A physics-based statistical mapping approach to understand and improve the performance of inertial confinement fusion implosions on OMEGA

by

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BIOGRAPHICAL SKETCH

The author was born in Tartu, Estonia. He moved to Tallinn, Estonia where he studied Engineering Physics at the Tallinn University of Technology. He graduated in 2013 with a Bachelor of Science degree. He worked as a computer-aided design engineer from 2013 to 2015 in Ehvert Mission Critical in Tallinn, Estonia. He moved to Rochester, New York, to study Mechanical Engineering in the University of Rochester. His research was supervised by Dr. Hussein Aluie and Dr. Riccardo Betti and was supported by the Horton fellowship from 2016 through 2022.
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ABSTRACT

Improving the performance of inertial confinement fusion implosions requires physics models that can accurately predict the response to changes in the experimental inputs. Good predictive capability has been demonstrated for the fusion yield using a statistical mapping of simulated outcomes to experimental data. [V. Gopalaswamy et al., Nature 565, 581-586 (2019)]. In this thesis, a physics-based statistical mapping approach is used to extract and quantify all the major sources of degradation of fusion yield and areal density in direct-drive implosions on the OMEGA Laser. The fusion yield is found to be dependent on the age of the deuterium tritium fill, the $\ell = 1$ asymmetry in the implosion core, the laser beam to target size ratio, and parameters related to the hydrodynamic stability. The inferred dependencies are compared to simulations and, in the case of the fill age dependency, a controlled set of dedicated experiments. Results from experiments designed with the aid of the statistical mapping model to optimize the target size on OMEGA are shown. A mapping of the areal density of OMEGA implosions reveals a strong degradation in highly convergent implosion designs. Dedicated experiments in both planar and spherical geometries were used to measure the kinetic energy of the shock release in order to assess the impact of shock release on the convergence of ICF implosions. No indication was found that the shock release could be responsible for the lack of convergence observed in experiments.
CONTRIBUTORS AND FUNDING SOURCES

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LIST OF SYMBOLS

\( v_a \)  Ablation velocity  
\( \Gamma \)  Adiabatic index  
\( \varepsilon_\alpha \)  Alpha particle energy  
\( \dot{W}_\alpha \)  Alpha-heating power  
\( \rho R \)  Areal density  
\( h_b \)  Bubble front coordinate  
\( \tau \)  Confinement time  
\( P_F \)  Fermi degenerate pressure  
\( \langle \sigma v \rangle \)  Fusion reactivity  
\( Y_{\text{no}\alpha} \)  Fusion yield without including contributions from alpha heating  
\( Y \)  Fusion yield  
\( P \)  Hotspot Pressure  
\( R_{\text{stag}} \)  Hotspot radius at stagnation  
\( R \)  Hotspot Radius  
\( V \)  Hotspot Volume  
\( \text{IFAR} \)  In-flight aspect ratio  
\( P_0 \)  Initial hotspot pressure  
\( R_0 \)  Initial target radius  
\( n \)  Ion number density  
\( T_i \)  Ion Temperature  
\( \gamma \)  Linear growth rate of Rayleigh-Taylor instability
MRS  Magnetic recoil spectrometer
\( \dot{m}_a \)  Mass ablation rate
NTOF  Neutron time-of-flight
\( \chi_{\text{no}\alpha} \)  Normalized Lawson parameter
\( \alpha_b \)  Parameter of non-linear Rayleigh-Taylor growth
\( k \)  Perturbation wavenumber
\( \dot{W}_{PdV} \)  Rate of PdV work
\( \dot{W}_{\text{cond}} \)  Rate of thermal conduction losses
\( g \)  Shell acceleration
\( \alpha_F \)  Shell adiabat
\( C_R \)  Shell convergence ratio
\( \rho_{sh} \)  Shell density
\( M_{sh} \)  Shell mass
\( P_{sh} \)  Shell pressure
\( \Delta \)  Shell thickness
\( U_{\text{shock}} \)  Shock velocity
\( I_\alpha \)  Short-wavelength Rayleigh-Taylor stability parameter
\( \ell \)  Spherical harmonic mode number
\( P_{\text{no}\alpha} \)  Stagnation pressure without including contributions from alpha heating
\( P_{stag} \)  Stagnation pressure
YOC  Yield over clean
CHAPTER 1
INTRODUCTION

The application of nuclear fusion to the production of clean energy continues to be one of humanity’s greatest challenges. The fusion reaction that offers the most promise, thanks to the favorable reaction cross-section, is that of deuterium (D) and tritium (T), which produces a neutron ($n$) and an $\alpha$-particle (a $^4\text{He}$ nucleus):

$$D + T \rightarrow n(14.03\text{MeV}) + \alpha(3.5\text{MeV}). \quad (1.1)$$

The key idea being that the charged $\alpha$-particles will stop and deposit their energy in the hot fusion plasma, thereby maintaining the fusion reactions. Whereas, the energetic neutrons will escape the plasma because of their lack of charge and ultimately the energy of the neutrons will be used in the form of heat to drive a steam cycle akin to a conventional power plant.

Over the years, many schemes have been proposed and worked on for laboratory scale confinement of the fusion fuel in the required extreme state of high temperature and pressure. The focus of this thesis is on inertial confinement fusion (ICF). Using powerful lasers to compress the fuel to the extreme conditions required for self-sustaining nuclear fusion reactions was first proposed in Ref. [1]. Focused laser light is used to drive the implosion of a millimeter-scale spherical capsule, typically made up of a layered hollow shell composed of an outer plastic (CH)
layer and an inner cryogenic DT ice layer with a low density DT vapor in the central region of the target. A schematic representation of an ICF target used in a high performance implosion experiment is shown in Fig. 1.1. The target is illuminated by a laser pulse with intensities reaching $\sim 10^{15}$ W/cm$^2$ over a nanosecond time-scale causing ablation of the outer layers of the capsule, which drives the implosion of the shell in a manner similar to rocket propulsion. The dense DT shell acts as a piston, reaching velocities of several hundred km/s compressing the low density central region, aptly called the hot spot, which can reach temperatures of around 5 keV or $5.8 \times 10^7$ K and pressures of 100 Gbar (for reference, the core of the sun has an approximate temperature of $1.6 \times 10^7$ K and a pressure of 250 Gbar [2]). The ultimate goal in ICF is achieving thermonuclear ignition of the target, where the alpha heating from the D+T fusion reactions in the hot spot that is being confined by the inertia of the dense shell will be sufficient to generate a propagating burn wave that will result in the burn-up of a significant fraction of the
DT fuel in the shell, thereby releasing more energy than was required to initiate the process.

Figure 1.2: The two main schemes employed in ICF experiments. (a) The direct-drive approach, where the laser beams arranged in a spherically symmetric configuration directly illuminate the target and (b) indirect-drive, where the laser beams are used to illuminate the inside of a cylindrical high-Z hohlraum producing x rays that drive the implosion of the capsule at the center of the hohlraum.

There are two main approaches to laser-driven ICF – the laser direct-drive (LDD) [3] approach pursued on the OMEGA laser [4] at the Laboratory for Laser Energetics (LLE) in Rochester, NY and laser indirect-drive (LID) [5] studied on the 2 MJ National Ignition Facility (NIF) [6] at the Lawrence Livermore National Laboratory. In the direct-drive scheme, as shown in Fig. 1.2(a), the target is directly illuminated by the laser
resulting in efficient coupling of the laser energy to the target, while in the indirect-drive approach shown in Fig. 1.2(b), the target is placed in a cylindrical radiation cavity or hohlraum made of a high-Z material, most commonly gold or depleted uranium. The laser is used to illuminate the walls of the hohlraum, producing x rays that in turn drive the implosion of the target. The main advantage of the indirect-drive approach is that the high frequency spatial modulations of the laser beam do not get imprinted onto the target but instead the target is driven by a very smooth illumination in the “bath” of x rays inside the hohlraum. The disadvantage of indirect-drive is much lower energy coupling efficiency as after the conversion into x rays only about 10% of the initial laser energy gets coupled into the target while around 70% of the laser energy is absorbed in the capsule in direct-drive.

The performance of an ICF implosion is measured by the Lawson triple product [7], which, when written in terms of parameters that are directly measurable in ICF experiments (see derivation in Sec. 1.1), is a function of the number of fusion reactions that occur in the target or fusion yield and areal density, which is a measure of the inertial confinement [8]. Recently, LID experiments on the 2 MJ NIF laser have demonstrated ignition by most metrics used in the literature [9]. The performance of the LDD experiments carried out on the much smaller 30 kJ OMEGA laser is usually calculated by projecting the results to a 2 MJ laser energy through a process called hydrodynamic scaling [10]. By these metrics, recent LDD implosions on OMEGA have achieved about 75% of the conditions required for ignition [11]. This suggests that only modest improvements are required to exceed the hydrodynamically
scaled ignition threshold. For instance, as shown in Ref. [11], a 30% increase in yield (to $2 \times 10^{14}$) at constant areal density and DT mass in OMEGA best-performing implosions would be sufficient to achieve conditions that hydrodynamically scale to a megajoule of fusion yield at 2 MJ of laser energy for symmetric illumination, while a 40% increase in both yield and areal density would scale to ignition.

The motivation behind the work on statistical modeling in this thesis is to develop an accurate predictive capability of the most important experimental outcomes of OMEGA ICF implosions in order to guide the efforts to incrementally improve the performance on OMEGA, with the goal of eventually reaching hydroequivalent ignition (i.e. attaining the conditions necessary for ignition when scaled to 2 MJ of laser energy). This requires that all of the dependencies are accounted for simultaneously because of the complex nonlinear nature of ICF implosions, where small changes to a few dominant parameters can easily hide the effect of the underlying weaker dependencies. Only when the dominant dependencies are accurately modeled can the subdominant dependencies be extracted from the experimental data which typically have simultaneous variations in a large number of parameters.

### 1.1 Hotspot ignition

The core conditions required for ignition can be derived by considering the hotspot energy balance. Certain simplifying assumptions will be made here to for the sake of simplicity of the derivation a more complete treatment of hotspot ignition can be found in Ref. [12]. Namely, the
alpha heating contribution to the hotspot pressure $P$ is taken to be negligible up to the time of peak compression $P(0) = P_{no\alpha}$, where the time of stagnation is denoted as time $t = 0$. Furthermore, the hotspot energy balance considered here does not include radiation losses, which are important in an ICF implosions but do not qualitatively alter the conclusions reached in this section. The hotspot energy balance then becomes

$$\frac{d}{dt} \left( \frac{3}{2} PV \right) = \dot{W}_\alpha + \dot{W}_{PdV} - \dot{W}_{\text{cond}},$$

(1.2)

where $P$ and $V$ are the pressure and volume of the hotspot, $\dot{W}_\alpha$ is the alpha-heating power, $\dot{W}_{PdV}$ is the rate of $PdV$ work done on the hotspot, which becomes a loss term after the hotspot begins expanding after shell stagnation and $\dot{W}_{\text{cond}}$ are the heat conduction losses. The $\dot{W}_{\text{cond}}$ loss can be neglected because the heat flux leaving the hotspot is recycled back into the hotspot through ablation of material from the inner surface of the shell [13].

Substituting in the alpha heating and $PdV$ work terms and writing in terms of the change in hotspot pressure the equation becomes

$$\frac{3}{2} \frac{dP}{dt} V = \varepsilon_\alpha \frac{n^2}{4} \langle \sigma v \rangle V - \frac{5}{2} P \frac{dV}{dt},$$

(1.3)

where $\varepsilon_\alpha = 3.5$ MeV is the energy of the alpha particles produced in the D+T fusion reactions, $n$ is the ion number density, and $\langle \sigma v \rangle$ is the reactivity of D+T fusion.

The inertial confinement time can be estimated by considering the dynamics of a hotspot confined by a shell with mass $M_{sh}$. From Newton’s
second law

\[ M_{\text{sh}} \ddot{R} = 4\pi P R^2, \]  

(1.4)

where \( R \) is the radius of the hotspot and \( \ddot{R} \) is the second derivative of the hotspot radius with respect to time. The characteristic confinement time can then be defined as

\[ \tau \sim \sqrt{\frac{R}{\ddot{R}}} = \sqrt{\frac{M_{\text{sh}}}{4\pi P R}}. \]  

(1.5)

By approximating the rate of change of hotspot volume using the confinement time as

\[ \frac{dV}{dt} \approx \frac{3}{5} \frac{V}{\tau}, \]  

(1.6)

the hotspot energy balance, after dividing by the hotspot volume, can be written as

\[ \frac{dP}{dt} = \frac{P^2}{24} \frac{\langle \sigma v \rangle}{T^2} - \frac{P}{\tau}, \]  

(1.7)

where the factor 3/5 in Eq. (1.6) was chosen such that the hotspot pressure in the absence of alpha heating decays over a timescale \( \tau \). Writing this in terms of the dimensionless parameters \( \hat{P} = P/P_{\text{no }\alpha} \) and \( \hat{t} = t/\tau \) results in

\[ \frac{d\hat{P}}{d\hat{t}} = \hat{P}^2 \chi_{\text{no }\alpha} - \hat{P}, \]  

(1.8)

where the normalized Lawson parameter [7] that determines the ignition condition is

\[ \chi_{\text{no }\alpha} = \frac{P_{\text{no }\alpha} \tau}{24 T^2 / \langle \varepsilon_\alpha \langle \sigma v \rangle \rangle} \]  

(1.9)

In a temperature range between 8 keV and 23 keV the fusion reactivity of D+T reactions scales like \( \langle \sigma v \rangle \sim T^2 \), which is not the best
approximation for typical ICF experiments with temperatures ranging between 3-5 keV. However, for the purposes of the current discussion this approximation greatly simplifies the math by ensuring that $\chi_{\text{no}\alpha}$ is constant over time.

Solving the ODE in Eq. (1.8) gives the hotspot pressure as

$$\hat{P} = \frac{1}{\chi_{\text{no}\alpha} + (1 - \chi_{\text{no}\alpha})e^t},$$  \hspace{1cm} (1.10)$$

where ignition of the hotspot occurs when $\chi_{\text{no}\alpha} > 1$ leading to $P \to \infty$. Because of this, the normalized Lawson parameter $\chi_{\text{no}\alpha}$ is useful as a metric of the overall performance of an ICF implosion when written in terms of experimentally measurable quantities.

The fusion yield $Y$ or number of fusion reactions in an ICF implosion is given by

$$Y = \int_0^\infty \frac{n^2}{4} \langle \sigma v \rangle V dt.$$  \hspace{1cm} (1.11)$$

Approximating the integral by using the burn-averaged hotspot volume $\langle V \rangle$ and writing in dimensionless form the fusion yield becomes

$$Y \approx \frac{3}{2} \frac{P_{\text{no}\alpha} \langle V \rangle}{\varepsilon_\alpha} \chi_{\text{no}\alpha} \int_0^\infty \hat{P}^2 d\hat{t}.$$  \hspace{1cm} (1.12)$$

Substituting in the solution of the hotspot pressure from Eq. (1.10) and taking the limit of $\chi_{\text{no}\alpha} \to 0$, where there is no contribution to the fusion yield from alpha heating (which is a good approximation for the ICF experiments on OMEGA), the fusion yield becomes

$$Y_{\text{no}\alpha} \approx \frac{3}{2} \frac{P_{\text{no}\alpha} \langle V \rangle}{\varepsilon_\alpha} \chi_{\text{no}\alpha}$$  \hspace{1cm} (1.13)$$
Using this relation between $P_{\text{no}\alpha}$ and yield and substituting the confinement time from Eq. (1.5) into the definition of $\chi_{\text{no}\alpha}$ in Eq. (1.9), the normalized Lawson parameter can be written in terms of the fusion yield, shell mass and hotspot radius as

$$\chi_{\text{no}\alpha} \sim \left( \frac{Y_{\text{no}\alpha} M_{\text{sh}}}{R^4} \right)^{1/3}. \quad (1.14)$$

In the limit of a thin shell, the shell mass can be approximated as $M_{\text{sh}} \approx 4\pi R^2 \rho R$, where $\rho R = \int_0^\infty \rho \, dr$ is the areal density of the shell. The Lawson parameter can then be rewritten as

$$\chi_{\text{no}\alpha} \sim \frac{Y_{\text{no}\alpha}^{1/3} (\rho R)^{2/3}}{M_{\text{sh}}^{1/3}} \quad (1.15)$$

The Lawson parameter in Eq. (1.15) provides a way to quantify the overall performance of an ICF implosion via measurements of the fusion yield and areal density (as well as the shell mass, which is not directly measurable in experiments but does not vary a great deal for implosion designs at a fixed scale). This fact directly motivates the search for predictive models of these two important observables in chapters 2 through 5 of this thesis.

1.2 Energy coupling in laser direct-drive

An example of the complex laser pulse shape used to drive the high performance implosion experiments on OMEGA is shown in Fig. 1.3. The main consideration that constrains the pulse shape design is the trade-off between compression and stability of the capsule. The details
of hydrodynamic instabilities in ICF implosions are discussed in Sec. 1.3; here, it suffices to note that the growth of shell perturbations can be mitigated by increasing the entropy of the shell as measured by the shell adiabat $\alpha_F = P_{sh}/P_F$, which is defined as the ratio of the shell pressure $P_{sh}$ to the Fermi degenerate pressure $P_F$. On the other hand, the highest possible target convergence can be achieved through adiabatic compression at the lowest possible shell adiabat. The front end of the laser pulse in Fig. 1.3 is therefore designed to launch a series of shocks into the target to set the adiabat of the shell that achieves the desired compromise between compression and stability. In the case of the pulse shape in Fig. 1.3 there are two shocks launched into the target – one launched by the initial Gaussian picket and a second shock launched by
the small “foot” feature just before the main drive. As the shocks break out of the inner surface of the shell, material is released into the central vapor region of the target. The kinetic energy of this release material limits the convergence of the target by increasing the initial hotspot pressure at the start of the deceleration phase. This phenomenon of shock release is studied experimentally in Ch. 6. The design of the main drive part of the pulse shape is governed by considerations related to maximizing the coupling efficiency to achieve a high implosion velocity while minimizing the adverse effects caused by hydrodynamic instabilities, laser-plasma interactions, coasting [14] and hot electron preheat [15].

During the initial picket of the laser pulse, the laser beams directly illuminate the outer surface of the target, causing the ablation of material off of the target surface. The ablated material quickly forms a “cloud” of plasma surrounding the target, where the density of the plasma decreases with distance from the target. At low plasma densities the laser light waves can propagate in the plasma where the laser energy is absorbed through collisional absorption (a process called inverse bremsstrahlung). The light waves propagate until they reach the critical density where the plasma frequency equals the laser frequency and all the incoming light gets fully reflected. The region of the plasma with a density below the critical density is referred to as the corona.

Beyond the critical density, the absorbed energy is transported to the ablation surface of the target by means of thermal conduction. This region of the plasma between the ablation surface and the critical surface is called the conduction zone. The temperature gradients in the conduc-
tion zone can become extremely large especially at early time when the conduction zone scale length is very short. The classical thermal transport models break down under these conditions, predicting unphysically large thermal fluxes. Instead, flux limited thermal transport or more accurate thermal transport models that include nonlocal effects [16] must be used to model the thermal conduction.

Figure 1.4: The basic mechanism of cross-beam energy transfer. The refracted outgoing beam depletes energy from the incoming beam thereby reducing laser absorption.

Energy coupling in ICF is further complicated by the process of cross-beam energy transfer (CBET) [17, 18], which is the transfer of energy between laser beams in the coronal plasma. Figure 1.4 shows the mechanism of CBET in direct-drive experiments. The rays in the outer portion of an incoming beam are refracted in the plasma thereby missing the target. These outgoing rays then interacts in the coronal plasma with the rays in the central region of an adjacent incoming beam, depleting their energy and thereby reducing the coupling efficiency. The reduced
absorption caused by CBET can be significant in implosion designs with a high degree of beam overlap. On the other hand, beam overlap is desired to achieve a more uniform illumination intensity on target. The size of the beam spots in relation to the target size must therefore be carefully chosen to optimize the trade-offs between uniformity and coupling efficiency.

1.3 Hydrodynamic instabilities in ICF implosions

In an ICF implosion, the target is subject to a number of hydrodynamic instabilities causing perturbations of the shell surfaces to grow, which can compromise the integrity of the shell as a piston compressing the hotspot and lead to inefficient conversion of the shell kinetic energy into hotspot internal energy. Most notably, the Rayleigh-Taylor (RT) instability\(^1\), which occurs when a heavy fluid is supported by a lighter fluid, can strongly degrade the performance of ICF implosions because of exponential growth of perturbation in the linear phase – the amplitude of small perturbations grows like \(e^{\gamma t}\), where \(\gamma\) is the growth rate. The fact that the shell is RT unstable is easy to intuitively understand by considering a reference frame co-moving with the imploding shell, where the outer surface of the dense shell is supported by the low density plasma of the conduction zone during shell acceleration. Whereas in the deceleration phase, the perturbations on the inner surface, seeded via feed-through from the outer surface, undergo RT growth as the dense shell is supported by the low density hotspot.

In ICF implosions, the classical RT growth rates are modified by the
effects of mass ablation. A simple approximation of the growth rate of a perturbation with wavenumber $k$ in the linear phase can be written as

$$\gamma \approx \alpha \sqrt{kg} - \beta kv_a,$$

(1.16)

where the values of the constants are $\alpha = 0.9$ and $\beta = 2.5-3$ and $v_a$ is the velocity of the ablation front moving through the shell and $g$ is the acceleration of the shell [21].

In comparison to the classical case, the ablative RT instability includes a stabilizing effect dependent on $v_a = \dot{m}_a/\rho_{sh}$ given by the ratio of the mass ablation rate and the shell density. The density of the shell for a given ablation pressure $P_A$ depends on the shell adiabat as $\alpha_F \propto P_A/\rho_{sh}^\Gamma$ and therefore $v_a \propto \alpha_F^{3/5}$ assuming an adiabatic index of $\Gamma = 5/3$.

In the nonlinear phase the RT bubble front $h_b$ advances according to the formula

$$h_b \approx \alpha_b Ag t^2,$$

(1.17)

where $A$ is the Atwood number and the parameter $\alpha_b$ is a function of the ablation velocity $v_a$ and therefore the shell adiabat, linearly decreasing with increasing $v_a$.

It is common to partition the perturbations in ICF implosions into categories of low, mid and high modes based on the qualitative differences in the dominant processes that govern the evolution of the perturbations.
1.3.1 Low mode seeds on OMEGA

Low modes in ICF are typically defined as those with wavelengths less than the shell radius $R$. In spherical geometry the wavelength $\lambda$ of a mode with wavenumber $\ell$ is $\lambda = 2\pi R/\ell$ and therefore perturbations with $\ell < 6$ are considered low modes. The linear RT growth rates for low modes are small (see the dependency on $k$ in Eq. (1.16) and therefore the amplitudes of the low mode perturbations at stagnation are dominated by convergence effect (Bell-Plesset growth) [22, 23].

On OMEGA the low mode perturbations are seeded by target offset from the target chamber center, laser beam mispointing and power imbalances between the beams. Additionally, the capsule wall as well as the DT ice layer can exhibit low mode asymmetries. Figure 1.5(a) shows an illustration of target offset, which produces a pure $\ell = 1$ asymmetry.
in the illumination intensity on the capsule surface. Figure 1.5(b) shows the signature of an $\ell = 1$ mode in the x-ray self emission image from a cryogenic OMEGA implosion that had a large target offset.

### 1.3.2 Mid mode seeds on OMEGA

Mid mode perturbations are of intermediate wavenumber that are affected both by the growth of RT instability as well as convergence effects in the deceleration phase. The range of wavenumbers that comprise the mid modes is typically taken to be $6 \leq \ell \lesssim 30$ where the wavelengths are short enough that beyond the early perturbations seeded during the picket and foot, the conduction zone smooths out the mid modes in the laser illumination. At the same time the wavelength of the perturbations is large enough that they feed through from the ablation front to the inner surface of the shell as the amplitude of the feed-through decays at a distance $x$ measured from the ablation surface as $\sim e^{-\ell x/R}$.

![Figure 1.6](image)

Figure 1.6: (a) Variation in laser illumination intensity caused by the OMEGA beam port geometry as a percentage of the mean illumination intensity calculated via hard-sphere projection for a case of $R_b/R_t = 0.75$. (b) Mode spectrum of the beam mode nonuniformity showing dominant contributions from the $\ell = 10$ and $\ell = 6$ modes.
Mid mode perturbations are seeded on OMEGA mainly by the laser illumination asymmetry caused by the beam geometry of the 60 beams arranged in a “soccer ball” pattern of pentagons and hexagons with the beam spots located at the vertices of the polygons. While the beam configuration on OMEGA has been designed to produce a very uniform overlapped illumination on target, some spatial asymmetry remains because of the finite number and size of the beams. A hard sphere projection of the nonuniformity in the illumination pattern created by the OMEGA beam geometry is shown in Fig. 1.6(a). The spherical harmonic mode spectrum of this asymmetry dominated by the $\ell = 10$ and $\ell = 6$ mid modes is shown in Fig. 1.6(b). The size of the laser beam spots is determined by the distributed phase plates [24]; the ratio of the beam radius to target radius $R_b/R_t$ can be used to characterize the degree of beam overlap and therefore the amplitude of the mid mode perturbation caused by the beam geometry.

1.3.3 High mode seeds on OMEGA

Perturbations with wavenumbers $\ell > 30$ are referred to as high modes in ICF. On OMEGA, the high modes are seeded by imprinting of the laser speckle pattern on the capsule surface at early time, as illustrated in the sketch in Fig. 1.7(a), as well as roughness of the capsule surface. Once a conduction zone has formed during the first picket in the laser pulse, the high frequency spatial modulations in the intensity across the beam spots are smoothed out and the nonuniformity amplitude is determined by the RT growth of the initial imprinted perturbation.

Short wavelength modes reach the nonlinear phase quickly and there-
Figure 1.7: Illustrative sketch of the process of laser imprinting. At early time, before a conduction zone has formed, the high frequency spatial modulation in the laser beam intensity seeds nonuniformities on the outer surface of the target.

Therefore in terms of their effect on ICF performance it is useful to consider the bubble front penetration fraction \( h_b / \Delta \), where \( \Delta \) is the in-flight shell thickness. It is evident from Eq. (1.17) that the bubble front advances proportional to the distance traveled by the shell \( gt^2 = 2(R_0 - R) \), where \( R_0 \) is the initial shell radius. The bubble front penetration fraction then becomes

\[
\frac{h_b}{\Delta} \approx 2\alpha_b \frac{R_0 - R}{\Delta}. \tag{1.18}
\]

This can be written in terms of the in-flight aspect ratio IFAR = \( \frac{R}{\Delta} \)

\[
\frac{h_b}{\Delta} \approx 2\alpha_b \text{IFAR}(C_R - 1), \tag{1.19}
\]

where \( C_R = \frac{R_0}{R} \) is the shell convergence ratio.

The exact relationship between IFAR and \( C_R \) can be somewhat complicated and dependent on the details of the laser pulse shape, however for the purposes of this discussion a good enough approximation is pro-
vided by the scaling \( \text{IFAR} \sim \text{IFAR}_{\text{max}}/C_R^2 \), where \( \text{IFAR}_{\text{max}} \) is the maximum aspect ratio at the start of the acceleration phase.

The bubble front penetration fraction then becomes

\[
\frac{h_b}{\Delta} \approx 2\alpha_b \text{IFAR}_{\text{max}} \frac{1}{C_R^2} (C_R - 1),
\]

(1.20)

where the maximum of the time-dependent term on the right-hand-side occurs at \( C_R = 2 \). From this relation, the critical value of \( \text{IFAR}_{\text{max}} \) for a given growth parameter \( \alpha_b \) can be found, which ensures that the bubble front does not penetrate through the shell, i.e. \( h_b/\Delta < 1 \).

As was mention earlier, the \( \alpha_b \) parameter that governs the nonlinear growth of ablative RT instability depends on the shell adiabat \( \alpha_F \) and therefore the bubble front penetration fraction depends on the combination of \( \text{IFAR}_{\text{max}} \) and \( \alpha_F \). The stability parameter \( I_\alpha \equiv \frac{\alpha_F}{3\text{IFAR}_{\text{max}}/20} \) was found in Ref. [25] and [26] to be a good parametrization of the effect of short wavelength modes on the performance of ICF implosions.

### 1.4 Filling and layering of ICF targets

The material composition of the DT fuel and the details of the target filling and layering procedures have been demonstrated to have an effect on the performance of ICF implosions [27]. The plastic shells used for cryogenic ICF experiments at OMEGA are permeatively filled with DT gas several days to weeks before shot time. During the filling procedure, the entire shell is subject to a high dose of radiation from the \( \beta \) decay of tritium both inside the shell and in the fill station. Furthermore, the inner shell surface also absorbs radiation from the DT inside the shell
after the target has been transferred out of the fill station [28]. The radiation can break up the chemical bonds in the plastic shell, resulting in defects that act as seeds to hydrodynamic instabilities. Such defects have been observed in plastic capsules after being exposed to radiation in combination with the stresses that the target experiences during the permeation fill procedure [29, 30]. The radiation dosage and, therefore the severity of the shell defects, depend on the tritium concentration in the fuel. As a result, the yield degradation from the instability growth cannot be modeled with just the outputs of 1-D codes but requires an extra parameter to account for the varying seeds.

After a target has been filled with DT gas, it is cooled down and a DT ice crystal is grown on the inner surface of the plastic capsule. The layer quality of the target deteriorates over time with the formation of so-called “baseball seams” in the ice, necessitating relayering of the target leading up to the shot time, which involves melting the existing layer and growing a fresh, more-uniform, ice layer in its place. As the tritium atoms in the target are continuously undergoing $\beta$ decay into $^3$He, the $^3$He atoms accumulate inside the target. Since the targets are often relayered several times, with the last relayering typically occurring within one day of the shot, the $^3$He, which has a lower freezing temperature than DT, accumulates in the vapor region of the target even though most of it is produced in the significantly more massive ice layer. The introduction of $^3$He in the vapor region reduces the fusion yield of an ICF implosion through two separate physical mechanisms: (1) increased radiation losses in the hot spot caused by the presence of the higher-Z $^3$He, and (2) increased initial density in the vapor region, leading to reduced
hot-spot pressure at stagnation[31].

Figure 1.8 shows the simulated yield degradation caused by the accumulation of $^3$He in the central vapor region as a function of fill age—the time between filling the target with DT and shot time. One-dimensional LILAC simulations of the high-performance OMEGA shot 90288 were carried out with varying fill ages and ice tritium concentrations $\theta_T$. The number of $^3$He atoms produced by the decay of tritium in both the ice layer and the vapor region over the course of the fill age was calculated and all of the $^3$He was assumed to be accumulated in the vapor as per the discussion above. As expected, the simulated yield degradation monotonically increases with longer fill ages and higher tritium concentrations.
1.5 Outline of the thesis

The ICF implosions are designed using sophisticated radiation-hydrodynamics simulation codes that include models for most of the important physics mentioned above, yet neither 1-D nor 3-D simulations have demonstrated the kind of predictive capability over any large ensemble of implosions that is sufficient for guided design improvements. That is to say that the simulations alone are unable to reliably suggest even small changes to current best-performing designs that would lead to incremental performance improvements toward the eventual goal of hydroequivalent ignition. Recently, statistical methods that aim to bridge the gap between simulation predictions and experimental results have been used to demonstrate excellent prediction accuracy and have led to a substantial increase in the fusion yield on OMEGA [11]. In Ch. 2 of this thesis, a physics-based statistical mapping model for the fusion yield is constructed that incorporates the outputs of 1-D simulations as well as certain inputs to experiments that are not included in the codes. The model is able to accurately predict the fusion yield over a large database of ICF implosions performed on OMEGA.

The main feature of the mapping model shown here is that all the factors that determine the predicted fusion yield can be expressed in an intuitive form, where each term in the model corresponds to a specific physical process that is degrading the yield with respect to simulations. In Ch. 3, this property of the model is used to identify and quantify all the major mechanisms of fusion yield degradation on OMEGA. The mapping model is used to extract each individual dependency from the database
so that they can be compared to simulations in isolation, thereby elucidating the shortcomings of current ICF codes.

In Ch. 4, results from recent experimental ICF campaigns are shown, where physics insights gained from the application of the statistical mapping model were used to improve the fusion yield of the highest-performing implosions on OMEGA.

In addition to the fusion yield, the Lawson parameter introduced in Sec. 1.1 that determines the performance of an ICF implosion depends on the areal density. In Ch. 5 the statistical mapping method is applied to construct a prediction model of the areal density in OMEGA implosions. The difficulty in modeling the areal density highlights the pressing need to understand the degradation of ICF performance in low adiabat implosions.

In Ch. 6, one of the possible explanations to poor performance at low adiabat - the process of shock release - is investigated in detail. Results are shown from a dedicated set of experiments that were designed to measure the kinetic energy of the shock release in order to validate the radiation-hydrodynamics codes used as inputs to the mapping model from Ch. 2 and 5.

Finally, concluding remarks are given in Ch. 7.
CHAPTER 2

A PHYSICS-BASED STATISTICAL MAPPING MODEL

2.1 The statistical mapping method

The statistical mapping method for building predictive models for ICF was first introduced in Ref. [11], where the framework for using outputs from 1-D simulations to model the effects of constant systematic nonuniformities was established. Here, a new version of the model is described, which allows for the existence of experimental inputs that are not included in simulations (e.g., details of the target-filling procedure) as well as variable nonuniformity seeds. This new model also provides physics understanding and quantification of each degradation mechanism affecting OMEGA implosions.

The core idea of the statistical mapping method is to find a mapping function $f_{\text{map}}$ that relates a set of simulation outputs $\mathbf{O}_{\text{sim}}$ to an experimental observable $\mathbf{O}_{\text{exp}}$, such as the fusion yield or areal density, over a large database of experiments:

$$f_{\text{map}} : \mathbf{O}_{\text{sim}} \rightarrow \mathbf{O}_{\text{exp}},$$

(2.1)

where the boldface notation indicates that the mapping is performed over vectors of observables from all shots in the database. This function can then be used to understand the causes of the variation of the observables measured in experiments and to search the available parameter
space in simulations to design implosions with improved experimental performance.

In the case where the simulations accurately capture all the relevant physics and all the experimental inputs are included, $f_{\text{map}}$ becomes the identity matrix and $O_{\text{sim}} = O_{\text{exp}}$. In practice, the physics models are always inaccurate to some degree and not all inputs are always perfectly characterized. The statistical mapping method allows for these discrepancies through increased complexity of $f_{\text{map}}$. As the complexity of the mapping function increases, so does the size of the training set required to achieve an out-of-sample prediction error similar to the in-sample training error \[32\]. This is problematic for ICF applications because of the scarcity of experimental data caused by the low shot rate of laser facilities. One possibility for dealing with this limitation is to train an initial neural network on a very large simulation database and then use transfer learning to re-train a subset of the neural network layers on the small experimental dataset \[33, 34\]. Here, as in Ref. \[11\], the mapping relies on the simplicity of the mapping functions and the fact that the dominant physics is captured reasonably well in the 1-D radiation-hydrodynamics code LILAC \[35\], which uses the first-principles equation of state \[36, 37\] and includes models for nonlocal thermal transport \[16\], multigroup radiation diffusion and cross-beam energy transfer (CBET) \[17, 18\]. In particular, the fusion yield is dominated by a very strong dependence on implosion velocity, which is well predicted by LILAC as indicated by shell trajectory measurements \[38\].

Using 3-D simulations should, in principle, result in a less-complex $f_{\text{map}}$ because nonuniformities are included. However, good agreement
between measured and simulated asymmetry signatures over a large database of implosions using self-consistent nonuniformity seeds has not yet been demonstrated. Even if that were the case, the computational cost of building a high-fidelity 3-D simulation database of \( \sim 200 \) implosions is prohibitively expensive. Furthermore, in practice, the simulation database is not static since any consideration of additional input parameters, improved physics models, or improved characterizations of existing input parameters requires a new iteration of the entire database. The advantage of using 1-D simulations is the low computational cost that allows iterative improvements to the simulations. Once a model with sufficient prediction accuracy has been built, the low cost of 1-D simulations allows for efficient scans of the parameter space in order to find new implosion designs with improved performance.

The observables in 1-D simulations simply depend on the 1-D inputs \( I_{1D} \) to the code (i.e., the target specifications and laser pulse shape)

\[
O_{1D}^{\text{sim}} = f_{\text{sim}}(I_{1D}). \tag{2.2}
\]

The experiments depend on these same 1-D inputs but also extra input parameters that are not included in the simulations \( I_{\text{sys,other}}^{\text{sys}} \). In addition, 3-D nonuniformity seeds are always present in experiments, which are categorized into three groups based on how they are handled in the statistical mapping framework: (1) \( S_{3D}^{\text{const}} \), systematic seeds that always have the same magnitude in all shots in the database (examples include the stalk holding the target, the target surface roughness, DT ice layer nonuniformity, and laser beam power balance); (2) \( S_{3D}^{\text{var}} \), sys-
tematic seeds that can be changed from shot to shot (e.g., the laser beam spot size $R_b$); and (3) $S_{3D}^{\text{ran}}$, random seeds that can be different in each shot (e.g., target offset from the center of the target chamber and laser beam mispointing). The experimental observables are therefore determined by

$$O_{\text{exp}} = f_{\text{exp}} (I_{1D}, I_{\text{sys}}^{\text{other}}, S_{3D}^{\text{const}}, S_{3D}^{\text{var}}, S_{3D}^{\text{ran}}).$$

(2.3)

If the simulated outputs are uniquely determined by the simulation inputs then the function in Eq. (2.2) can be inverted $I_{1D} = f_{\text{sim}}^{-1} (O_{1D}^{\text{sim}})$ and replaced into Eq. (2.3) since the 1-D inputs to both the simulations and experiments are identical

$$O_{\text{exp}} = f_{\text{exp}} [f_{\text{sim}}^{-1} (O_{1D}^{\text{sim}}), I_{\text{sys}}^{\text{other}}, S_{3D}^{\text{const}}, S_{3D}^{\text{var}}, S_{3D}^{\text{ran}}].$$

(2.4)

A mapping function can now be defined that maps the simulated outputs and other experimental inputs to the experimental observables

$$O_{\text{exp}} = f_{\text{map}} (O_{1D}^{\text{sim}}, I_{\text{sys}}^{\text{other}}, S_{3D}^{\text{var}}, S_{3D}^{\text{ran}}),$$

(2.5)

where the constant systematic seeds $S_{3D}^{\text{const}}$ have been absorbed into $f_{\text{map}}$. Equation Eq. (2.5) holds regardless of the accuracy of the physics models in the 1-D codes; however, inaccuracies in the codes will lead to increased complexity of the mapping function $f_{\text{map}}$.

The experimental observables that are at the focus of this work are the fusion yield $Y_{\text{exp}}$ and areal density $\rho R_{\text{exp}}$ which together make up the Lawson parameter $\chi$ used as the metric of ICF implosion performance.
In the following section, a mapping model is constructed for the fusion yield over a large database of OMEGA experiments, whereas the areal density is modeled in Ch. 5.

2.2 Mapping the fusion yield of OMEGA implosions

2.2.1 Selecting the model parameters

In order to find $f_{\text{map}}$, the major parameters that control the fusion yield in all the categories of Eq. (2.5) must be selected. The experimental inputs $I_{\text{sys}}^{\text{other}}$ that have been found to have a significant effect on the fusion yield on OMEGA are parameters related to the target filling procedure: (1) $\theta_T$ and (2) $\theta_D \equiv (1 - \theta_T)$—the concentrations of tritium and deuterium, respectively, in the DT ice layer and; (3) $\tau_{\text{fill}}$—the DT fill age defined as the time between filling the target and shot time. The physical justification of the yield dependence on these parameters is given in Sec. 3.1, where it is shown that their effect can be combined into a single parameter $\xi(\tau_{\text{fill}}, \theta_T)$. This is by no means a full list of all the extra inputs that influence the implosion experiments, but these appear to have the largest impact on the fusion yield. Any extra parameters that have an effect on the yield, which have been excluded at this stage of the analysis, will result in increased prediction error of the constructed mapping model.

The variable systematic nonuniformity seed ($S_{\text{var}}^{3D}$) included in the yield model is the ratio $R_b/R_t$ of the beam spot size $R_b$ to the target radius $R_t$, which will be explored in more detail in Sec. 3.3 and is related to the nonuniformities seeded by the beam port geometry and laser imprint [39], as well as laser-plasma interactions in the corona. Another exam-
ple of a parameter that could be considered a variable nonuniformity seed would be the bandwidth of smoothing by spectral dispersion (SSD) [40]. This thesis is concerned with high-performance implosion designs that utilize the maximum available SSD bandwidth. For this reason, all the lower bandwidth shots have been excluded from the database, thereby making the SSD bandwidth a constant nonuniformity seed included in $S_{3D}^{\text{const}}$. If one were to use the statistical mapping approach to quantify the effect of SSD on fusion yield, a parameter that relates the SSD bandwidth to the fusion yield would have to be included in such a model.

Random seeds $S_{3D}^{\text{ran}}$ change from shot to shot relatively frequently by such a degree that they can have significant effect on the performance of repeated shots of the same implosion design. On OMEGA, the random seeds that have been observed to impact the fusion yield are laser-beam pointing and target offset. Since there is currently no accurate method to relate the magnitudes of these perturbations to their effect on the fusion yield, these nonuniformities have to be accounted for post-shot by including measured asymmetry signatures in the mapping model. This can be accomplished by using the inferred ion-temperature asymmetry $\hat{R}_T$—a signature of the dominant $\ell = 1$ mode generated by these particular seeds [41, 42]. The exact definition of $\hat{R}_T$ is given in Sec. 3.2. No further random asymmetry signatures have been identified at OMEGA; therefore, all other nonuniformity seeds are considered to be systematic. The inclusion of these random seeds in the statistical mapping model greatly expands the available training data and enables fair comparison between existing implosion designs; whereas, new designs can
simply be assumed to nominally have no degradation caused by random effects and corrected post-shot, if necessary.

2.2.2 Fitting to the OMEGA database

With the experimental inputs and nonuniformity seeds identified, the mapping relation for the fusion yield becomes

\[ Y_{\text{exp}} \approx f_{\text{map}} \left[ O_{1D}^{\text{sim}}, \xi(\tau_{\text{fill}}, \theta_T), \frac{R_b}{R_t}, \tilde{R}_T \right]. \tag{2.6} \]

An often used metric in ICF research is the yield-over-clean or the ratio of measured fusion yield to the simulated yield \( YOC_{\text{exp}} \equiv \frac{Y_{\text{exp}}}{Y_{\text{sim}}} \). Writing the mapping relation \([\text{Eq. (2.6)}]\) in terms of \( YOC_{\text{exp}} \) lends itself to an intuitive interpretation of the measured yield as an ideal implosion that has been degraded by a number of different mechanisms:

\[ YOC_{\text{exp}} = \bar{f}_{\text{map}} \left[ O_{1D}^{\text{sim}}, \xi(\tau_{\text{fill}}, \theta_T), \frac{R_b}{R_t}, \tilde{R}_T \right]. \tag{2.7} \]

Note that since \( YOC_{\text{exp}} \) is defined in terms of the simulated 1-D yield, the “degradation” with respect to the 1-D code is not necessarily caused only by 3-D effects but can also be caused by inaccurate 1-D physics models in the simulation.

Approximating the mapping function with power laws, as done in Ref. [11], yields

\[ YOC_{\text{exp}} \approx \mu_0 \prod_{i=1}^{N} \left( O_{1D}^{\text{sim}} \right)_i^{\mu_i} [\xi(\tau_{\text{fill}}, \theta_T)]^{\mu_i N+1} \left( \frac{R_b}{R_t} \right)^{\mu_{N+2}} \tilde{R}_T^{\mu_{N+3}}, \tag{2.8} \]

where \( N \) is the number of simulated observables included in the model.
The exponents $\mu_i$ are obtained by nonlinear regression over a large database of ICF implosions.

An accurate fit of $\text{YOC}^{\text{exp}}$ with separable power laws as in Eq. (2.8) enables a very powerful application of the statistical mapping method. The model can be rewritten in terms of a product of individual degradations, each driven by a specific physical mechanism. These individual degradation mechanisms can then be studied in isolation by comparing simulated degradation to the experimentally inferred degradation, and new experimental campaigns can be designed with the aid of the mapping model to precisely isolate the degradation mechanisms in experiments. The final form of the mapping model is written as

$$\text{YOC}^{\text{exp}} \approx \text{YOC}_h \text{YOC}_f \text{YOC}_b \text{YOC}_{\ell = 1} \text{YOC}_{\text{res}}, \quad (2.9)$$

where $\text{YOC}_h$ represents the yield degradation caused by hydrodynamic effects, $\text{YOC}_f$ accounts for the effect of fill age and target composition, the effect of finite beam size is denoted as $\text{YOC}_b$ and the degradation caused by random $\ell = 1$ asymmetry is included in $\text{YOC}_{\ell = 1}$. The residual scaling term $\text{YOC}_{\text{res}}$ is a weak dependence that is necessary to explain the performance of a handful of hydro-equivalently downscaled implosions with thin capsule wall thicknesses [43, 44]. The $\text{YOC}_{\text{res}}$ term has no effect on the high-performance implosions that are in the focus in this work and therefore is not discussed here.

Figure 2.1 shows the overall accuracy of the mapping model in Eq. (2.9), comparing the predicted $\text{YOC}$ with the measured $\text{YOC}^{\text{exp}}$. The power-law exponents along with the 95% confidence intervals of all the map-
Figure 2.1: Experimental yield-over-clean (YOC)-measured neutron yield normalized to the 1-D simulated yield vs. the predicted YOC from the statistical mapping model [Eq. (2.9)]. The dashed line acts as a guide to the eye and indicates a perfect prediction accuracy.

Fitting parameters are shown in Table 2.1 and exact definitions of all YOC terms are given in Ch. 3.
Table 2.1: Power indices and confidence intervals for all degradation terms, as well as values of the critical parameters, resulting from fitting the model in Eq. (2.9) to the OMEGA database.

<table>
<thead>
<tr>
<th>Parameter</th>
<th>Coefficient</th>
<th>95% confidence interval</th>
</tr>
</thead>
<tbody>
<tr>
<td>$I_{\mu_1}$</td>
<td>$\mu_{1,&lt;} = 0.97$</td>
<td>$\mu_{1,&lt;} = 0.80 \ldots 1.13$</td>
</tr>
<tr>
<td></td>
<td>$\mu_{1,&gt;} = 0.42$</td>
<td>$\mu_{1,&gt;} = 0.37 \ldots 0.47$</td>
</tr>
<tr>
<td></td>
<td>$I_{\text{crit}} = 0.8$</td>
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</tr>
<tr>
<td></td>
<td>$I_{\text{cut}} = 2.4$</td>
<td></td>
</tr>
<tr>
<td>$C_{R}^{\mu_2}$</td>
<td>$\mu_2 = -0.78$</td>
<td>$\mu_2 = -0.87 \ldots -0.68$</td>
</tr>
<tr>
<td>$\bar{D}_{\mu_3}$</td>
<td>$\mu_3 = -2.95$</td>
<td>$\mu_3 = -3.82 \ldots -2.10$</td>
</tr>
<tr>
<td>$\left(YOC_{\text{He}}^{\text{sim}}\right)^{\mu_4}$</td>
<td>$\mu_4 = 1.30$</td>
<td>$\mu_4 = 1.19 \ldots 1.40$</td>
</tr>
<tr>
<td>$\bar{R}_{T}^{\mu_5}$</td>
<td>$\mu_5 = -1.37$</td>
<td>$\mu_5 = -1.53 \ldots -1.21$</td>
</tr>
<tr>
<td>$R_{T}^{\text{min}} = 1.14$</td>
<td></td>
<td></td>
</tr>
<tr>
<td>$\left(R_{b/t}\right)^{\mu_6}$</td>
<td>$\mu_6 = 2.59$</td>
<td>$\mu_6 = 2.25 \ldots 2.93$</td>
</tr>
<tr>
<td>$\left(\bar{R}_{b/t}\right)^{\mu_7}$</td>
<td>$\mu_7 = 2.00$</td>
<td>$\mu_7 = 1.27 \ldots 2.75$</td>
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<tr>
<td>$R_{b/t}^{\text{crit}} = 0.86$</td>
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CHAPTER 3
INFERRED DEPENDENCIES OF FUSION YIELD IN OMEGA EXPERIMENTS

One of the key elements of the statistical mapping method is that all of the major degradation terms are fit to the experimental data simultaneously. This is necessary when dealing with the complex datasets of ICF experiments, which always include varying contributions from several different degradation sources. This is often true even for dedicated campaigns where the experiments have been carefully constructed to vary only one design parameter at a time because different degrading effects often depend on the same design parameters. For example, changing the shell adiabat to control the degree of ablative stabilization of short-wavelength perturbations also influences the behavior of low and mid modes by changing the target convergence. Furthermore, the dominant degradation effects can vary considerably through the shot-to-shot variability of the facility and therefore must be accurately modeled so as to not overshadow the effects of weaker degradations.

Once a model of the type in Eq. (2.9) has been constructed by mapping all of the terms to a large database of experiments, the power-law formulation employed here lends itself to a very straightforward manner of extracting the individual dependencies from the data,

$$YOC_{j}^{exp} = \frac{YOC_{i}^{exp}}{\prod_{i \neq j} YOC_{i}},$$

(3.1)
where \( \text{YOC}^{\text{exp}}_j \) is the experimentally inferred yield-over-clean caused only by the degradation source denoted by \( j \). In the following subsections, Eq. (3.1) is used to isolate and visualize the dependencies of each degradation in Eq. (2.9) and to compare with simulations.

### 3.1 Filling and layering targets

The parameter used in the statistical mapping model to account for the effect of target filling and layering effects is the simulated yield-over-clean caused by \(^3\text{He} \) accumulation, which is defined as:

\[
\xi(\tau_{\text{fill}}, \theta_T) \equiv \text{YOC}^{\text{sim}}_{\text{He}} = \frac{Y^{\text{sim}}_{1\text{D}, \text{He}}}{Y^{\text{sim}}_{1\text{D}}},
\]

(3.2)

where \( Y^{\text{sim}}_{1\text{D}} \) is the fusion yield in a simulation with no \(^3\text{He} \) contamination and \( Y^{\text{sim}}_{1\text{D}, \text{He}} \) is the yield in a simulation with the \(^3\text{He} \) included in the vapor.

Figure 3.1 shows the experimentally inferred yield degradation caused by \(^3\text{He} \) accumulation, extracted from the experimental database using the formula in Eq. (3.1), compared to the predicted degradation \( \text{YOC}_f \sim (\text{YOC}^{\text{sim}}_{\text{He}})^{1.30} \). The experimental data are in good agreement with the prediction; however, the fact that the confidence interval of the exponent to \( \text{YOC}^{\text{sim}}_{\text{He}} \) does not include unity suggests a stronger degradation with fill age than seen in 1-D simulations. This could be caused by a stronger-than-predicted effect of the \(^3\text{He} \) accumulation or the existence of a further degradation mechanism that depends on the fill age such as the damage to the ablator discussed previously.

The predictive capability of the \( \text{YOC}^{\text{sim}}_{\text{He}} \) term is best demonstrated by the two extremely long-fill-age shots 94717 and 97141, indicated with
Figure 3.1: Experimentally inferred degradation caused by accumulation of $^3\text{He}$ in the vapor, extracted from the data using Eq. (3.1). The red diamonds show results from two experiments with extremely long fill ages. The dashed curve corresponds to the mapping model relation in Eq. (2.9).

red diamonds in Fig. 3.1, with fill ages of 45 and 97 days, respectively. The experimentally inferred yield-over-clean caused by $^3\text{He}$ accumulation for these shots is $40 \pm 3\%$ and $25 \pm 2\%$, compared to predicted values of $34\%$ and $24\%$. Accurate prediction of shots that lie so far outside the usual range of fill ages increases confidence that $\text{YOC}_{\text{He}}^{\text{sim}}$ correctly captures the physics of the degradation mechanism.

3.2 $\ell = 1$ Asymmetry

The yield degradation caused by $\ell = 1$ asymmetry generated by random nonuniformities is represented by the $\text{YOC}_{\ell=1}$ term in the mapping model Eq. (2.9). On OMEGA, the $\ell = 1$ mode can be seeded by target
offset from the center of the target chamber, laser-beam mispointing, and power imbalance, as well as $\ell = 1$ asymmetry in the target. Within the statistical mapping framework, the effect of random nonuniformities must be accounted for post-shot through parameters related either to the amplitude of the random seeds or to any measured signatures of the random mode. The experimental signature associated with an $\ell = 1$ asymmetry is the ratio $R_T = \frac{T_{\text{max}}}{T_{\text{min}}}$ between the maximum and minimum ion temperatures $T_i$ when measured over multiple lines of sight (LOS). The ion temperature is experimentally inferred from the broadening of the D+T fusion peak in the neutron spectrum measured by neutron time-of-flight (NTOF) detectors [45]. In addition to thermal motion, the neutron peak is broadened by fluid flows in the hot spot [46, 41, 42], where isotropic flows generated by asymmetric compression increase the apparent $T_i$ in all LOS, while anisotropic flows generated by certain low modes can lead to significantly higher apparent $T_i$ in certain LOS than others. In particular, the $\ell = 1$ mode produces a large-scale jet in the hot spot with an apparent maximum $T_i$ along the axis of the jet and a minimum perpendicular to it. Three-dimensional simulations of single-mode perturbations in the deceleration phase of ICF implosions in Ref. [47] show that the asymmetry in apparent $T_i$ of a mode $\ell = 1$ is significantly larger relative to other modes and the associated yield-over-clean can be approximated as a power law of the $T_i$ asymmetry $\text{YOC}_{\ell=1}^{\text{sim}} \sim R_T^{-1.53}$. The presence of low modes other than $\ell = 1$ in implosion experiments complicates the relationship between the yield degradation and $R_T$ at small values of $R_T$ because the sum of the contributions of the various modes is not a unique power law of $R_T$. Furthermore, the $T_i$ measure-
ment error of \(\approx 10\%\), as well as the imperfect coverage for measuring \(T_{\text{max}}\) and \(T_{\text{min}}\) of the six NTOF LOS available on OMEGA [48], necessitates the introduction of a threshold parameter \(R_T^{\text{min}}\) below which the yield degradation cannot be detected by the mapping model. The yield-over-clean caused by the \(\ell = 1\) mode therefore becomes

\[
\text{YOC}_{\ell=1} \sim \hat{R}_T^{\mu_5}, \quad \hat{R}_T \equiv \max \left(1, \frac{R_T}{R_T^{\text{min}}} \right),
\]

(3.3)

where \(\mu_5\) is obtained through the regression over the entire database as described in Sec. 2.2, while \(R_T^{\text{min}}\) is found by minimizing the cross-validation error of the model.

Figure 3.2: Experimentally inferred yield degradation due to ion temperature asymmetry extracted from the data using Eq. (3.1). The dashed curve corresponds to the mapping model relation in Eq. (2.9).

Figure 3.2 shows the inferred yield degradation caused by the \(\ell = 1\) mode extracted from the experimental database using the formula in
Eq. (3.1) compared to the predicted degradation $\text{YOC}_{\ell=1} \sim \hat{R}_{T}^{-1.37}$. The experimental data are in good agreement with the power-law approximation at increasing $\hat{R}_{T}$, while more scatter is seen as $\hat{R}_{T} \to 1$ for the reasons outlined above. The exponent $-1.37$ of $\hat{R}_{T}$, inferred by the mapping model, is similar to the simulated value of $-1.53$ for an $\ell = 1$ asymmetry [47]. The best-fit value of the threshold parameter $R_{T,\text{min}} = 1.14$ is consistent with the $\approx 10\%$ measurement uncertainty of $T_i$. Since no other signatures of random asymmetries are observed on OMEGA and the random variations in yield are well explained by $\text{YOC}_{\ell=1}$, all nonuniformity sources other than the $\ell = 1$ seeds listed above can be considered systematic with approximately constant seeds.

### 3.3 Finite beam size

The $\text{YOC}_b$ term represents the yield degradation caused by finite beam size, which is determined by the characteristics of the distributed phase plates [24] being used. The OMEGA laser illuminates the target with 60 partially overlapping beams arranged in a “soccer ball” geometry of pentagons and hexagons with the beams located at the vertices of the polygons. This beam geometry facilitates a uniform symmetric illumination suitable for high performance direct-drive ICF experiments; however, because of the finite number and size of the beams, some residual nonuniformity remains. The amplitude of this mid-mode nonuniformity seed can be parametrized by the ratio of beam radius to target radius $R_b/R_t$, which measures the degree of overlap of the beams. Figure 3.3 shows the beam mode amplitude, measured by the rms of the illumina-
tion nonuniformity, as a function of $R_b/R_t$. At small $R_b/R_t$, the beam mode rms decreases quickly with increasing beam overlap characterized by increasing $R_b/R_t$ until saturating at $R_b \approx 0.82R_t$, where the nonuniformity on target is already very small and further increasing the beam overlap has no effect.

Figure 3.3: The relative rms nonuniformity of the laser illumination as a function of beam-to-target size ratio according to hard-sphere projection.

Further degradation of fusion yield with reduced beam overlap can be related to changes in laser imprint [39] and cross-beam energy transfer (CBET) [17, 18]. Laser imprint is a short wavelength perturbation caused by the “imprinting” of the laser speckle pattern on the capsule surface at early time. Once a sufficiently extensive plasma layer has formed around the target, the conduction zone between the target ablation surface and laser absorption region acts to smooth out the short-wavelength perturbations in the laser. The imprint nonuniformity is greatly reduced by overlapping beams; therefore, part of the yield degradation at small $R_b/R_t$ could be explained by increased imprint seeds. While CBET is modeled in LILAC, as discussed in Sec. 2.2,
any inaccuracies in the 1-D physics models will manifest in the mapping model as additional degradations in yield, which must be approximated with functions of 1-D parameters. Since the amount of CBET strongly depends on the degree of beam overlap, the observed YOC\textsubscript{b} could be related to less than predicted CBET mitigation at smaller beam sizes.

![Figure 3.4](image)

Figure 3.4: Experimentally inferred fusion yield degradation caused by finite beam size extracted from the data using Eq. (3.1). The dotted orange curve indicates the power-law relation from a mapping model with only a simple $R_b/R_t$ dependence. The dashed black curve shows the more-complex relationship from the mapping model with both terms Eq. (3.4) and Eq. (3.5) included.

As the beam radius is increased beyond the target radius, the beam overlap on the target reaches a maximum where further increasing the beam size has no additional effect on the degree of overlap. Such behavior cannot be captured by a simple power-law dependence; instead
a cutoff is applied at \( R_b = R_t \) such that

\[
\overline{R}_{b/t} = \begin{cases} 
\frac{R_b}{R_t}, & \text{if } R_b < R_t \\
1, & \text{if } R_b \geq R_t
\end{cases}
\] (3.4)

While the existence of asymptotic behavior at large \( R_b/R_t \) is physically justified, the exact threshold is unclear, and \( R_b = R_t \) was chosen here as the threshold to fit the handful of shots in the database with \( R_b > R_t \). Since the data at \( R_b > R_t \) are scarce, the threshold function behaves as a switch that only acts on a small number of shots, increasing the likelihood of overfitting the model in this region; therefore, the confidence in prediction accuracy is low. This is not detrimental to predicting the fusion yield of high-performing implosion designs, which use a beam radius smaller than the target radius to couple more of the laser energy to the target.

Figure 3.4 shows the experimentally inferred yield degradation caused by finite beam size, extracted from the database using the formula in Eq. (3.1), compared to the predicted degradation in Eq. (2.9). The dotted orange line in Fig. 3.4 shows the inferred dependence on the simple beam-to-target ratio defined in Eq. (3.4), where the degradation scales like \( \text{YOC}_b \sim \left( \overline{R}_{b/t} \right)^{3.01} \). The dashed black line shows the dependence of fusion yield on a more-complex parameter that takes into account the physics of the various degradation sources as outlined below.

Three dimensional simulations with the radiation-hydrodynamics code \textit{ASTER} [49] indicate a power-law regime at sufficiently large seeds of the beam mode in highly convergent targets where the yield-over-clean
can be approximated as $\text{YOC} \sim C_R^{-0.7} (R_b/R_t)^{1.9}$ (Ref. [50]), where the convergence ratio $C_R$ is the ratio of the stagnated shell radius to the initial target radius. Since the experimentally inferred dependence on $R_b/R_t$ (as shown in Fig. 3.4) is significantly stronger and starts affecting the fusion yield at higher values of $R_b/R_t$, the observed degradation cannot be explained by the beam mode alone.

To uncover the physical mechanisms responsible for the observed dependence on $R_b/R_t$, the parameter in the mapping model must be re-defined in such a way as to be able to discern between the effects of the different possible degradation mechanisms. Motivated by the saturation of the beam-mode seed amplitude at $R_b/R_t \approx 0.82$, as measured by the rms of the hard-sphere projection in Fig. 1.6(b), a threshold parameter $R_{b/t}^{\text{crit}}$ is defined, which divides the parameter space into two distinct regions: (1) a region at $R_b/R_t < R_{b/t}^{\text{crit}}$ where the fusion yield is degraded by the beam mode, and (2) a region at $R_b/R_t \geq R_{b/t}^{\text{crit}}$ where the beam-mode effect is deemed small. This can be described in the mapping model by the introduction of the beam mode parameter,

$$
\text{YOC}_b \sim (\hat{R}_{b/t})^{\mu_a} \left(\hat{R}_{b/t}\right)^{\mu_7}, \quad \text{where}
$$

$$
\hat{R}_{b/t} = \begin{cases} 
\frac{R_b}{R_t R_{b/t}^{\text{crit}}}, & \text{if } R_b/R_t < R_{b/t}^{\text{crit}} \\
1, & \text{if } R_b/R_t \geq R_{b/t}^{\text{crit}}
\end{cases}
$$

(3.5)

where the value of the critical parameter $R_{b/t}^{\text{crit}}$ is found by minimizing the cross-validation error. The predicted dependency on the YOC$_b$ term in Eq. (3.5) is shown in Fig. 3.4 with the dashed black line. The $R_b/R_t$
degradation is independent of all the other terms in the model and, as a result, the two parametrizations under comparison here do not meaningfully affect the exponents of the rest of the terms in Eq. (2.9). Therefore, by freezing all the other exponents, a direct comparison can be made between the two formulations for the YOC\(_b\) term, which can be seen in Fig. 3.4. As expected, the major difference between the two models occurs below the \(R_b/R_t = R_{b/t}^{\text{crit}}\) threshold where the more-complex model predicts a stronger degradation with decreasing \(R_b/R_t\). Still, the prediction confidence is fairly low because of the scarcity of data and with a maximum difference of only \(\sim 30\%\) in the predicted yield of the two models in the most extreme case, more data are needed to determine which of the models is more accurate. Increasing the predictive capability of the statistical mapping model in this region is of high importance to the performance campaign on OMEGA since the highest performing implosion designs operate near the \(R_{b/t}^{\text{crit}}\) threshold. Furthermore, understanding the exact physical mechanisms behind the observed \(R_b/R_t\) degradation is crucial to future mitigation methods. For example, techniques such as beam zooming [51] can be used to reduce the beam size at late time in order to increase absorbed laser energy. For the part of the observed \(R_b/R_t\) dependence that is caused by changes in the laser-imprint seeds, a late-time reduction in \(R_b/R_t\) would increase energy coupling without any additional penalty caused by imprint since the perturbations are seeded early in the laser pulse when there is no conduction zone to smooth out the nonuniformities. The same is true to a lesser degree in the case of beam-mode degradation [50]. If, however, there are inaccuracies in modeling CBET mitigation at smaller
$R_b/R_t$, then zooming may not reduce the degradation with respect to simulations.

### 3.4 Hydrodynamic effects

The most powerful aspect of the statistical mapping approach for ICF applications is the capability to predict the effect of systematic 3-D nonuniformities by using only 1-D parameters. Furthermore, the 1-D code does not have to be perfectly accurate because any discrepancies can also be captured by the mapping function. Of course, as pointed out in Sec. 2.1, the further away the 1-D codes are from experimental reality, the more complex the mapping function becomes and may prove practically impossible to find given the scarcity of the data. The effect of systematic nonuniformities in the OMEGA database as well as the inaccuracies in the 1-D code LILAC are reflected in the YOC$_h$ term.

The parameters chosen to model the degradation caused by systematic nonuniformities are guided by physical considerations. As explained in Sec. 1.3.3, the critical value of IFAR where the bubble front of short wavelength perturbations completely penetrates the shell is a function of the shell adiabat $\alpha_F$; therefore, one can construct a single parameter $I_\alpha \equiv \left(\frac{\alpha_F}{3}\right)_{1.1}^{1.1}$ that describes the effect of short-wavelength perturbations [25, 26]. At a sufficiently high adiabat and low IFAR, the extent of the bubble front penetration into the shell becomes small and further increasing the stability of the implosion is not expected to significantly improve the yield-over-clean; therefore, a piecewise constant exponent
is used in the power-law approximation of \( \text{YOC}_h \):

\[
\mu_1 = \begin{cases} 
\mu_{1, <}, & \text{if } I_\alpha < I_{\text{crit}} \\
\mu_{1, >}, & \text{if } I_{\text{crit}} < I_\alpha < I_{\text{cut}} \\
0, & \text{if } I_\alpha > I_{\text{cut}}
\end{cases}
\]  

(3.6)

where the parameters \( I_{\text{crit}} \) and \( I_{\text{cut}} \) are determined by minimizing the cross-validation error of the model. A power-law exponent of this form was chosen to discriminate between three distinct stability regimes based on physical considerations: (1) a highly unstable regime at \( I_\alpha < I_{\text{crit}} \), where the fusion yield is strongly degraded by the bubble front penetrating through the shell inner surface; (2) a moderately stable regime at \( I_{\text{crit}} < I_\alpha < I_{\text{cut}} \), where the bubble front undergoes significant growth but an inner layer of the shell remains intact; and (3) a stable regime at \( I_\alpha > I_{\text{cut}} \), where the bubble front does not reach far enough into the shell to significantly affect the fusion yield of the implosion.

The convergence ratio \( C_R \equiv R_0 / R_{\text{stag}} \), where \( R_{\text{stag}} \) is the radius of the shell at stagnation, is used to model the yield degradation caused by low and mid modes that can undergo significant Bell-Plesset growth in the deceleration phase. Finally, a weak dependence is found on the dimensionless parameter \( \hat{D} \equiv R_{\text{out}} / R_{\text{in}} \), the ratio of the initial shell outer radius to the inner radius, which is proportional to the shell thickness. This dependence is hypothesized to be related to inaccuracies in modeling shock transit through the shell. Combining all the parameters together,
the hydrodynamic degradation term becomes

\[ \text{YOC}_h \sim I^\mu_1 C^\mu_2 \hat{D}^\mu_3. \] (3.7)

Figure 3.5: Experimentally inferred yield degradation caused by convergence effects extracted from the data using Eq. (3.1) as a function of the convergence ratio $C_R$. The dashed curve corresponds to the mapping model relation in Eq. (2.9).

Figure 3.5 shows the inferred yield degradation as a function of only the convergence ratio $C_R$ compared with the prediction from the mapping model Eq. (2.9). Here, the YOC is normalized to the maximum YOC available in the OMEGA database and therefore the value quoted here serves as an upper bound to YOC$_{C_R}$. Since the mapping model cannot infer absolute degradation, it cannot be ruled out that even the shots with $C_R < 15$ suffer from some degradation related to convergence effects. The figure shows that the fusion yield of the most convergent shots is
degraded by \(\approx 40\%\), well-predicted by the mapping model. The highest performing shot 96806 exhibits a \(\text{YOC}_{C_R} \approx 79\%\) with \(C_R = 19.8\). Note that the exponent of the convergence ratio term \((-0.78, \text{see Table 2.1})\) is in good agreement with the dependence \(\text{YOC}^{\text{sim}} \sim C_R^{-0.7}\) seen in 3-D simulations with mid-mode perturbations seeded by OMEGA port-geometry illumination nonuniformities [50].

![Figure 3.6: Experimentally inferred yield degradation caused by short wavelength hydrodynamic instabilities extracted from the data using Eq. (3.1). The dashed curve corresponds to the mapping model relation in Eq. (2.9). The dotted lines correspond to \(I_{\text{crit}}\) and \(I_{\text{cut}}\) where the switches in the mapping model power-law exponent occur.](image)

*Figure 3.6* shows the experimentally inferred yield degradation from Eq. (3.1) caused by only the growth of short-wavelength modes as a function of the stability parameter \(I_\alpha\). The piecewise power-law prediction of the mapping model is shown with a dashed line and the dotted lines correspond to the critical value \(I_{\text{crit}} = 0.8\) and the cutoff value
$I_{\text{cut}} = 2.4$, where the switches in the mapping model exponent occur. In previous studies, a similar critical threshold was found for the shell areal density at $I_\alpha = 1$ in Ref. [25] and experimental signatures of fuel-ablator mix were found to be correlated with the stability parameter $I_\alpha$ (Ref. [52]). The highest fusion yields on OMEGA have been achieved by implosions in the moderately stable regime with $I_\alpha \approx 1.5$–2.0. It is clear from Fig. 3.6 that only a very small number of shots determine the behavior of the model in the stable regime and therefore the ability to generalize the model predictions to new designs in this parameter region is low.
4.1 Fill age scan

To study the inferred dependence on fill age discussed in Sec. 3.1, a series of experiments was carried out at five different fill ages ranging from 3 to 34 days while keeping all other parameters of the target and laser pulse shape fixed. The results are given in Table 4.1. Shot 94717 showed a significant ion-temperature asymmetry while the rest of the shots had asymmetries below the threshold \( R_{T}^{\text{min}} \). Therefore the corrected YOC\(_{\text{corr}}^{\text{exp}} = YOC_{\text{exp}}^{\text{corr}} / \tilde{R}_{T}^{-1.37} \) is provided, where the measured yield-over-clean is corrected for the ion-temperature asymmetry using the relation from the mapping model [Eq. (2.9)] to enable a fair comparison between all the shots. The asymmetry correction along with the very small variation in the simulated 1-D performance of all the shots ensures that this set of experiments clearly isolates the effect of the fill age. Furthermore, shots 96806 and 96808 were taken on the same day with different fill ages, removing the effect of any day-to-day variations in the system and thereby providing the cleanest possible A/B comparison of the fill-age effect on the fusion yield. The mapping model predictions for the yield degradation caused by \(^3\text{He}\) accumulation in the vapor, normalized to a three-day fill are also provided in Table 4.1, showing excellent agreement with the measurements. The good agreement is further highlighted by the inferred YOC\(_{\text{He}}^{\text{exp}}\) in Fig. 4.1, where this set
<table>
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<th>Fusion Yield ($\times 10^{14}$)</th>
<th>Fill age (days)</th>
<th>$\hat{R}_T$</th>
<th>$\text{YOC}^{\text{exp}}$</th>
<th>$\text{YOC}^{\text{corr}}$</th>
<th>Predicted $\text{YOC}_{\text{He}}$</th>
<th>Measured $\text{YOC}_{\text{He}}$</th>
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<td>90288</td>
<td>1.52</td>
<td>5.0</td>
<td>1.0</td>
<td>0.38</td>
<td>0.38</td>
<td>0.93</td>
<td>0.94</td>
</tr>
<tr>
<td>96808</td>
<td>1.33</td>
<td>8.1</td>
<td>1.0</td>
<td>0.37</td>
<td>0.37</td>
<td>0.85</td>
<td>0.92</td>
</tr>
<tr>
<td>94350</td>
<td>0.83</td>
<td>14.3</td>
<td>1.0</td>
<td>0.24</td>
<td>0.24</td>
<td>0.69</td>
<td>0.60</td>
</tr>
<tr>
<td>94717</td>
<td>0.45</td>
<td>45.1</td>
<td>1.36</td>
<td>0.11</td>
<td>0.17</td>
<td>0.38</td>
<td>0.41</td>
</tr>
</tbody>
</table>

Table 4.1: Results from a controlled set of experiments where all the input parameters except for the fill age were held constant to within normal shot-to-shot variations. $\text{YOC}^{\text{corr}}$ represents the yield-over-clean corrected for the effect of the random $\ell = 1$ asymmetry. The predicted $\text{YOC}_{\text{He}}$ is the YOC predicted by the mapping model normalized to an implosion with a three-day fill; the measured $\text{YOC}_{\text{He}}$ is the $\text{YOC}^{\text{corr}}$ normalized to shot 96806.

of shots is shown with the yellow diamonds.
Figure 4.1: Experimentally inferred degradation caused by accumulation of $^3\text{He}$ in the vapor, extracted from the data using Eq. (3.1). The yellow diamonds show results from repeated experiments of the same implosion design shot on targets with varying fill ages. The dashed curve corresponds to the mapping model relation in Eq. (2.9).

### 4.2 Large diameter targets

A major application of the statistical mapping method is to guide the design of future experiments by accurately modeling the response of the fusion yield to small changes in input parameters. Here, a path to higher yields through the use of larger diameter targets is shown. The advantage of increasing the target size is the mitigation of CBET caused by reduced beam overlap, which increases the efficiency of coupling laser energy to the target. This leads to a significantly higher implosion velocity, which causes a large increase in the 1-D yield. On the other hand, increasing the target size will reduce the beam-to-target ratio $R_b/R_t$, which strongly degrades the yield-over-clean according to the mapping model [Eq. (2.9)]. In fact, the best performing OMEGA implosions operate close to the threshold of $R_b/R_t = 0.86$, therefore in larger targets the YOC degrades very rapidly as $\text{YOC} \sim (R_b/R_t)^{1.59}$. 
The thickness of the DT ice layer is a further knob that can be used to tune the implosion in this potential path toward higher performance. For a fixed target mass, larger targets have thinner ice and therefore a higher IFAR, which degrades the YOC according to the mapping model [Eq. (2.9)]. Thanks to the increased energy coupling, however, the ice thickness for larger targets can be increased while still maintaining a higher implosion velocity compared to smaller targets. The mapping model suggests, therefore, that the optimum target size can be found by balancing the competing effects of energy coupling, $R_b/R_t$ degradation, and hydrodynamic stability (IFAR).

Table 4.2 shows data from the best performing OMEGA experiment with a 960-$\mu$m-outer-diameter (OD) target (shot 96806) as well as data from implosions with 1010-$\mu$m-OD targets, which were predicted to achieve higher performance by the mapping model. All the shots in the table have ion-temperature asymmetry below the $R_t^{\text{min}}$ threshold; however, because of varying fill ages, the fusion yield must be corrected for a fair comparison of the implosion designs. In Table 4.2, the “corrected yield” refers to the fusion yield normalized to a three-day fill using the dependence on fill age in Eq. (2.9). The table shows that the 1010-$\mu$m-OD shot 100535, with a significantly more massive target, achieved a higher neutron yield ($1.77 \times 10^{14}$) than the best performing 960-$\mu$m-OD target with a corrected yield of $1.56 \times 10^{14}$, whereas the older shot 91312 with a 1010 $\mu$m-OD target and 42.4-$\mu$m-ice filled 14 days prior to the shot would have achieved a 31% yield improvement at $2.05 \times 10^{14}$ over the 960-$\mu$m-OD target, had it been shot with a short fill of three days.
Table 4.2: Data from experiments with larger target outer diameters. The corrected yield refers to the fusion yield normalized to a three-day DT fill age using the relation from the statistical mapping model [Eq. (2.9)].

<table>
<thead>
<tr>
<th>Shot number</th>
<th>96806</th>
<th>100535</th>
<th>91312</th>
</tr>
</thead>
<tbody>
<tr>
<td>Neutron yield</td>
<td>$1.56 \times 10^{14}$</td>
<td>$1.72 \times 10^{14}$</td>
<td>$1.41 \times 10^{14}$</td>
</tr>
<tr>
<td>Outer diameter ($\mu$m)</td>
<td>958</td>
<td>1018</td>
<td>1007</td>
</tr>
<tr>
<td>Ice thickness ($\mu$m)</td>
<td>40.6</td>
<td>47.4</td>
<td>42.4</td>
</tr>
<tr>
<td>Fill age (days)</td>
<td>3.1</td>
<td>3.3</td>
<td>14.3</td>
</tr>
<tr>
<td>Corrected yield</td>
<td>$1.56 \times 10^{14}$</td>
<td>$1.77 \times 10^{14}$</td>
<td>$2.05 \times 10^{14}$</td>
</tr>
</tbody>
</table>
CHAPTER 5
STATISTICAL MAPPING OF AREAL DENSITY

In this chapter the same statistical mapping framework that was developed and applied to the fusion yield on OMEGA is used to construct a statistical mapping model of the areal density $\rho_R$.

5.1 The difficulty of predicting areal density

Developing a predictive model of the areal density is significantly more challenging than predicting the fusion yield for several reasons. The relative range of the measured areal densities in the OMEGA database is small, varying by a factor of 4, compared to the factor of 30 variation in the measured fusion yields. The areal density in the experiments is not strongly dependent on a single dominant parameter analogous to the implosion velocity in the case of fusion yield but appears to be instead determined by weak contributions from a number of different parameters. This low responsiveness of the $\rho_R$ to changes in the implosion design makes it very difficult to identify the key parameters that determine the outcome of experiments. As a result, the prediction accuracy of the model that is developed in this chapter can be reproduced by several other formulations of the statistical mapping model with different choices of parameter sets. This calls into question the interpretation of the power-law terms as effects of specific physical degradation mechanisms as was done for the fusion yield model in Ch. 3 because multiple
potentially conflicting explanations can not be ruled out based on the current dataset.

Another issue with predicting the \( \rho R \) is the inaccuracy of the \( \rho R \) measurement, which relies on an accurate understanding of complicated physics. The areal density in OMEGA cryogenic implosions is calculated by taking the ratio of D+T neutrons scattered by the areal density to the number of primary D+T neutrons. This ratio is inferred from measurements of the neutron energy spectrum. The primary D+T neutrons will have an energy of 14.03 MeV (with some variance caused by thermal and fluid motion in the hotspot), while the scattered neutrons will have a lower energy, depending on the scattering angle. The neutron energy spectrum is measured on OMEGA using either the magnetic recoil spectrometer (MRS) \cite{53} or neutron time-of-flight (NTOF) \cite{54} detectors.

The MRS diagnostic is configured to detect forward-scattered neutrons at a relatively small scattering angle, thereby sampling a region of the energy spectrum that contains very little contribution from sources other than the down-scattered D+T neutrons. Still, the measurement uncertainty of the MRS diagnostic is typically around 10 %, which already limits the ability to increase \( \rho R \) by implementing the same type of incremental guided design improvements that have proved so successful in increasing the fusion yield on OMEGA as highlighted in Ch. 4 and Ref. \cite{11}. The \( \rho R \) measurement is further complicated by the fact that in the presence of strong low modes the amplitude of modulation of the shell areal density at the time of peak convergence can reach a large fraction of the \( \rho R \) that would be measured in a completely symmetric implosion. This means that multiple \( \rho R \) detectors that sample neutrons
scattered by different regions of the shell can infer substantially different values for the areal density. Note that this is not the case for the fusion yield in the presence of low modes because the spatially varying attenuation of the D+T primary neutrons in the case of the low areal densities on OMEGA is negligibly small. For this reason, the measured areal density used in the $\rho R$ mapping model is defined as the geometric mean

$$\rho_{\text{measured}} = \sqrt{\rho_{\text{MRS}} \times \rho_{\text{NTOF}}}, \quad (5.1)$$

where $\rho_{\text{MRS}}$ is the areal density measured with the MRS detector and $\rho_{\text{NTOF}}$ is measured by the NTOF detector along a different line of sight. While such averaging clearly improves the accuracy of the $\rho R$ measurement, the two lines of sight are not enough to eliminate the significant error with respect to a $4\pi$ averaged $\rho R$ in a distorted implosion [55].

Furthermore, the uncertainty of the NTOF measurement is around twice that of the MRS at $\approx 20\%$ because the NTOF instrument detects back-scattered neutrons in a much lower energy range. The signal in this region of the neutron spectrum includes significant background contributions from both the D+D and T+T primary neutrons as well as several neutron scattering events and even environmental scattering of neutrons off objects in the target chamber. Extracting the areal density from the neutron spectrum therefore requires an accurate model of the background signal, which increases the uncertainty of the $\rho R$ measurement.
5.2 Mapping the areal density of OMEGA implosions

Using the statistical mapping method outlined in Ch. 2, a power-law model was fit to the $\rho R$-over-clean ($\rho$ROC = $\rho R_{\text{measured}}/\rho R_{\text{LILAC}}$) measurements in the OMEGA database. The results of the fit are shown in Fig. 5.1, where the relatively large uncertainty of the $\rho R$ measurements discussed above is clearly visible in the magnitude of the error bars.

The parameters chosen for the $\rho$ROC model were all previously defined in Ch. 2 with the exception of $\rho$ROC$_{\text{He}}^{\text{sim}}$, which is defined as the simulated $\rho R$ degradation due to accumulation of $^3$He in the vapor re-
Table 5.1: Exponents and confidence intervals for all mapping terms resulting from fitting the model in Eq. (5.2) to the areal density measurements in the OMEGA database.

<table>
<thead>
<tr>
<th>Parameter</th>
<th>Coefficient</th>
<th>95% confidence interval</th>
</tr>
</thead>
<tbody>
<tr>
<td>$I_{\alpha}^\mu_1$</td>
<td>$\mu_{1,&lt;} = 0.28$</td>
<td>$\mu_{1,&lt;} = 0.21 \ldots 0.34$</td>
</tr>
<tr>
<td></td>
<td>$\mu_{1,&gt;} = 0.18$</td>
<td>$\mu_{1,&gt;} = 0.13 \ldots 0.23$</td>
</tr>
<tr>
<td></td>
<td>$I_{\text{crit}} = 0.8$</td>
<td></td>
</tr>
<tr>
<td>$C_R^{\mu_2}$</td>
<td>$\mu_2 = -0.98$</td>
<td>$\mu_2 = -1.10 \ldots -0.84$</td>
</tr>
<tr>
<td>$(\bar{R}_{b/t})^{\mu_3}$</td>
<td>$\mu_3 = 0.27$</td>
<td>$\mu_3 = 0.04 \ldots 0.50$</td>
</tr>
<tr>
<td>$\hat{R}^{\mu_4}_T$</td>
<td>$\mu_4 = -0.35$</td>
<td>$\mu_4 = -0.50 \ldots 0.18$</td>
</tr>
<tr>
<td>$R_{\text{min}}^{\mu_5}$</td>
<td>$\mu_5 = 1.14$</td>
<td></td>
</tr>
<tr>
<td>$(\rho_{\text{ROC}}^{\text{sim}}_{\text{He}})^{\mu_5}$</td>
<td>$\mu_5 = 0.23$</td>
<td>$\mu_5 = -0.82 \ldots 0.41$</td>
</tr>
<tr>
<td>$\theta^{\mu_6}_T$</td>
<td>$\mu_6 = 0.21$</td>
<td>$\mu_6 = 0.01 \ldots 0.40$</td>
</tr>
</tbody>
</table>

The data indicate a correlation of $\rho R$ with the stability parameter $I_{\alpha}$. Similarly to the yield model in Ch. 2, the model was allowed to fit the data in two distinct regions above and below the critical value of the stability parameter $I_{\text{crit}} = 0.8$. Asymptotic behavior at large values of the stability parameter was not considered here. The dependence of areal density on the stability parameter is consistent with earlier empirical results in Ref. [25].

A dependency of $\rho R$ on the beam to target size ratio was also found,
similarly to the yield model. The inferred $\rho R$ dependence is considerably weaker as reflected in the 95 % confidence interval that nearly includes a zero exponent. The more sophisticated approach employed in the yield model with several threshold parameters was not attempted here because of significantly more noise present in the $\rho R$ database.

The areal density was found to be inversely correlated with the ion temperature asymmetry $\tilde{R}_T$. Even though a $\rho R$ measurement along a single line of sight could either increase or decrease in the presence of strong low mode asymmetries, the average $\rho R$ measured along the two lines-of-sight available on OMEGA was found to be able to capture, on average, the degradation of the average shell areal density.

No dependence was found of $\rho R$ on the simulated degradation caused by $^3$He accumulation in the vapor region with the wide 95% confidence firmly including the zero exponent. A lack of dependence on the fill age is not entirely unexpected given that the simulated degradation $\rho \text{ROC}_{\text{He}}^{\text{sim}}$ is rather weak to begin with. The simulated $\rho R$-over-clean for a target with a 14-day-old fill is $\approx 0.94$ while the simulated YOC for the same target is $\approx 0.69$. A very weak dependence on the tritium concentration of the ice layer was found in the $\rho R$ model in Eq. (5.2), which could indicate some dependence on the details of the target filling and layering procedure, however the wide confidence interval of this exponent shows a large uncertainty.

The above dependencies are all rather weak and, as mentioned, other mapping formulas can be found that reproduce similar prediction accuracy and therefore there is very little confidence in any physics interpretation of the $\rho R$ degradation with respect to simulations as inferred
from the statistical mapping model. The strongest dependency by far in the $\rho R$ model of Eq. 5.2 is the dependency on the target convergence ratio. The dependence on $C_R$ is not as replaceable by other parameters as the weak dependencies listed above, however it is still not clear how one should interpret the physics of the $\rho R$ degradation related to this term. Since the convergence ratio in the simulations is very strongly correlated to the areal density this cannot be simply interpreted as, for example, a $\rho R$ degradation caused by growth of mid modes in highly convergent targets. This is because the strong correlation prevents de-coupling the effects of 1-D physics and instability growth. What is clear is that the $C_R$ term imposes a strong penalty to designs with high convergence i.e. shots with low shell adiabat. This reinforces the well known problem in ICF that shots with low adiabat do not converge as much as expected based on simulations, which is one of the main issues why it is so difficult to reach ignition. Since the $\rho R$ mapping model in Eq. (5.2) is not capable of uncovering the physical processes that are causing lower than expected areal density in experiments, one must turn to dedicated physics studies that aim to investigate these issues in a more controlled fashion. One of the possible explanations for lower than expected convergence at low adiabats is that there could be inaccuracies in the current radiation-hydrodynamics codes in modeling the dynamics of the material released into the hotspot at shock breakout, which influences the convergence of implosions by setting the initial hotspot energy. The experiments described in the next section were designed to produce data that can be used to benchmark the simulation of shock release.
CHAPTER 6
SHOCK RELEASE EXPERIMENTS

6.1 The role of shock release in ICF

Figure 6.1: The mechanism of shock release in an ICF target. (a) the initial condition of a dense shell surrounding a low density vapor region; (b) a strong shock is driven inside the shell with the laser; (c) at shock breakout a weak shock is transmitted into the vapor region and a rarefaction wave is reflected back toward the outer surface of the shell. The material in the tail of the rarefaction wave that converges toward the center makes up the shock release.
As discussed in Sec. 1.2, the initial part of the laser pulse in an ICF implosion is designed to launch a series of shocks into the shell to set the adiabat, which is one of the most important parameters that affects the performance of the implosion. In the following sections, the dynamics of the material that is released from the inner surface of the shell after shock breakout (i.e. the shock release) are studied extensively.

Figure 6.1 shows the density and pressure profiles of a spherical ICF target driven symmetrically by a laser, the initial profile of a dense shell surrounding a low density vapor region is shown in Fig. 6.1(a). While a typical high performance ICF pulse shape is designed to launch two or more shocks into the target, here the case of a single shock launched by a picket [Fig. 6.1(b)] is considered. At the time the shock reaches the inner surface of the shell (commonly referred to as the shock breakout), a weak converging shock is transmitted into the low density vapor region and a rarefaction wave travels back toward the outer surface of the shell [Fig. 6.1(c)]. The shell material that comprises the tail of the rarefaction wave converges toward the center of the target and is called the shock release.

Prior experiments designed to measure the dynamics of the shock release have focused on probing the very low density tail of the rarefaction wave. Haberberger et al. [56] used images of the 4ω probe [57] to infer the trajectories of material in the shock release at number densities between $10^{19}$ and $10^{20}$ cm$^{-3}$. At these densities, it was found that the leading edge of the release was moving significantly faster than predicted by radiation-hydrodynamics codes. A pre-expansion of the rear surface of the target caused by radiation preheat was hypothesized to
be the cause of the observed discrepancy and simulations with an imposed density gradient on the rear surface were shown to reproduce the measured trajectories of the low density material. Separately, in Ref. [26], the separation of the carbon and hydrogen species of the target at shock breakout was suggested as a mechanism to explain the measurements in Ref. [56] based on results from molecular dynamics simulations. The shock release experiments discussed in this chapter are designed to measure the kinetic energy of the bulk of the release material, which is converted into the internal energy of the hotspot upon convergence of the release material. Since the hotspot pressure at the start of deceleration influences the final target convergence, it is the kinetic energy of the release material that is the aspect of shock release that is most relevant to ICF implosion performance.

6.1.1 Effect on stagnation pressure

The shock release is important in ICF experiments because it strongly influences the final convergence of the implosion. As the release material stagnates in the center of the target, the kinetic energy of the material is converted into internal energy thereby setting the initial hotspot pressure $P_0$.

Consider the dynamics of an imploding shell at the start of the deceleration phase, meaning that the shell has reached peak implosion velocity $V_i$, the laser has turned off and the shell begins to decelerate because of the pressure $P$ inside the shell. From Newton’s second law, the mass of the shell $M_{sh}$ times the acceleration is equal to the force of
the hot spot pressure applied on the inner surface of the shell:

\[
M_{sh} \frac{d^2 R}{dt^2} = 4\pi R^2 P, \tag{6.1}
\]

where \( R = R(t) \) is the radius of the imploding shell. Assuming adiabatic compression and the ideal gas EOS, the energy conservation equation becomes:

\[
P V^{5/3} = P_0 V_0^{5/3}, \tag{6.2}
\]

where \( V \) is the volume of the shell, the subscript 0 refers to conditions at the start of the deceleration phase and an adiabatic index \( \gamma = 5/3 \) was assumed. Eventually the shell will stagnate at its minimal radius \( R_{stag} \). The pressure of the hot spot at the time of shell stagnation – the stagnation pressure \( P_{stag} \) – can be found by integrating the system of equations [Eq. (6.1) and Eq. (6.2)] from the start of the deceleration phase to stagnation. After substituting the pressure from Eq. (6.2) into Eq. (6.1) and multiplying both sides for convenience by \( \frac{dR}{dt} \) we find

\[
\frac{dR}{dt} \frac{d^2 R}{dt^2} = \frac{4\pi P_0 R_0^5}{M_{sh}} \left( \frac{1}{R} \right)^3 \frac{dR}{dt}. \tag{6.3}
\]

This can be rewritten as

\[
\frac{d}{dt} \left[ \left( \frac{dR}{dt} \right)^2 \right] = -\frac{4\pi P_0 R_0^5}{M_{sh}} \frac{d}{dt} \left( \frac{1}{R^2} \right). \tag{6.4}
\]

Integrating both sides from the start of deceleration to stagnation, keep-
ing in mind that the shell velocity \( \frac{dR}{dt} = 0 \) at stagnation, we find

\[
V_i^2 = \frac{4\pi P_0 R_0^5}{M_{sh}} \left( \frac{1}{R_{stag}^2} - \frac{1}{R_0^2} \right).
\]  

(6.5)

Rearranging this and using the energy equation once again to rewrite in terms of the stagnation pressure leads to

\[
P_{stag} = P_0 \left( \frac{M_{sh} V_i^2 + 4\pi P_0 R_0^3}{4\pi P_0 R_0^3} \right)^{5/2}
\]

(6.6)

where the two terms in the numerator are the shell kinetic energy and a multiple of the hot spot internal energy at the start of deceleration. Because the hot spot internal energy is much smaller than the shell kinetic energy at the start of deceleration, the internal energy term can be neglected in the numerator resulting in an inversely proportional relationship between the stagnation pressure and the initial hot spot pressure

\[
P_{stag} \propto P_0^{-3/2}.
\]

(6.7)

Therefore, the stagnation pressure is determined by the kinetic energy of the shock release, which, as mentioned, is what sets the initial hot spot pressure.

### 6.2 Experimental setup for measuring the kinetic energy of shock release

A novel platform was designed to facilitate the measurement of the kinetic energy of the bulk of the release material which ultimately would
be converted to the internal energy of the hot spot at the start of deceleration in an ICF implosion. This is in contrast to prior release experiments that have largely focused on the dynamics of the very low density leading edge of the release which carries a negligible fraction of the kinetic energy that influences the implosion convergence.

6.2.1 Target configuration

Experiments were carried out in both planar and spherical geometry on the OMEGA EP [58] and the OMEGA laser, respectively; the principle and direct measurable quantities were identical in each case. Sketches of the targets used in the experiments are shown in Fig. 6.2(a) and Fig. 6.2(b) for the planar and spherical cases respectively.

![Diagram of experimental setup](image_url)

**Figure 6.2:** Schematic of the experimental setup in (a) planar and (b) spherical geometries. In both cases the laser is used to (1) drive a shock in the polystyrene target, releasing mass at shock breakout from the rear surface. The release material traverses the (2) vacuum gap and collides with the (3) quartz witness. The VISAR probe is used to measure the shock velocity in the witness. Large plastic washers are used in the planar target to shield the foils from radiation from the corona.
The planar target consists of two foils separated by a vacuum gap as illustrated in Fig. 6.2(a):

1. A laser-driven polystyrene foil with an aluminum coating. A short laser pulse is used to drive a shock through this foil, which breaks out of the rear surface and launches a rarefaction wave traveling the opposite way. The tail of the rarefaction wave consisting of material originating from near the rear surface forms the shock release.

2. The release material is allowed to freely expand into a vacuum gap between the two foils.

3. Having traversed the vacuum gap, the release material collides with the quartz foil (referred to as the witness), driving a strong shock inside the foil with a shock pressure of several Mbar).

Upon stagnation of the release material against the witness foil, the kinetic energy of the release is converted into internal energy. This internal energy manifests as the pressure of the shock that is driven inside the witness foil by the release collision. The shock pressure can be related via the Hugoniot relations to the velocity of the shock front (see Sec. 6.4). The velocity interferometer system for any reflector (VISAR) is used as indicated in figure 6.2(a) to measure the velocity of the release-driven shock inside the witness. In addition, the planar target features large washers separated by spacers that serve as structural support to ensure a precise gap between the two foils. The washers also shield the foils from radiation preheat originating from the corona – the
hot plume of plasma that forms from the material that is ablated from the laser-driven side of the polystyrene foil.

The basic principle of the experiment remains the same in the spherical setup [Fig. 6.2(b)], where the role of the laser-driven foil is assumed by the outer polystyrene shell (marked by number 1 in the figure) and the solid fused silica hemisphere (marked by number 3) at the center of the target acts as the witness. The target is arranged in a cone-in-shell setup where the gold cone provides a clear line of sight to the witness hemisphere for the VISAR probe beam as well as structural support for the placement of the hemisphere to the center of the target.

6.2.2 Differences to shock release in ICF implosions

The experimental platform described here is used to study the isolated rarefaction wave dynamics in a simplified form. Several key differences compared to the shock release in ICF implosions are present. In a typical ICF implosion the major shock release of concern is that of DT into a low density gas, whereas the experiments at hand are used to study the release of polystyrene into vacuum. Granted, in an ICF implosion there does exist the shock release of plastic into DT ice, however it is not related to setting the initial hot spot pressure. The difference in the material composition of the release material could be relevant to the release behavior, e.g. the contribution of species separation is certainly more important in a case with a large difference in the atomic masses of the species such as carbon and hydrogen in polystyrene compared to the more equal mass distribution of DT. Furthermore, differences in equation-of-state and opacity tables reduce the relevance of measuring
the shock release of polystyrene to ICF implosions.

In ICF implosions there are typically multiple shocks launched into the shell that are designed to merge in the vapor region. Therefore, some of the release material is shocked multiple times, which complicates the in-flight dynamics. ICF implosions are driven by the laser for an extended time after the shock breakout occurs and therefore the shock release is not a “ballistic” rarefaction wave but instead is influenced in-flight by the driving pressure from the laser. Such in-flight dynamics can be further complicated by a spurious shock wave (often referred to as the “$N + 1$ shock” in ICF literature) that can form in the shell as a consequence of the interaction of the outgoing rarefaction wave with an acoustic wave originating from the ablation surface. The pulse shapes used in this study were carefully designed to avoid the complications arising from driving the target post shock breakout in order to isolate the fundamental dynamics of a pure rarefaction wave.

6.2.3 Laser setup and target variations

The pulse shapes for the release experiments were designed in such a way that the shock generated in the polystyrene foil (outer shell in spherical geometry) would impart enough kinetic energy to the release to drive a measurable shock in the witness. At the same time, as alluded to in the previous chapter, the pulse shapes were designed to cut off prior to the shock breakout from the polystyrene foil (outer shell) to generate a simple rarefaction wave that is not influenced by the laser after shock breakout.

In planar experiments the design objectives were achieved by the
use of a set of square pulse shapes with the duration and intensity of the pulses given in table 6.1. To preserve the planar approximation of the release, SG8-1100 distributed phase plates (DPP) [24] were used to produce a laser beam profile with a spot size of 483 $\mu$m (defined by the radius of the $1/e$ intensity contour) and a super-Gaussian exponent of 8.57. The chosen DPPs were the largest available on OMEGA EP that were capable of delivering the laser intensity required to drive shocks with the necessary strength. In all experiments the targets were driven with two diagonally opposed beams (i.e. either OMEGA EP beams nr. 1 and 2 or beams nr. 3 and 4) at an incidence angle of 23.81°.

The planar experiments were performed on a total of four different target types to study the release under a variety of conditions. Both the thickness of the polystyrene foil and the vacuum gap size were varied and laser pulse shapes with high and low intensities were used to provide a variety of plasma conditions at shock breakout and to control the level of radiation preheat. The specifications of the different target types used in the planar experiments are shown in Table 6.1 along with the corresponding laser pulse shape parameters for each target type.

The thicker 125 um targets were driven with longer pulses to ensure an adequate shock strength at breakout while the shorter pulses were used on the thinner 50 um foils. Low intensity experiments were performed on only two of the target types. Two repeats were performed of the high intensity experiments on all four target types showing excellent repeatability of the measurements (see Sec. 6.5.3) while only one shot on each target type was taken at low intensity.

The spherical targets were driven by picket pulses (see parameters
### Table 6.1: Target and pulse shape configurations used in planar release experiments. PS stands for polystyrene. The laser intensities and pulse durations are given for both high intensity and (low intensity) cases.

<table>
<thead>
<tr>
<th>Target type</th>
<th>PS foil thickness ($\mu$m)</th>
<th>Gap width ($\mu$m)</th>
<th>Laser intensity ($\times 10^{14}$ W/cm$^2$)</th>
<th>Pulse duration (ps)</th>
</tr>
</thead>
<tbody>
<tr>
<td>Type IA</td>
<td>50</td>
<td>100</td>
<td>3.0 (1.2)</td>
<td>600 (600)</td>
</tr>
<tr>
<td>Type IB</td>
<td>50</td>
<td>170</td>
<td>3.0 (1.2)</td>
<td>600 (600)</td>
</tr>
<tr>
<td>Type IIA</td>
<td>125</td>
<td>100</td>
<td>3.6 (1.7)</td>
<td>1300 (2000)</td>
</tr>
<tr>
<td>Type IIB</td>
<td>125</td>
<td>170</td>
<td>3.6 (1.7)</td>
<td>1300 (2000)</td>
</tr>
</tbody>
</table>

### Table 6.2: Target and pulse shape configurations used in spherical release experiments.

<table>
<thead>
<tr>
<th>Description</th>
<th>Shell thickness ($\mu$m)</th>
<th>Radiation shielding</th>
<th>Picket energy (J/beam)</th>
</tr>
</thead>
<tbody>
<tr>
<td>Low intensity</td>
<td>30</td>
<td>-</td>
<td>22</td>
</tr>
<tr>
<td>High intensity</td>
<td>30</td>
<td>-</td>
<td>30</td>
</tr>
<tr>
<td>Preheat target</td>
<td>18.6</td>
<td>0.9 $\mu$m Au</td>
<td>30</td>
</tr>
</tbody>
</table>


in Table 6.2) with the same SG5-850 DPPs that are used in high performance ICF experiments. As is typical to a cone-in-shell configuration, 39 on-target beams were used instead of the full 60 beams available on OMEGA with the beams pointed symmetrically about the cone axis providing maximum illumination uniformity in the region of the target surface opposite the VISAR probe.

Both low and high intensity experiments were performed on the spherical release target with 30 \( \mu \)m shell thickness as seen in Table 6.2. In addition, a separate preheat mitigation experiment with a buried gold layer target driven by a picket pulse was performed using the same spherical cone-in-shell setup (see details in Sec. 6.5.5).

6.2.4 Dedicated shock breakout targets

![Figure 6.3](image)

Figure 6.3: Schematic of the dedicated shock breakout targets in (a) planar and (b) spherical geometries. The VISAR probe is used to measure the shock velocity in the polystyrene target as indicated.

As a consequence of radiation preheat of the witness (see Sec. 6.5.2
and Sec. 6.5.5) dedicated experiments on specialized targets had to be performed in order to measure the time of shock breakout from the laser-driven polystyrene foil (polystyrene shell in spherical geometry). Fig. 6.3 shows the target configuration for the shock breakout experiments to measure the time of shock breakout from the laser-driven foil. The planar shock breakout targets [Fig. 6.3(a)] consist of a standalone polystyrene foil of the same thicknesses used in the release experiments. The shock breakout experiments were performed with all combinations of target types and laser pulse shapes listed in table 6.1. Similarly, in spherical geometry, cone-in-shell targets [Fig. 6.3(b)] in all the shell configurations with the same picket energies as the release experiments in table 6.2 were used to measure the shock transit in the shell and the time of break-out. In all cases, the VISAR diagnostic was used to measure the time of shock breakout and where possible the velocity of the laser-driven shock.

6.3 The Velocity Interferometer System for Any Reflector (VISAR)

The VISAR diagnostic [59] at the OMEGA facility is an active shock detection system, which can infers the velocity of a shock front traveling through a transparent medium as a function of time by using the changes in the signal of a reflected probe beam. In contrast, passive systems have typically been used to detect the flash of light produced by the shock breaking out of an opaque medium. Stepped targets can then be used to infer the shock velocity by considering the different shock breakout times from different regions of the target. It is impor-
tant to note that in case of opaque targets, VISAR can still be used to passively detect the flash from shock breakout. This property is especially useful in many of the experiments described in this chapter where the ionizing radiation from the laser-heated corona causes the targets to become opaque to the VISAR probe beam.

Figure 6.4: A sample VISAR streak camera image from one of the release experiments (OMEGA EP shot 35495). Both upward (accelerating shock) and downward (decelerating shock) shifts in the fringe pattern can be seen along with a discontinuous jump when the signal that is reflected off of the shock front first appears.

The VISAR instrument works by shining a 532 nm probe beam on the target, which reflects back from the first reflecting surface in the tar-
get. In the case of a strong shock traveling in a material that is transparent to the probe beam, this reflecting surface is the ionized shock front. The reflected signal is passed through an interferometer where the beam is split between two interferometer arms (detailed schematics of the VISAR diagnostic at OMEGA are given in Ref. [59]). In the interferometer, an etalon with a known thickness and refractive index is used in one of the arms to apply a delay to the signal. A sinusoidal fringe pattern is then imposed on the output signal. Doppler shifts of the light reflecting from the target – caused by the movement of the reflecting shock front in the experiments discussed here – result in a change in the phase of the fringe pattern. The temporally resolved fringe pattern is recorded with a streak camera; an example of one such image from OMEGA EP shot 35495 is shown in Fig. 6.4. In the figure, an upward shift in the fringes corresponds to an acceleration of the shock front while a downward shift corresponds to a deceleration. By extracting the phase of the fringe pattern in the streak camera image, the velocity of the moving reflector can be inferred, since the velocity sensitivity, typically given as a velocity per fringe (VPF), can be calculated based on the properties of the etalon used in the interferometer. The most common technique to extract the phase from the image is the Fourier Transform Method (FTM) introduced in Ref. [60]. The AnalyzeVISAR code from the Lawrence Livermore National Laboratory, which implements the FTM, was used to analyze the streak camera images in this work.

An ambiguity in the absolute velocity magnitude remains in cases where there are very fast changes in the velocity of the reflecting sur-
face that are not temporally resolved in the streak camera images and appear as discontinuous jumps. For example, such jumps can occur when a shock is transmitted through an interface between two materials with different properties, or when a trailing shock overtakes a leading shock, or as is the case in the experiments discussed here, in the initial moments when the shock front emerges from an opaque material and the reflected signal first appears on the image (see an example in Fig. 6.4). In the presence of these discontinuities, the phase of the fringe pattern is ambiguous by multiples of $2\pi$ and the extent of the velocity jump is not known. To resolve the ambiguity, a second channel is used in the VISAR diagnostic where the reflected signal passes through a different interferometer with a different etalon, resulting in a different VPF in the second channel. Since the same velocity must be inferred from the two channels with different VPF-s, the set of $2\pi$-multiples that can produce a mathematically possible solution is greatly reduced. The etalons for the two channels are then chosen in such a way that only one physically plausible solution remains.

In experiments where the shock velocity cannot be measured because the probe beam is absorbed by ionized material in front of the shock, the VISAR diagnostic can be used to passively measure the time of shock breakout. Light emitted as the shock breaks out of an opaque medium is recorded in the VISAR streak camera image; an example image of a shock breakout measurement can be seen in Fig. 6.5(a). Since signal passes through the two arms of the interferometer, the etalon-imposed delay in one of the arms causes the flash of light to appear as a “double-blip” in the streak camera image, where the peaks are sepa-
Figure 6.5: (a) A sample VISAR streak camera image from a shock breakout experiment (OMEGA EP shot 35492) showing the characteristic “double-blip” signature of a shock breaking out of an opaque medium; (b) horizontal lineout of the streak camera image showing the two peaks separated by the etalon delay of 80 ps.

rated in time by the etalon delay [see the line-out in Fig. 6.5(b)]. The shock breakout times in Sec. 6.5.1 and Sec. 6.5.6 were extracted from signals of this kind by fitting a Gaussian to the second peak and subtracting the etalon delay from the time where the fitted signal reaches 2% of the peak intensity.

6.4 Radiation-hydrodynamics simulations

Post-shot simulations of all the release experiments were carried out with the 1-D radiation-hydrodynamics code LILAC [35]. The physics options used in the code were chosen to be identical to standard ICF simulations (see Sec. 2.1) in all simulations in spherical geometry while certain compromises as mentioned below had to be made in planar simulations because of limitations of LILAC. In all cases, the nonlocal model
for thermal transport [16] was used. Similarly, multi-group diffusion with 48 energy groups was used as the radiation transport model in all simulations in this chapter except where explicitly stated otherwise.

The planar release experiments feature a high degree of expansion in the tail of the rarefaction wave, which leads to some of the computational cells having extremely low densities that are poorly extrapolated in current FPEOS [36, 37] tables causing unphysical behavior in the simulations. Therefore the SESAME equation of state [61] was used in the polystyrene foil in planar simulations while the spherical simulations used FPEOS as is the case in ICF simulations.

A model for cross-beam energy transfer (CBET) [17, 18] was included in the spherical simulations while no CBET model was used in planar simulations because the planar 1-D setup does not allow correctly ray-tracing the actual beam geometry.

The accuracy of the EOS model used in the quartz witness can be validated by comparing the relationship between the shock velocity $U_{\text{shock}}$ and the post-shock pressure $P_1$ to empirical results available in the literature. The Rankine-Hugoniot (RH) relations can be derived by applying the conservation laws of mass, momentum and energy to a control volume that encompasses the shock front. The resulting equations are

$$\begin{align*}
\frac{\rho_1}{\rho_0} &= \frac{U_{\text{shock}}}{U_{\text{shock}} - u_1} \\
P_1 &= \rho_0 U_{\text{shock}} u_1 \\
E_1 - E_0 &= \frac{1}{2} P_1 \left( \frac{1}{\rho_0} - \frac{1}{\rho_1} \right)
\end{align*}$$

(6.8)

where the subscripts 0 and 1 refer to conditions in front of and behind
the shock front, respectively; the particle velocity is denoted by $u$, mass density is $\rho$ and $e$ is the specific internal energy. A model fitted to a combination of experimental data and molecular dynamics simulations of the shock velocity as a function of post-shock particle velocity in quartz is given in Ref. [62]:

$$U_{\text{shock}} = a + bu_1 - cu_1 e^{-du_1},$$

(6.9)

where the coefficients $a$, $b$, $c$ and $d$ are obtained from a fit to experimental measurements. Substituting this relation into the RH momentum equation in Eq. (6.8) yields the relationship between shock pressure and shock velocity.

Figure 6.6: A comparison between an empirical fit to data from dedicated EOS measurements in quartz [62] and the relationship between the post-shock pressure and shock velocity from a LILAC simulation using the SESAME EOS.

Figure 6.6 shows a comparison of the empirical fit for quartz given above and the relationship between shock pressure and shock velocity
from a LILAC simulation of one of the release experiments using the SESAME EOS in quartz. As can be seen in the figure, the quartz EOS in LILAC is in good agreement with prior experiments [62] and is not a significant source of error when comparing the simulated shock velocity to measurements.

Figure 6.7: A comparison of the release-driven shock velocity in the witness between FPEOS (solid black curve) and SESAME EOS (dashed blue curve) used in the polystyrene shell in spherical simulations.

Since FPEOS tables were unavailable for use in the planar simulations, a study was conducted in spherical geometry to determine the sensitivity of the simulation results to the choice of EOS of the polystyrene shell. Figure 6.7 shows the simulated velocities of the release-driven shock in the witness from spherical simulations using FPEOS (solid black curve) and SESAME EOS (dashed blue curve). The figure shows some discrepancy in the shock velocity in the witness, especially during the
early transient time. This indicates that some error is likely introduced in the planar simulations from the choice of EOS, however the magnitude of the velocity discrepancy in Fig. 6.7 is considerably lower than some of the discrepancies with respect to experimental data shown in Sec. 6.5.3 and therefore would not be explained by inaccuracies in the EOS of the polystyrene foil.

6.4.1 Mitigating the limitations of planar simulations

In addition to the aforementioned difficulties with modeling CBET, there are several other multi-dimensional effects that complicate the simulation of planar geometry compared to the more symmetric spherical geometry. For example, transverse thermal conduction effectively reduces the drive with respect to a pure 1-D case. Also, self-generated magnetic fields in the coronal plasma can reduce the thermal conductivity thereby reducing the drive. While magnetic fields can be generated in the spherical case as well, the high degree of radial symmetry should lead to less vorticity in the corona and therefore weaker magnetic fields.

To account for the effect of transverse thermal conduction the 1-D simulations were calibrated by reducing the laser drive to match 2-D DRACO [63] simulations. Figure 6.8 shows with solid lines the shock trajectories simulated in DRACO in both the 50 \( \mu \text{m} \) (blue) and 125 \( \mu \text{m} \) (black) polystyrene foils. It was found that a 20\% reduction in the laser drive was necessary to match to the DRACO results. The results from the calibrated LILAC runs are shown in Fig. 6.8 with the dashed lines for both target thicknesses.

To account for other multi-dimensional effects, or indeed missing or
Figure 6.8: Calibrating LILAC to DRACO to account for transverse thermal conduction. Shows the shock trajectories simulated by DRACO with solid lines for both the 50 \( \mu \text{m} \) (blue) and 125 \( \mu \text{m} \) (black) polystyrene foils. The dashed lines indicate the calibrated LILAC simulations for the 50 \( \mu \text{m} \) and 125 \( \mu \text{m} \) case in blue and black, respectively.

inaccurate 1-D physics in the code, a set of dedicated shock breakout measurements were performed corresponding to each laser pulse shape and foil thickness combination that was used in the release experiments. The results of these measurements were used to further calibrate the LILAC simulations by additional adjustments to the laser drive. The experimental results and subsequent calibration of LILAC are discussed in Sec. 6.5.1.

6.4.2 Simulation results

The dynamics in the release experiment are illustrated in Fig. 6.9, which shows four snapshots of the density and pressure profiles taken at key
Figure 6.9: Density (black lines) and pressure (red lines) profiles from LILAC at different times in the simulation: (a) the initial conditions; (b) at $t = 0.8$ ns; (c) at $t = 1.7$ ns; and (d) at $t = 4.0$ ns. The case shown is that of a 50 $\mu$m foil driven by the high intensity pulse (see Table 6.1). The front of the return shock that defines the material responsible for driving the shock inside the witness is called out in subfigure (d).

Figure 6.9(a) shows the initial setup at $t = 0$ with the 50 $\mu$m polystyrene foil on the right of the figure, separated by a 100 $\mu$m vacuum gap from the quartz witness foil on the left. In this view the foil is driven by the laser from the right-hand side with a 500 ps duration square pulse with an intensity of $3.0 \times 10^{14}$ W/cm$^2$.

At $t = 0.8$ ns, in Fig. 6.9(b), the strong laser-driven shock is seen inside the polystyrene foil. At this moment in time, well before the shock breakout time at 1.0 ns, the laser is already off, ensuring that the shock release remains a pure rarefaction wave. A pre-expansion caused by the radiation preheat of the laser-side surface of the quartz foil is visible in Figure 6.9(b). It is important to minimize this pre-expansion since it can affect the VISAR signal as will be discussed in Sec. 6.5.2 and 6.5.5.
At $t = 1.7 \text{ ns}$ (Fig. 6.9(c)) the shock has broken out of the rear surface of the polystyrene foil, generating a rarefaction wave travelling back toward the front surface. The shock release - the tail of the rarefaction wave - is seen traversing the vacuum gap toward the quartz witness foil. Due to the limitations of LILAC, the vacuum region is modeled as a low density gas. The accuracy of this approximation was checked via a convergence study with respect to the gas density as well as comparisons of the dynamics of the release traveling in the low density gas with a release into vacuum without the witness foil present. While the presence of the gas has a small effect on the profile of the leading edge of the release, the error in the shock dynamics inside the witness foil was found to be negligible (see Sec. 6.4.3).

The profiles at $t = 4.0 \text{ ns}$ are shown in figure 6.9(d), where the release material has collided with the quartz witness foil and generated a 6 Mbar shock inside the witness. The strength of the shock is such that the ionized shock front is expected to become a good reflector of the VISAR probe beam thereby enabling measurement of the shock velocity - typically a shock pressure greater than 1 Mbar is required in quartz.

Figure 6.10 shows the simulated time-history of the shock velocity inside the quartz witness foil. As more and more release material collides with the witness, the shock initially experiences a period of rapid acceleration of around 1 ns until the velocity reaches a peak and starts to slowly decelerate. The shock breaks out of the rear side of the 100 $\mu$m witness foil at 8 ns.

Since the time-dependent shock velocity is sensitive to the amount of release material that has collided with the witness foil at a particular
Figure 6.10: Velocity of the release-driven shock inside the witness as a function of time. The shock velocity is extracted from a planar simulation of a 125 $\mu$m thick target driven by the low intensity pulse shape (see Table 6.1).

moment in time, the experimental platform can be used to probe the kinetic energy contained in different spatial regions of the release. The relationship between the shock strength and the kinetic energy distribution in the release is illustrated in Fig. 6.11, where the release-driven shock pressure is plotted against the amount of release material that is responsible for driving the shock. This stagnated release mass is defined as the material bounded by the reflected shock from the collision with the witness, which is traveling back to the right as seen in figure 6.9(d). In figure 6.11 the stagnated release mass is given in terms of the thickness of the region of the initial polystyrene foil where the stagnated material originates from. As can be seen in figure 6.11, the experimental platform probes the kinetic energy in the release of several microns.
Figure 6.11: Pressure of the release-driven shock inside the witness as a function of the amount of release mass that is responsible for driving the shock given as the thickness at the rear surface of the initial foil where the material originated from. The defining criterion for the mass responsible for the shock is shown in Fig. 6.9. The shock pressure is extracted from a simulation of a 50 \( \mu \text{m} \) thick target driven by the high intensity pulse shape (see Table 6.1).

of the foil. This is in contrast to previous experiments where only the material released within \( \sim 0.1 \mu \text{m} \) from the surface was studied.

6.4.3 Simulation convergence

As mentioned in Sec. 6.4.2, the LILAC simulations of the release experiments used low density gas as a surrogate for vacuum. This was done because of a lack of capability to model vacuum cells in LILAC. Convergence studies were conducted in order to validate the accuracy of this approximation, the results of which are shown in Fig. 6.12.

Figure 6.12(a) shows the convergence of density profiles in simulations of a release into a low density gas compared to a release into
Figure 6.12: Convergence of simulations with decreasing pressure in the low density gas cells that act as a surrogate to vacuum. (a) Simulated density profiles of shock release into vacuum (solid black line) and low density gas at varying initial gas pressures. (b) Simulated velocity of the release-driven shock in the witness foil at varying initial gas pressures.

vacuum. Note that even though LILAC is not capable of simulating vacuum cells, vacuum boundary conditions are implemented and therefore the density profile of shock release into vacuum can be simulated. The figure shows that the density profile are sufficiently converged at low gas pressures except at very low densities where the kinetic energy of the release material is not enough to materially influence the shock velocity in the witness foil. A gas pressure of 0.2 atm was used in all the simulations of the release experiments discussed in this thesis.

A qualitative difference between having a vacuum gap or one filled with low density gas is the weak transmitted shock in the case with gas in the gap. It is therefore important to ensure that the density of the gas is low enough that the transmitted shock does not affect the predicted velocity time-history of the release-driven shock in the witness, which is the important measurable quantity in the experiments. Figure 6.12(b)
shows the predicted shock velocity in the quartz witness for varying initial gas pressures, showing that the simulated velocity is converged at 0.2 atm gas pressure.

![Figure 6.13: Convergence of the shock velocity time-history in the witness as a function of the grid size used in the simulations.](image)

Convergence studies with respect to the grid size were also performed. In order to ensure that the high gradients that occur at the ablation surface, at the rear surface of the polystyrene foil during shock breakout and at the laser-side surface of the witness foil during the release collision, a non-uniform grid was used in the simulations with highly refined regions near the foil surfaces. The results from a convergence study with varying cell sizes are shown in Fig. 6.13, where the ratios of sizes of neighbouring cells was kept constant while the number of cells was varied. The curves are labeled by the size of the smallest cell in each simulation. The results show that the velocity of the release-
driven shock in the witness foil is converged at the 0.01 µm cell size that was used in all the simulations shown in the rest of this chapter.

### 6.5 Experimental results

#### 6.5.1 Shock breakout calibration in planar geometry

Dedicated shock breakout experiments were performed in order to isolate the shock release physics from the modeling of shock transit. As pointed out in Sec. 6.4.1, there are numerous reasons that complicate simulating the transit of laser-driven shocks, especially in planar geometry. For this reason, LILAC simulations were calibrated by adjusting the laser drive to match the results of the shock breakout experiments.

Typically for VISAR measurements of laser-driven shocks, the VISAR signal “blanks” while the laser is turned on because of the radiation from the corona ionizing the material in front of the shock, which absorbs the probe beam. If the laser pulse is short enough, the material “recovers” from the preheat after some time has elapsed from laser cut-off and the velocity of the shock can be measured (See, for example, Ref. [64]). In all but one of the experiments discussed here the recovery did not occur and only the flash caused by shock breakout (see Sec. 6.3) was observed by the VISAR instrument. Therefore, the measurement used to calibrate the LILAC simulations was the time of shock breakout and not the shock velocity time-history.

Figure 6.14 shows the measured shock breakout time from all shock breakout experiments, which were carried out for each combination of target type and laser pulse shape used in the release experiments. The
laser drive in LILAC was subsequently adjusted to calibrate the simulations. A reduction in laser intensity of 30 % over the entire duration of the pulse was found to best reproduce the measured shock breakout times. Recall that this calibration is applied on top of the correction for transverse thermal conduction determined from 2-D DRACO simulations (see Sec. 6.4.1). The results from the calibrated LILAC simulations are compared to the measurements in Fig. 6.14. The measured shock breakout times are reproduced in the simulations to within 50 ps.

As mentioned, measurement of the shock velocity was unsuccessful in most of the shock breakout experiments. The single exception was shot 35498 - the low intensity shot with a 50 µm polystyrene foil. The VISAR measurement of the shock velocity in the polystyrene foil in shot

Figure 6.14: Comparison of VISAR measurements of shock breakout times \( t_{sb} \) from dedicated calibration experiments to simulated values from LILAC. The dashed black line is a guide to the eye indicating perfect agreement between simulations and experiments.
Figure 6.15: VISAR measurement of shock velocity in the polystyrene foil in the calibration shot 35498 compared to the post-shot LILAC simulation.

35498 is shown in Fig. 6.15 with the solid black curve and the simulated shock velocity is shown with the dashed blue curve. The calibration from Fig. 6.14 has been applied to the LILAC simulation shown here. The simulated shock velocity shows good agreement with the measurement in both the velocity magnitude and the rate of deceleration of the unsupported shock. Overall the calibrated LILAC simulations can reproduce the measurements, although in light of the results shown in Sec. 6.5.3, the determination of the shock velocity in the high intensity cases would greatly reduce the uncertainty that the plasma conditions at shock breakout are indeed accurately captured in the calibrated simulations at high intensity.
6.5.2 Radiation preheat issues in planar geometry

In high intensity planar release experiments (with laser intensities at \(3.0 \times 10^{14}\) W/cm\(^2\) and above) the VISAR signal was found to be significantly compromised by radiation preheat of the witness foil. The radiation emitted from the corona preheats the material at the laser-side surface of the quartz foil, which causes expansion prior to the release collision with the witness as seen in the simulated density profiles in Fig. 6.9(b) and Fig. 6.9(c).

![Figure 6.16: A side-on 4\(\omega\) probe image from a planar release experiment showing the leading edge of the shock release traversing the gap between the foils. The image was taken at a time after the shock breakout from the polystyrene foil but prior to the collision with the quartz witness foil. The blue and red annotated lines indicate the initial positions of the polystyrene and quartz foils, respectively. The release material can be seen in the vacuum gap as indicated and the pre-expansion of the quartz witness caused by radiation preheat is clearly visible.](image)

The pre-expansion of the witness was also observed in experiments as seen in Fig. 6.16, which shows a side-on 4\(\omega\) probe image. The image was taken after the shock breakout from the polystyrene shell while the release material was traversing the vacuum gap. Clear expansion of the witness foil prior to the release collision can be seen in Fig. 6.16,
where the initial surface position of the quartz witness is marked by the red rectangle. As a result of preheating, free electrons in the expanded layer cause absorption of the VISAR probe beam – “blanking” the signal until the release-driven shock has passed through the preheated region of the witness. This can be seen in Fig. 6.17 that shows a comparison of VISAR streak camera images from experiments with low and high intensity pulse shapes. The low intensity case in Fig. 6.17(a) shows an initial period of rapid acceleration of the shock as an increasing amount of material accumulates in the collision against the witness, reaching a peak velocity $\approx 1$ ns after the collision, followed by a slow deceleration until the shock finally reaches the rear surface of the quartz foil. This velocity time-history is in good qualitative agreement with the LILAC simulation in Fig. 6.10.

However, in the high intensity case shown in Fig. 6.17(b), where the longer-duration pulse shape was used with a 125 $\mu$m foil, the acceleration phase is not visible because of the aforementioned blanking issue. In this case the late-time decelerating shock velocity can still be compared to simulations while the kinetic energy in the leading part of the shock release cannot be measured.

Figure 6.17(c) shows a VISAR streak camera image from a high intensity experiment with a 50 $\mu$m foil, where the measurement was less-affected by the radiation preheat than in the 125 $\mu$m case of Fig. 6.17(b). This is because of less radiation energy per unit mass delivered due to the considerable shorter pulse shape used. While the preheated region does not completely absorb the probe beam in Fig. 6.17(c), the signal is still distorted during the time when the release-driven shock is trav-
elging through the pre-expanded region. It is likely that the existence of a gradient in density, and therefore a gradient in the refractive index in the preheated region causes the appearance of fringe shifts in the image that do not correspond to a change in the actual velocity of the shock front. Here, as in Sec. 6.5.1, a more careful analysis [65] that takes into account the reflectivity in a nonuniform target would be required and therefore only the shock velocity measured while the shock is propagating inside an undisturbed region of the witness is compared to simulation results.
Figure 6.17: The effect of varying levels of radiation preheat in VISAR streak camera images: (a) an initial period of acceleration is seen in the low intensity experiment on the 125 \( \mu \text{m} \) target indicating low preheat of the witness; (b) the early signal after the release collision with the witness is lost as the preheated quartz layer absorbs the VISAR probe in the high intensity experiment on the 125 \( \mu \text{m} \) target; and (c) the initial signal after the release collision is significantly distorted as the shock travels through the preheated quartz layer in the high intensity experiment on the 50 \( \mu \text{m} \) target.
6.5.3 Release measurements in planar geometry

As discussed in Sec. 6.1, the shock release influences the outcome of an ICF implosion by setting the initial hot spot pressure as the kinetic energy of the release material is converted to internal energy of the hot spot upon stagnation of the material. Similarly, in the release experiments described in this chapter, the kinetic energy of the shock release is converted into internal energy as the material stagnates against the witness. The pressure behind the shock driven by the release collision is a measure of this internal energy and therefore an indirect measure of the release kinetic energy – the relevant quantity that determines the performance of an ICF implosion.

VISAR measurements of the release-driven shock velocity, which can be related to the shock pressure via the Hugoniot relations as discussed in Sec. 6.4, were obtained for all the planar target and laser pulse shape configurations listed in Sec. 6.2.3. The results are shown in Fig. 6.18, where the shock velocities extracted from the VISAR streak camera images are shown with solid black curves. Here, all results shown are from the experiments with a 100 µm vacuum gap, no qualitative difference was found in the cases with the larger 170 µm gap size. Simulated velocity curves from the calibrated LILAC simulations (see Sec. 6.5.1) are shown in Fig. 6.18 with the dashed blue curves.

As mentioned in Sec. 6.3, the velocity magnitude resulting from analyzing the VISAR image is ambiguous in a way where potential solutions differ by specific multiples of the VPF. If the distance traveled by the reflecting surface that produces the VISAR signal (in this case the shock
front in the witness) is known, then the integral of the velocity curve with respect to time can be compared to the known thickness of the target to resolve the ambiguity. For this reason, the dotted orange curves in Fig. 6.18 are provided as a rough guide to the eye, indicating an area under the curve, which, combined with the integral of the measured velocity, matches the 100 $\mu$m thickness of the quartz foil.

Figures 6.18(a), 6.18(b) and 6.18(c) all indicate good agreement between the simulations and experiments with the kinetic energy of the release material $E_k \sim U_{\text{shock}}^2$ within 15% of the measurements. These measurements correspond to both high and low intensity experiments with the thinner 50 $\mu$m target as well as the case of a 125 $\mu$m target driven by the low intensity laser pulse. A short period of seemingly rapid deceleration can be seen in the high intensity, thin foil case in Fig. 6.18(b) caused by the shock front traveling through a preheated quartz layer as discussed in Sec. 6.5.2. This part of the signal is disregarded from the analysis since the observed fringe movement does not directly correspond to the shock velocity. The dotted orange curve shown as a visual aide further confirms the unphysical nature of the early rapidly decelerating signal because a naive integration of the VISAR signal yields a distance traveled that is greater than the thickness of the witness foil.

In contrast to the cases discussed in the previous paragraph, the experiments performed with a high intensity pulse on a thick 125 $\mu$m target showed significant disagreement with the simulation predictions as seen in Fig. 6.18(d). Some of the more likely causes for this observed discrepancy are considered in the next section and sensitivity of LILAC simulations to modifications in the physics models is used to judge the
feasibility of the hypothesized explanations where possible.

To assess the repeatability of the experiments, repeated shots were performed with the high intensity pulse shapes using both 50 \( \mu \text{m} \) and 125 \( \mu \text{m} \) foils. In addition, repeated shots of the shock breakout experiments were performed with high intensity on both target types. The measured shock breakout times of the repeated experiments are included in the data in Fig. 6.14. Repeated experiments with low intensity pulse shapes were not performed because of a lack of shot availability. Figure 6.19 shows the excellent repeatability seen in the release experiments. The figure shows the measured shock velocities from two high intensity experiments on 125 \( \mu \text{m} \) targets with a 170 \( \mu \text{m} \) vacuum gap. The good agreement between the repeated experiments shown in Fig. 6.19 is representative of all the performed repeats.
Figure 6.18: VISAR measurements of the shock velocity (solid black curves) in
the quartz witness foil in (a) the low intensity experiment on the 50 \(\mu\)m target;
(b) the high intensity experiment on the 50 \(\mu\)m target; (c) the low intensity
experiment on the 125 \(\mu\)m target; and (d) the high intensity experiment on the
125 \(\mu\)m target. Post-shot predicted shock velocities from LILAC simulations
are shown with the dashed blue curves. The dotted orange lines indicate an
area under the curve that corresponds to the 100 \(\mu\)m thickness of the witness
foil.
Figure 6.19: VISAR measurements of shock velocity in the quartz witness from the repeated high intensity experiments on the 125 μm target with a 170 μm vacuum gap.
6.5.4 Discussion of the results of the planar experiments

The material properties and laser intensities involved in the release experiments, as in most high energy density physics experiments, are rather extreme and require the inclusion of complex physics models in simulations. This section explores the possibility that the discrepancy between the simulation and experiment seen in Fig. 6.18(d) could be the result of inaccuracies in the details of the radiation transport and the coupling of laser energy to the target.

![Shock Velocity vs Time](image)

Figure 6.20: Sensitivity to changes in physics models of the velocity time-history in the witness foil in LILAC simulations. The results from the standard simulation are shown (solid blue line) compared to a simulation with an extended laser pulse (dashed orange line) and a simulation using the $S_n$ radiation transport model (dotted green line).

A series of simulations was used to investigate the effect of variations in the physics models by switching the models and parameters involved with the effects under consideration. The sensitivity of the results is
assessed by observing the response of the velocity of the release-driven shock in the witness – the most important observable quantity in the experiments. The results of these numerical studies are shown in Fig. 6.20. The blue solid line indicates the predicted shock velocity from the standard simulations that were compared to the experimental results in Sec. 6.5.3.

As was noted in Sec. 6.5.1, the calibrated LILAC simulations were not confirmed to produce the correct plasma conditions at shock breakout in the high intensity case because of VISAR blanking in the experiments. It is possible that the density and pressure profiles at shock breakout could be substantially different from LILAC predictions. To illustrate this fact, the following numerical experiment was performed. In the simulation of the 125 µm foil driven by the high intensity laser pulse, the velocity of the laser-driven shock is not influenced by changes in the conditions at the ablation surface after the time of laser cutoff because ion acoustic waves generated at that time, traveling at the speed of sound, are unable to reach the shock front before breakout occurs. Therefore, the pulse shape can be extended in the simulation without influencing the time of shock breakout. The dashed orange line in Fig. 6.20 shows the shock velocity in the quartz witness from a simulation where the pulse shape was artificially extended by 400 ps. Even though the shock breakout time in this simulation remained constant with respect to the standard simulation and the post-shock pressure right at the breakout surface did not increase, there was significantly more kinetic energy in the shock release as evidenced by the large increase in the release-driven shock velocity. This is because in the case with the extended
Figure 6.21: Comparison of density and pressure profiles at shock breakout in the case of the standard pulse shape (solid line) and a pulse shape extended by 400 ps (dashed line).

pulse shape the laser-driven shock remains supported by the ablation pressure for a longer time and that results in dramatically different density and pressure profiles at shock breakout as shown in Figure 6.21, where the profiles are compared to ones from the standard simulation. Essentially, the figure shows that more energy is imparted to the material located some distance away from the breakout surface resulting in more kinetic energy in the bulk of the shock release. This study is not meant to suggest that there would be any physics reason why the pulse shape in the simulations ought to be extended but merely to illustrate the sensitivity of the release kinetic energy to the plasma conditions at shock breakout. While there is no known way to mitigate the preheat issues that prevent the measurement of the laser-driven shock velocity in
the high intensity experiments, future experiments with stepped targets could be performed to obtain a better estimate of the shock velocity in the same way as done in the passive velocity measurements mentioned in Sec. 6.3.

The sensitivity of the release-driven shock velocity to the choice of radiation transport model was studied by running simulations using the $S_n$ transport model [66] instead of the standard multigroup diffusion approximation. Differences in radiation transport could modify the radiation preheat of the laser-driven foil, which can affect the dynamics of shock release. The results of the simulations are shown in Fig. 6.20 with the dotted green curve. As evident in the figure, the kinetic energy of the bulk of the shock release is not sensitive to the radiation transport model used.

Kinetic effects, such as species separation have been proposed in the literature (e.g. Ref. [26]) to explain the results from prior shock release experiments [56]. It is not clear if such effects could explain the large difference in the kinetic energy of the relatively large amount of material (compared to prior experiments) probed in the experiments in Sec. 6.5.3. To study the sensitivity of the release-driven shock velocity to various kinetic effects requires running sophisticated simulations including more complex physics than were present in the simulations performed as a part of this thesis. At this time, such effort has not been made, but could be considered in future work.
6.5.5 Radiation preheat in spherical geometry

Figure 6.22: VISAR streak camera image from shot 104206 showing a well-defined acceleration phase after the release collision with the witness. The existence of a low level of radiation preheat is indicated by the absence of both the signals of shock motion inside the polystyrene shell as well as the flash of light caused by shock breakout from the shell.

Because of the convergence of the release material in the spherical target the same shock pressure in the witness can be achieved with comparatively less energy in the laser drive than in the planar case. The use of short picket pulses with 100 ps duration in place of the square pulses with several hundred picoseconds duration results in significantly less
preheating of the witness. The distortion of the VISAR signal caused by radiation preheat as highlighted in Sec. 6.5.2 in planar geometry were not observed in any of the spherical experiments.

Evidence of some preheat can still be seen in the VISAR images, an example of which is shown in Fig. 6.22. As was the case in the planar experiments, the reflection from the laser-driven shock front traversing the polystyrene shell is not visible in the image. The most likely explanation being once more that the VISAR probe beam is absorbed in the preheated outer layer of the fused silica hemisphere. Furthermore, the flash of light expected to accompany the breakout of the laser-driven shock from the rear surface of the polystyrene shell is only visible in the shock breakout experiments absent the fused silica hemisphere but not visible in the release experiments. For this reason, dedicated shock breakout experiments were still required in spherical geometry using targets without the fused silica hemisphere as described in Sec. 6.2.4.

Figure 6.23: Sketch of the target used in the preheat experiment. A high-Z gold layer is buried in the polystyrene shell to shield the fused silica witness hemisphere from the coronal radiation.
Since radiation preheat was a concern prior to the spherical experiments based on the signals observed in the planar case, special targets with a buried gold layer were fielded to study the effectiveness of the high-Z gold layer in shielding the witness from the coronal radiation. The release target with a buried gold layer is sketched in Fig. 6.23 where a 0.9 \( \mu \text{m} \) thick gold layer is shown inside a 18.6 \( \mu \text{m} \) thick polystyrene shell. Simulations performed with LILAC show a significant reduction in the radiation energy absorbed in the fused silica witness because of the presence of the gold layer.

Figure 6.24 shows the VISAR streak camera image from OMEGA shot 104216, where the target with a buried gold layer shown in Fig. 6.23 was used. For a fair comparison of preheat levels, since the buried layer alters the shock transit in the shell, a 30 J/beam picket pulse was applied instead of the 22 J/beam picket used on the 30 \( \mu \text{m} \) pure polystyrene target. The pulse shape was designed to achieve the same plasma conditions at shock breakout as in the pure polystyrene experiments. A clear reduction in the radiation preheat of the witness hemisphere can be seen in Fig. 6.24 indicated by the presence of the fringe motion at 200 to 300 ps that corresponds to the reflection of the VISAR probe beam off of the laser-driven shock traveling in the shell as well as the shock breakout signal seen at around 300 ps in the image.

The clear evidence of a reduction of radiation preheat of the witness seen in Fig. 6.24 confirms the effectiveness of using a buried high-Z layer in the release experiments. Therefore, targets with such a shielding layer should be considered for future release experiments in planar geometry while the VISAR signal in spherical experiments is not signif-
icantly affected by the preheat even without the high-Z shielding layer. The introduction of a high density buried layer has a significant impact on the shock transit in the shell (or the laser-driven foil in planar geometry), introducing complex shock interactions inside the target and the additional uncertainty of having to model the EOS of a new material. Care must then be taken to acquire the appropriate data on shock transit in the target with a buried layer in order to ensure proper calibration of the simulations.

While the VISAR images of the shock release experiments shown in this section were successfully used to assess the qualitative effect of the buried gold layer on the radiation preheat, these data can unfortunately not be used to validate the modeling of kinetic energy of the shock release in LILAC simulations. Because of an unfortunate incident of data loss, the measurements of the laser energy delivered on target in shot 104216 were not stored and therefore accurate post-shot simulations could not be run.
Figure 6.24: VISAR streak camera image from the preheat experiment (OMEGA shot 104216). The acceleration phase is seen in the reflection of the probe beam off of the shock front in the witness hemisphere similarly to the experiment with the polystyrene shell in Fig. 6.22. Because of the reduced preheat due to the addition of the shielding gold layer the signal of the shock motion inside the polystyrene shell is also visible.
6.5.6 Shock breakout measurements in spherical geometry

Dedicated shock breakout experiments were performed on cone-in-shell targets with a polystyrene shell thickness of 30 $\mu$m. Both the low and high intensity picket pulses as used in the release experiments specified in Sec. 6.2.3 were used in the shock breakout experiments. The reduced preheat with respect to the planar case, as mentioned in the previous section, enabled measurements of shock velocity while in transit inside the shell in the spherical shock breakout experiments. The resulting temporal velocity curves are shown with solid lines in Fig. 6.25 for the experiment with a low intensity pulse.

Because of technical difficulties during the experiments, the mea-
Table 6.3: Measured shock breakout times from spherical experiments compared to predictions from LILAC simulations

<table>
<thead>
<tr>
<th>Description</th>
<th>Measured (ps)</th>
<th>LILAC (ps)</th>
</tr>
</thead>
<tbody>
<tr>
<td>Low intensity</td>
<td>480</td>
<td>550</td>
</tr>
<tr>
<td>High intensity</td>
<td>370</td>
<td>510</td>
</tr>
</tbody>
</table>

Measurements from the two different channels of VISAR had to be acquired on two separate shots. While the repeatability of the experiments was found to be excellent just as in the planar experiments in Sec. 6.5.3, the subtle shot-to-shot variations in the system resulted in an ambiguity when extracting the velocity curve from the VISAR images. For this reason, two possible solutions corresponding to the two experiments are shown in Fig. 6.25. The LILAC simulated shock velocities are shown in Fig. 6.25 with the dashed blue curve. Based on the fact that the shock breakout occurs earlier in the experiments than predicted by LILAC, it is more likely that the correct velocity measurement ought to have higher shock velocity than LILAC, corresponding to the solid black curve in Fig. 6.25.

The measured and LILAC-simulated shock breakout times for these experiments are given in Table 6.3. In both the low and high intensity cases, LILAC predicts the shock breakout to occur later than what was measured in the experiment, with a larger discrepancy seen in the high intensity case. This is in contrast to the planar experiments in Sec. 6.5.1, where the laser drive had to be adjusted lower in order to match the shock breakout times. Many of the possible causes for the inaccuracies in modeling shock transit in the planar case are either already taken into account in the spherical simulations (e.g., a model for CBET
is included) or are not expected to have a significant impact in spherical geometry (e.g., transverse thermal conduction losses and magnetic fields in the corona are much weaker in the more symmetric spherical experiments). For this reason, the spherical simulations were not calibrated to match the shock breakout measurements, but were instead set up identically to standard ICF simulations.

6.5.7 Release measurements in spherical geometry

![Graphs](a) and (b)

Figure 6.26: VISAR measurements of the shock velocity (solid black curves) in the fused silica witness hemisphere in experiments using: (a) the low intensity picket pulse; and (b) the high intensity picket pulse. Post-shot predicted shock velocities from LILAC simulations are shown with the dashed blue curves.

Shock release experiments were carried out with 30 \( \mu \text{m} \) thick polystyrene cone-in-shell targets using both low and high intensity pulse shapes as specified in Sec. 6.2.3. VISAR measurements of the release-driven shock velocity in the fused silica witness are shown with solid black curves in Fig. 6.26(a) and Fig. 6.26(b) for the low and high intensity cases, respectively. The VISAR streak camera images from which the
shock velocity was extracted are shown in Fig. 6.26(c) for the low intensity case and in Fig. 6.26(d) for the high intensity case. The lack of significant preheat of the witness is apparent in the clear acceleration phase at the beginning of the signal in both the low and high intensity cases, indicating that the VISAR instrument is able to capture the shock motion over nearly the entire initial period after the release collision with the witness.

Note how, in Fig. 6.26(c) and Fig. 6.26(d), the spatial extent of the signal that is reflected off of the shock front in the witness is very narrow, consisting of 100 pixels (compare to the images from the planar experiments in Sec. 6.5.2). This is caused by the VISAR probe beam reflecting off of a curved shock front which means that only the reflection from the central region that is perpendicular to the probe beam makes it’s way back to the instrument. Based on these figures, the curvature of the hemisphere with a 250 $\mu$m radius is close to the limit where any increase in curvature could make it very difficult to extract useful information out of the signal. It is therefore not advisable to use smaller hemispheres in future experiments with a similar setup.

The LILAC-predicted velocity curves are shown in Fig. 6.26 with the dashed blue lines. Considering the discrepancy between the measured and simulated shock breakout times, the agreement in the velocity of the release-driven shock in Fig. 6.26 is surprisingly good. The time of the collision of the release material with the witness appears especially well-predicted as well as the overall shape of the velocity curve, while the magnitude of the shock velocity and therefore the kinetic energy in the release material is somewhat overpredicted by LILAC. The magni-
tude of this discrepancy in the kinetic energy of 20% is less than that seen in the poorly predicted planar case in Sec. 6.5.3. However, as seen in Sec. 6.5.6, LILAC predicts a significantly slower shock transit through the polystyrene shell, suggesting that if the plasma conditions at shock breakout were to be matched then that would increase the discrepancy seen in Fig. 6.26. For this reason, it cannot be concluded that the shock release dynamics in these experiments are accurately captured by LILAC. Interestingly, the kinetic energy in the spherical release experiment is overpredicted by LILAC while in the anomalous planar case the kinetic energy was significantly underpredicted, suggesting that the discrepancies in the two cases could be caused by different physical effects.

6.6 Conclusions

Shock release experiments to measure the kinetic energy of the bulk of the release material where conducted in both planar and spherical geometry. The release dynamics were studied over a range of ICF relevant conditions by varying the target parameters and laser intensity.

Good agreement between measurements and calibrated simulations was found in most cases in planar geometry with the important exception of the high intensity experiment on a 125 µm thick foil, where the measured kinetic energy significantly exceeded the simulation predictions. The sensitivity of the simulations was explored for several proposed sources of the discrepancy. The kinetic energy of the release was found to be insensitive to details of the EOS and radiation transport mod-
els, contrary to results from prior experiments which focused on the low density leading edge of the shock release. The simulations exhibited a high sensitivity to plasma conditions at shock breakout which could be inaccurately captured by the codes because of the complicated physics involved in the absorption of the laser drive and thermal conduction in the coronal plasma, especially in planar geometry. More sophisticated dedicated experiments are required in the future to develop a better understanding of the shock transit in the release experiments.

Shock release experiments in spherical geometry were found to be in relatively good agreement with simulations using the physics models included in standard ICF simulations. Discrepancies in the predictions of the shock breakout time suggest, however, that a more accurate understanding of the shock transit in the shell is necessary. Considering that the kinetic energy of the shock release in LILAC simulations was predicted to be higher than what was measured in the experiments the results from the spherical release campaign indicate that the lower than expected convergence observed in ICF experiments is not caused by the shock release.
CHAPTER 7
CONCLUDING REMARKS

“...My propositions serve as elucidations in the following way: anyone who understands me eventually recognizes them as nonsensical, when he has used them – as steps – to climb up beyond them. (He must, so to speak, throw away the ladder after he has climbed up it.)”

- Ludwig Wittgenstein

The statistical mapping model as first presented in Ref. [11] began as a black box model, unable to elucidate the underlying physics of the degradation mechanisms in ICF implosions, but was nevertheless able to provide unprecedented prediction accuracy, especially in the region of parameter space in close vicinity to prior experiments. The OMEGA ICF performance campaign, guided by the early statistical mapping model, accomplished a quadrupling of fusion yield at a moderately high areal density thereby making great progress toward the goal of hydroequivalent ignition. Still, the major problem with a purely statistical approach is the tendency of the predictions to suddenly fail in unexpected ways when venturing too far from prior data. This is often caused by spurious correlations between parameters due to inherent biases in the data. As a real-life example from the OMEGA database, a statistical mapping model could be fit to the measured fusion yield using the length of the pulse shape from $t = 0$ to laser cut-off as one of the parameters of the model. The resulting mapping would indicate
a degradation in fusion yield with increasing pulse shape length. The resulting conclusion that there is a degradation mechanism related to the time-scale of the implosion would, however, be a mistake. There is a very strong inverse correlation in the OMEGA database between the shell adiabat and the pulse shape length, which is caused by the way that nearly all ICF pulse shapes are designed. In order to set a low adiabat in the shell, the initial shock wave launched by the picket must be weak and therefore the shock will travel slowly through the shell, which affects the optimal timing of the rest of the pulse shape. As a result, pulse shapes are significantly longer in low-adiabat designs. The detailed analysis in Sec. 3.3 revealed that the observed degradation is not related to the pulse shape length but instead is caused by the effects of hydrodynamic instabilities at low adiabats.

The physics based nature of the mapping model developed in Ch. 2 means that the predictions are more likely to generalize beyond the training data. Dependencies that are based on physical processes can be more readily extrapolated outside of the existing data and should not be as susceptible to breaking down completely compared to models based on naive correlations in a large complex dataset. The statistical mapping method described in this thesis was successfully applied to a large database of OMEGA implosions thereby uncovering the dominant degradation mechanisms of the fusion yield. A strong yield dependency on the magnitude of the $\ell = 1$ mode was found with the inferred power-law exponent in good agreement with results from previous 3-D simulations [47]. The degradation caused by $^3$He accumulation in the vapor region was found to be correlated to 1-D simulation results, although
the experimental data suggests a stronger effect than predicted by the codes. A strong degradation caused by the finite beam size was inferred, which is not fully explained by 3-D simulations of the beam mode [50]. The data suggest additional degradation sources related to the beam size such as inaccuracies in the CBET models or laser imprinting. Finally, a strong degradation caused by effects related to hydrodynamic instabilities was inferred in best performing implosions on OMEGA.

A series of controlled experiments where only the DT fill age was varied were found to be in excellent agreement with the mapping model predictions. The statistical mapping method was used to identify a path toward higher performance on OMEGA utilizing larger diameter targets to optimize the trade-offs between degradation from hydrodynamic instabilities and finite beam size effects and the increased coupling of laser energy.

Another advantage of deconstructing the measured degradations with respect to simulations in terms of specific physical mechanisms is that it enables an iterative process, where the insight gained from the statistical mapping model can be used to improve the simulations, which will in turn enhance the accuracy of the mapping model. The statistical mapping method was also applied to the areal density in the OMEGA experimental database. Because of the relatively large uncertainty of the areal density measurement the conclusions about the degradation mechanisms of $\rho R$ are not as clear as in the case of fusion yield. The mapping model suggests a strong degradation of areal density in implosion designs that have high convergence in simulations, prompting further investigation into the physical causes of the lack of convergence.
observed in the experiments. The experiments described in Ch. 6 were designed to study the effect of shock release on convergence in ICF implosions. The measurements from shock release experiments in both planar and spherical geometries were overall in relatively good agreement with simulation results with the exception of high intensity experiments performed with thick planar targets. No indication was found that the shock release could be responsible for low convergence in ICF implosions.

To end on a hopeful note, one could envision a scenario, where the iterative application of the statistical mapping approach and subsequent improvements to the codes could eventually lead to a situation where the simulations themselves possess good enough predictive capability that the mapping model would no longer be required. It can then be “thrown away” much like Wittgenstein’s metaphorical ladder in the epigraph above this chapter. Until parity in prediction accuracy has been achieved in simulations the statistical mapping model will remain a powerful tool for both design optimization as well as furthering the understanding of the physics in ICF implosions and, perhaps, in other complex systems where accurate simulations are not available.
REFERENCES


