experiment produced encouraging results.

This material is based upon work supported by the Department of Energy National Nuclear Security Administration under Award Number DE-NA0003856, the University of Rochester, and the New York State Energy Research and Development Authority.

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A Highly Efficient, 10-J Output Signal Amplifier for Ultra-Intense All-OPCPA Systems

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The highest-energy beams and the shortest pulses are always in the mainstream of laser science and engineering to create peta- and exawatt lasers.¹ Optical parametric chirped-pulse–amplification (OPCPA) systems, pumped by high-energy Nd:glass lasers, have the potential to produce ultra-intense pulses (> 10^{23} W/cm²). Existing large-scale glass lasers could be used to pump a large all-OPCPA system.² While front-end OPCPA stages are mainly focused on spectral, temporal, and phase characteristics of beams, the final OPCPA stages additionally need to be energy efficient. We report on the performance of the final high-efficiency amplifier in an OPCPA system based on large-aperture (63×63 -mm²), partially deuterated potassium dihydrogen phosphate (DKDP) crystals.

The experiment was performed on the Multi-Terawatt (MTW) Laser System, which is a hybrid OPCPA and Nd:glass laser. For all-OPCPA,² the MTW laser was switched to a narrowband mode at 1053 nm with energy up to 60 J and pulse-length variation from 1.2 to 1.6 ns. This radiation was converted to 526 nm using second-harmonic generation with an efficiency of 70%. The "green" beam at 526 nm travels by way of two large-scale periscopes into the next room through a series of vacuum spatial filters to maintain high beam quality. Finally, this 45-mm × 45-mm beam pumps a 48-mm or 52-mm-long DKDP crystal on the final stage synchronously with the seed beam. The chirped seed beam ($\tau = 1.5$ ns, E = 240 mJ, $\Delta\lambda$ from 830 to 1010 nm) was created using a sequence of four lower-energy amplifiers seeded by a white-light continuum. The 11.8-J output signal was compressed to 19 fs.

The maximum pump-to-signal conversion efficiency of 37% was achieved with a 52-mm-long DKDP crystal (deuteration level of 70%) and 40 J of pump energy at 527 nm due to the flattop super-Gaussian pump beam profile and flat-in-time pulse. The seed and the pump were precisely synchronized in time. The shape of the 1.6-ns pump pulse was precompensated in the front end to reach a top-flat shape on the final amplifier. The input pump pulse and the residual pump pulse are shown in Fig. 1(a). The resulting hole in the residual pulse corresponds to the place where the seed pulse is located. It also demonstrates how deeply the pump pulse is depleted.

A deep saturation regime in the 52-mm-long crystal smooths output intensity modulation of the seed beam. Figure 1(b) shows a much better output signal spectrum at higher pump energies, even with a moderate quality of the initial seed spectrum. Figure 2 shows saturation of amplification with much better quality of the signal beam (d) than the seed beam (a) and deep depletion of the residual beam at 2ω (b) compared to the pump beam (c).

The maximum conversion efficiency from the pump beam into a signal of 37% (Fig. 3) was achieved with the 52-mm-long DKDP crystal (70% deuteration level) and pulse duration of 1.2 ns (FWHM).



Figure 1

(a) An oscillogram with the input pump pulse and output residual pump pulse, both at 526 nm; (b) spectrum of the input seed beam (black curves) and the output signal beam (red curves) at the maximum pump beam energy.



Figure 2

Images of typical interaction beams in the 52-mm-long DKDP: (a) the input seed beam, (b) the residual pump beam at 2ω , (c) the input pump beam, and (d) the output signal beam. NOPA: noncollinear optical parametric amplifier.



Figure 3 Maximum pump-to-signal conversion efficiency and depletion of the pump energy for the 1.2-ns pump pulse.

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Broadband Sum–Frequency Generation of Spectrally Incoherent Pulses

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High-energy nanosecond solid-state laser systems operating in the near infrared require frequency conversion to improve the efficiency of laser–matter interaction. This is generally done with a sequence of two nonlinear crystals, one for frequency doubling from 1ω (1053 nm) to 2ω (526.5 nm) and one for mixing of the resulting 2ω with the remaining 1ω to generate 3ω pulses (351 nm) (Ref. 1). The spectral acceptance of the tripling stage can be increased using two crystals or angular dispersion at 1ω (Refs. 2 and 3), but neither scheme allows for efficient operation beyond ~1 THz with incoherent nanosecond pulses. Simulations show that spectrally incoherent broadband pulses can mitigate the detrimental laser–plasma instabilities and on-target beam imprint, therefore increasing the coupling efficiency of energy into the target.⁴

A novel sum-frequency generation (SFG) scheme based on a noncollinear interaction between a 1 ω broadband angularly dispersed pulse and a narrowband 2 ω pulse allows for efficient frequency conversion into broadband 3 ω pulses. Experimental results are in excellent agreement with simulations, demonstrating the generation of spectrally incoherent 3 ω pulses with bandwidths larger than 10 THz in a relatively thick 1-cm KDP crystal. This scheme can be implemented with commercially available large-aperture diffraction gratings and nonlinear crystals to support a new generation of high-energy laser facilities delivering spectrally incoherent pulses.

The wave-vector mismatch for SFG of a broadband angularly dispersed pulse (frequency $\omega + \Omega$, angular dispersion *D*) with a narrowband pulse (frequency 2ω) in a noncollinear geometry [Fig. 1(a)] along the wave vector at ω is

$$\Delta k(\Omega, \alpha, D, \theta) = k_{o}(\omega + \Omega)\cos(D\Omega) + k_{o}(2\omega)\cos(\alpha) - k_{e}(3\omega + \Omega, \theta)\cos[\beta(\Omega, \alpha, D, \theta)],$$
(1)

where α is the internal noncollinear angle between the 1 ω and 2 ω beams, θ is the frequency-dependent angle between the crystal axis and the wave vector at $3\omega + \Omega$, and β is the frequency-dependent angle between 1 ω and $3\omega + \Omega$ beams. There is a continuum of combinations of the three degrees of freedom (D, α , and θ) that cancels the phase mismatch and its frequency derivative at $\Omega = 0$, therefore yielding broadband SFG. For example, operation with $\alpha = 1.7^{\circ}$ and $\Delta = -0.59$ mrad/nm ($\Delta = -2\pi c D/\lambda^2$) in a Type-I KDP crystal allows for the conversion of ~10 THz of bandwidth from 1053 nm to 351 nm in a 1-cm crystal, i.e., $10 \times$ larger than in a collinear scheme [Fig. 1(b)]. Crystal-angle detuning allows for SFG of frequency components symmetrically located relative to 1ω , e.g., the signal and idler resulting from parametric amplification of a 1ω signal close to spectral degeneracy with a pump at 2ω (Ref. 5).

The SFG demonstration follows the principle described in Fig. 1(a). A collinear optical parametric amplifier (OPA) seeded with either a monochromatic tunable signal or a spectrally incoherent signal originating from an amplified spontaneous emission source at wavelengths below 1053 nm is pumped by a 1.5-ns pulse at 526.5 nm, leading to a combined signal and idler symmetric relative to 1053 nm (Ref. 5). The OPA 1 ω output is spectrally dispersed by an 802.5-1/mm transmission grating at Littrow, which is re-imaged onto a 1-cm KDP crystal, itself re-imaged onto a 2305-1/mm transmission grating that compensates for the 3 ω angular dispersion resulting from the 1 ω angular dispersion and noncollinear SFG geometry. The OPA 2 ω pump is separately re-imaged to the SFG crystal. For the fixed 1 ω angular dispersion *D*, the noncollinear angle α is optimized for frequency conversion of a



Figure 1

(a) Sum-frequency generation of a broadband angularly dispersed 1 ω pulse with a narrowband 2 ω pulse in a noncollinear geometry; (b) relative SFG efficiency for $\alpha = 1.7^{\circ}$ and $\Delta = -0.59$ mrad/nm versus crystal angle detuning.

monochromatic signal at 1030 nm and the corresponding idler at 1077 nm at the same crystal angle, therefore ensuring symmetric phase matching relative to 1053 nm.

The spectral acceptance measured with a monochromatic tunable 1ω OPA seed shows that noncritical SFG is obtained at one specific crystal angle, while detuning matches the SFG to a pair of signal and idler beams at opposite frequencies relative to 1ω [Figs. 2(a) and 2(b)], in excellent agreement with the simulations [Fig. 1(b)]. Broadband spectrally incoherent light at 1ω is obtained by seeding the OPA with an amplified spontaneous emission pulse covering ~10 nm at 1030 nm. With the combined signal and idler, SFG-crystal tuning allows for the generation of more than 10 THz of bandwidth either centered at 351 nm or in two symmetric side lobes, depending on the crystal angle (Fig. 3).



Figure 2

Spectral acceptance characterization of a 1-cm KDP crystal with a tunable monochromatic 1ω signal: relative SFG energy versus (a) frequency relative to 3ω and crystal angle and (b) lineouts at four crystal angles.



Figure 3

Spectral density of a generated broadband spectrally incoherent pulse at 3ω (a) as a function of crystal angle and (b) for four crystal angles. In (b), the spectral density of the 1ω input to the SFG stage is plotted with a black line.

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Generation of Strong Magnetic Fields for Magnetized Plasma Experiments on a 1-MA Pulsed-Power Machine

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Introduction

The development of techniques for the generation of strong magnetic fields provides an opportunity to investigate plasmas in megagauss (MG) fields. Strong magnetic fields change basic properties of hot and dense plasma. Studies of magnetized plasmas are relevant to basic and applied plasma physics, controlled fusion research, and astrophysics. Astrophysical magnetic fields can reach a value of 10⁹ MG in magnetars¹ and a value of 1 to 100 MG in white dwarf plasma.² High magnetic fields also provide an increased neutron yield in inertial confinement fusion.³ A 30- to 40-MG magnetic field plays a key role in the magnetized liner inertial fusion (MagLIF) approach to fusion.⁴ Magnetic fields change the dynamics of plasma expansion, ^{5–7} the development of plasma instabilities, and parametric effects. Laser–plasma interactions in external magnetic fields display unusual plasma expansion such as the generation of disk-like plasma in a 2- to 3-MG transverse magnetic field.⁷ Narrow plasma jets are generated in the longitudinal magnetic field.^{5,6} Astrophysical magnetized plasmas can be scaled to laboratory plasmas.⁵ Megaampere-class pulsed-power machines routinely generate MG magnetic fields.

Plasma in an Azimuthal Magnetic Field

Unusual dynamics of plasma expansion in the azimuthal 1- to 3-MG magnetic field of a rod load were observed in Ref. 7. Here, the results of an additional series of shots are presented. Figure 1(a) shows a scheme of the laser–plasma interaction (LPI) experiment with a laser pulse focused on the surface of the Al rod load 0.9 mm in diameter. A current in the load generates an azimuthal magnetic field. The laser intensity in the focal spot is 3 to 5×10^{15} W/cm². Without the laser pulse, a strong magnetic field contains plasma that arises on the surface of the load. After the laser shot, side-on laser-imaging diagnostics show the formation of two plasma jets on the front and rear sides of the load. The shadowgrams in Figs. 1(b) and 1(c) present jets of laser-produced plasma at 6 ns after the laser pulse. The magnetic field on the surface of the rod load. Diagnostics with a tilted probe [Fig. 1(c)] explain the focal spot; the second smaller jet was seen on the rear side of the rod load. Diagnostics with a tilted probe [Fig. 1(c)] explain the front and rear jets as parts of a plasma disk around the load. Plasma propagated along the magnetic field and formed a thin disk around the load with ring structures in it.

The radial size of the disk is longer and the plasma density is higher in the front half-disk, so the disk is not symmetric. Plasma expansion is observed during >10 ns after the laser pulse. The electron plasma density in the rings is 6 to 8×10^{18} cm⁻³. The formation of the disk happens only in the presence of the strong azimuthal B field. Two-frame shadowgrams and interferograms show that the disk expands radially with a velocity of ~250 km/s. The electron temperature of the plasma is measured from the x-ray Al K-shell spectra to be about 400 eV in the area of interaction. Plasma disks were observed in LPI with Al, Cu, and Ti rod loads.

Two-dimensional cylindrical magnetohydrodynamic (MHD) simulations of the plasma in the strong transverse magnetic field were performed in Ref. 8. The simulations with the current in the rod that resulted in a magnetic field of 3 MG on the rod



showed the formation of the density wave localized in the axial direction and moving in the radial direction. The propagation of the density wave continues after the end of the laser pulse. The azimuthal magnetic field strongly changes in time in the region of the density wave due to magnetic-field generation by crossing density and temperature gradients. The change in the azimuthal magnetic field due to the density wave is comparable to the magnitude of the azimuthal field generated by the current in the rod. The magnetic β parameter at 1 to 2 ns after the laser pulse is about 1 and decreases after the pulse. The thermal pressure is responsible for the motion of the plasma. When plasma moves in the radial direction away from the rod, the thermal pressure decreases and the magnetic pressure has the main role in the plasma expansion. The density of the plasma in the MHD simulations is shown in Fig. 2 and is in agreement with the side-on shadowgrams in Fig. 1(b).



Plasma in the Longitudinal Magnetic Field

Plasma takes the shape of a jet in the longitudinal magnetic field of 0.1 to 0.2 MG (Ref. 5). A 1-MA pulsed-power machine allows for the investigation of plasma jets in higher magnetic fields.⁶ Plasma in the B field of the coil was produced by the Leopard laser operated with a 0.4-ps pulse. Intensity of the laser beam in the focal spot was 2 to 6×10^{18} W/cm². The laser target was placed at 1 mm from the copper coil 2.5 to 3 mm in internal diameter. The axial B field at this point was 0.5 to 0.7 MG, depending on the timing relative to the current pulse. A Si laser target was used to avoid the influence of eddy currents. The size of the target was $2 \times 2 \text{ mm}^2$ and 10 μ m thick. The interferogram and shadowgram in Fig. 3 illustrate the collimation of plasma and the formation of the long plasma jet on the front side and the shorter rear jet in the B field of 0.7 MG. A plasma plume begins focusing at ~1 mm from the target and forms a narrow jet. The velocity of propagation of the jet tip calculated is ~200 km/s. Laboratory plasma jets of this type can be scaled to the astrophysical jets.⁵ In MHD simulations, plasma expands, forming a jet, and the magnetic field in the plasma is much weaker than the external magnetic field. At the same time, the magnetic field at plasma edges increases and becomes larger than the external field that results in collimation due to magnetic pressure. While the magnetic field is compressed at the edges of the jet, it is much smaller inside the jet forming the magnetic-field envelope. This



Figure 3

(a) Interferogram and (b) shadowgram of the Si target during the current and the magnetic field in the coil load at 7 ns after the laser pulse; (c) a schematic of the laser beam (L) and target (T) near the coil load.

envelope maintains the collimated jet during all its evolution. Later the diffusion of the magnetic field into the envelope makes the field inside the envelope close to the field on its edge. A plasma jet in a magnetic field depends on the magnetic β parameter being small. Figure 4 shows the expansion of the jet tip in simulations (the dashed line), and the tip positions in experiment and simulations are in a good agreement.



Figure 4 Positions of the jet tips (squares). MHD simulation (dashed curve).

Two-Plasmon Decay in the MG Magnetic Field

Two-plasmon decay (TPD) plays an important role in LPI. TPD occurs near one quarter of the plasma critical density n_c for the laser frequency ω_0 , and the resulting two Langmuir waves (plasmons) have "blue" and "red" spectral shifts compared to the $\omega_0/2$ frequency. Wave conversion involving these two plasmons generates new light waves with frequencies around $\omega_0/2$. Wave conversion involving TPD plasmons can also generate light with frequency around $3/2\omega_0$, which easily leaves the plasma and makes a robust diagnostic. The strong magnetic field produces a shift proportional to the square of the electron Larmor frequency $\omega_e: -(\omega_{ce})^2/\omega_0$ in addition to the thermal shifts of the "red" and "blue" spectral components.⁹ The narrowband Nd:glass laser used in the TPD experiments generate pulses at 1053 nm with $\Delta\lambda_0 \sim 10$ pm and an energy of 6 J at 2 ns. Al and Ni rod loads 1 mm in diameter were used to generate the magnetic field of 2 to 3 MG in the surface plasma. A laser pulse was focused on the rod surface with an intensity of 1 to 3×10^{14} W/cm². An intensified charge-coupled–device (ICCD) camera was used to record the $3/2\omega_0$ emission.

Figure 5(a) presents $3/2\omega_0$ spectra from the Al rod load. The strong 2- to 3-nm widening and 2- to 4-nm shift of "red" and "blue" $3/2\omega_0$ spectral components were observed. Both red and blue $3/2\omega_0$ components are clearly seen but the blue component is weaker. The $3/2\omega_0$ emission was not seen in Ni and Cu loads. $3/2\omega_0$ emission can be observed only if TPD instability develops, and the TPD threshold is inversely proportional to the density scale length. The strong $3/2\omega_0$ emission in Fig. 5(a) is an indication of extended plasma with a gradual density profile, and the absence of $3/2\omega_0$ emission is an indication of more-localized plasma with a steep density profile.





Conclusion

In this work, it was shown that the 1-MA pulsed-power machine provides a robust platform for experiments with plasma in MG magnetic fields. The dynamics of expansion of the laser-produced plasma in the strong transverse and longitudinal magnetic fields were studied with the rod and coil loads. The expanding plasma takes the shape of a thin plasma disk in the azimuthal field of the rod load. Plasma is confined in the vertical direction by the 2- to 3-MG magnetic fields. In the longitudinal magnetic field, laser-produced plasma generates narrow 3- to 4-mm jets with a density of 10^{19} to 10^{20} cm⁻³. The TPD parametric instability generates wide and shifted "red" and "blue" components of $3/2\omega_0$ emission in the 2- to 2.7-MG field. Finally, pulsed-power technology provides a capability for the investigation of plasmas and laser-matter interaction in 1- to 4-MG magnetic fields at the university-scale machine.

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FY21 Q2 Laser Facility Report

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During the second quarter of FY21, the Omega Facility conducted 295 target shots on OMEGA and 202 target shots on OMEGA EP for a total of 497 target shots (see Tables I and II). OMEGA averaged 10.2 target shots per operating day, averaging 91.5% Availability and 97.8% Experimental Effectiveness. OMEGA EP averaged 7.8 target shots per operating day, averaging 88.4% Availability and 94.1% Experimental Effectiveness.

	<u> </u>	<u>`</u>	
Program	Laboratory	Planned Number of Target Shots	Actual Number of Target Shots
ICF	LLE	60.5	61
	LANL	22	24
	LLNL	11	14
ICF Subtotal		93.5	99
HED	LANL	33	34
	LLNL	27.5	31
HED Subtotal		60.5	65
LBS	LLE	22	23
	LLNL	11	9
LBS Subtotal		33	32
AIBS		27.5	21
NLUF		55	51
Calibration	LLE	0	27
Grand Total		269.5	295

Table I: OMEGA Laser System target shot summary for Q2 FY21.

Program	Laboratory	Planned Number of Target Shots	Actual Number of Target Shots
ICF	LLE	28	40
	LLNL	7	7
ICF Subtotal		35	47
HED	LLE	7	8
	LANL	14	19
	LLNL	14	14
HED Subtotal		35	41
LBS	LLE	21	31
	LLNL	35	42
LBS Subtotal		56	73
LNet		7	6
NLUF		24.5	23
Calibration	LLE	0	12
Grand Total		157.5	202

Table II: OMEGA	EP Laser System	target shot summar	v for O2 FY21
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For OMEGA and OMEGA EP shot planning efforts, the Experimental Proposal and Shot Request Form (SRF) systems were upgraded. The SRF now requires an association to the Proposal template during creation and uses this information to determine the date of the SRF. With this upgrade all SRF's can follow schedule changes with a single update.

The OMEGA Stage-F Alignment Sensor Package upgrade project has now completed 30 of 60 beamline systems. The cameras are being replaced with higher-resolution digital charge-coupled-device (CCD) equipment.

The OMEGA EP beam apodization system that ensures that the gaps between gratings in the pulse compressors do not see damaging laser fluence (known as the "gapodizers") had position sensors relocated for enhanced system safety. Two optics replacements have occurred of interest to the PI community: The OMEGA EP lower compressor deformable mirror was replaced due to laser damage accumulated over years of operation, and the OMEGA EP "backlighter" beam's off-axis parabola (OAP) focusing optic was replaced with a reworked OAP.

Publications and Conference Presentations

Publications

I. A. Begishev, S.-W. Bahk, C. Dorrer, C. Feng, M. J. Guardalben, C. Jeon, R. G. Roides, M. Spilatro, B. Webb, D. Weiner, J. D. Zuegel, and J. Bromage, "A Highly Efficient, 10-J Output Signal Amplifier for Ultra-Intense All-OPCPA Systems," Proc. SPIE **11666**, 1166607 (2021).

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-Forthcoming Publications-

D. H. Edgell, A. M. Hansen, J. Katz, D. Turnbull, and D. H. Froula, "Unabsorbed Light Beamlets for Diagnosing Coronal Density Profiles and Absorption Nonuniformity in Direct-Drive Implosions on OMEGA," to be published in Review of Scientific Instruments.

T. M. Johnson, A. Birkel, H. E. Ramirez, G. D. Sutcliffe, P. J. Adrian, V. Yu. Glebov, H. Sio, M. Gatu Johnson, J. A. Frenje, R. D. Petrasso, and C. K. Li, "Yield Degradation due to Laser Drive Asymmetry in D³He Backlit Proton Radiography Experiments on OMEGA," to be published in Review of Scientific Instruments.

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Conference Presentations

J. L. Shaw, M. V. Ambat, S.-W. Bahk, I. A. Begishev, R. Boni, J. Bromage, G. Bruhaug, S. Bucht, E. M. Campbell, A. Davies, C. Dorrer, P. Franke, R. K. Follett, V. N. Goncharov, D. Haberberger, A. Howard, G. W. Jenkins, J. Katz, T. J. Kessler, B. Kruschwitz, M. M. McKie, L. Nguyen, J. B. Oliver, J. P. Palastro, D. Ramsey, M. A. Romo-Gonzalez, T. T. Simpson, D. Turnbull, L. J. Waxer, H. Wen, D. H. Froula, P. M. King, N. Lemos, G. J. Williams, H. Chen, F. Albert, M. D. Sinclair, C. Joshi, F. A. Hegmann, D. Purschke, S. Stoller, N. Vafaei-Najafabadi, B. Malaca, J. L. Martins, J. Vieira, S. Jolly, F. Quéré, A. Di Piazza, G. Gregori, Z. Li, T. M. Antonsen Jr., A. Arefiev, K. Weichman, and R. Bingham, "New Developments in Laser Wakefield Acceleration at the Laboratory for Laser Energetics," presented at the Oxford Hilary Series Seminar, virtual, 25 January 2021. J. L. Peebles, J. R. Davies, D. H. Barnak, M. J. Bonino, T. Cracium, R. Betti, and P.-Y. Chang, "Probing Strong Electric and Magnetic Fields using Axial Proton Radiography of Laser-Driven Coils," presented at Charged Particle Radiography in High-Energy-Density Laboratory Plasmas, virtual, 25–28 January 2021.

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O. M. Mannion, K. S. Anderson, R. Betti, E. M. Campbell, D. Cao, C. J. Forrest, V. Yu. Glebov, V. N. Goncharov, V. Gopalaswamy, I. V. Igumenshchev, S. T. Ivancic, D. W. Jacobs-Perkins, J. P. Knauer, A. Lees, F. J. Marshall, Z. L. Mohamed, D. Patel, S. P. Regan, H. G. Rinderknecht, R. C. Shah, C. Stoeckl, W. Theobald, K. M. Woo, B. D. Appelbe, A. J. Crilly, J. Chittenden, W. Taitano, P. Adrian, J. A. Frenje, N. Kabadi, and M. Gatu Johnson, "Neutron Spectroscopy in Laser-Direct-Drive Inertial Confinement Fusion Implosions," presented at the Imperial College Seminar, virtual, 10 February 2021.

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R. B. Spielman, K. M. Woo, P. Tzeferacos, A. Reyes, C. R. Stillman, I. A. Begishev, C. Mileham, and J. Bromage, "Experiments and Modeling of SPL Absorption and Ablation with Picosecond Pulses," presented at the Pulsed-Laser Workshop, virtual, 18 February 2021.

J. D. Zuegel, "Ultra-Short Pulse Laser (USPL) Science and Technology at LLE," presented at the High Energy Laser Joint Technology Office Ultrashort Pulse Laser Workshop, virtual, 18 February 2021.

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The following presentations were made at Photonics West 2021, virtual, 6–11 March 2021:

I. A. Begishev, S.-W. Bahk, C. Dorrer, C. Feng, M. J. Guardalben, C. Jeon, R. G. Roides, M. Spilatro, B. Webb, D. Weiner, J. D. Zuegel, and J. Bromage, "A Highly Efficient, 10-J Output Signal Amplifier for Ultra-Intense All-OPCPA Systems."

T. Z. Kosc, T. J. Kessler, H. Huang, and S. G. Demos, "Minimizing Risk of Laser Damage due to Transverse Stimulated Raman Scattering in Large-Aperture KDP/DKDP Plates."

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The following presentations were made at the APS March Meeting, virtual, 15–19 March 2021:

S. X. Hu, P. M. Nilson, V. V. Karasiev, R. Paul, M. Ghosh, J. Hinz, S. Zhang, D. Mihaylov, V. Recoules, N. Brouwer, M. Torrent, I. E. Golovkin, and T. Walton, "Understanding Matter at Superdense and Warm Conditions."

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R. Paul, S. X. Hu, V. V. Karasiev, R. Dias, "Phase Diagram of Ternary Carbon-Sulfur-Hydrogen System up to 300 GPa."

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D. H. Froula, J. P. Palastro, R. Boni, M. Ambat, P. Franke, J. Oliver, D. Ramsey, J. L. Shaw, T. T. Simpson, D. Turnbull, S. Jolly, F. Quéré, C. Benedetti, E. Esarey, C. Geddes, C. Schroeder, R. Bingham, S. Stoller, N. Vafaei-Najafabadi, G. Gregori, B. Malaca, A. Helm, J. Vieira, A. DiPiazza, A. Howard, A. Arefiev, T. M. Antonsen Jr., and Z. Li, "Spatiotemporal Pulse Shaping for Plasma-Based Applications," presented at the Industrial Associates Bi-Annual Meeting, virtual, 17–19 March 2021.

E. M. Campbell, "Nuclear Energy, Today and Tomorrow," presented at Oklahoma University, virtual, 19 March 2021.

The following presentations were made at the 16th Direct-Drive and Fast-Ignition Workshop, virtual, 22–24 March 2021:

V. N. Goncharov, I. V. Igumenshchev, T. J. B. Collins, P. B. Radha, S. X. Hu, S. P. Regan, D. H. Froula, D. Harding, W. Theobald, M. J. Rosenberg, A. Shvydky, S. F. B. Morse, J. D. Zuegel, C. Dorrer, T. C. Sangster, and E. M. Campbell, "Expanding Ignition Parameter Space with the Dynamic Shell Formation Concept."

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The following presentations were made at the High Energy Density Plasma Diagnostics Course, virtual, 30 March–3 June 2021:

K. Churnetski, M. Michalko, and S. T. Ivancic, "The Single Line-of-Sight Time-Resolved X-Ray Imager."

S. T. Ivancic, "Omega Overview."

- S. T. Ivancic, "Test Project."
- S. T. Ivancic, "TRXI Install."

H. McClow, O. M. Mannion, and Z. Mohamed, "Neutron Time-of-Flight Diagnostics."

