LLE 1997 Annual Report

October 1996 – September 1997

Inertial Fusion Program and
National Laser Users' Facility Program

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Contents

Executive Summary	v
Two-Dimensional SSD on OMEGA	1
Areal Density Measurement of Laser Targets Using Absorption Lines	11
Modeling the Temporal-Pulse-Shape Dynamics of an Actively Stabilized Regenerative Amplifier for OMEGA Pulse-Shaping Applications	18
Multiple Scale Derivation of the Relativistic Ponderomotive Force	24
Subpicosecond Imaging Based on Electro-Optic Effect	36
Nuclear Diagnostics for High-Density Implosions	46
Status of Optics on the OMEGA Laser after Eighteen Months of Operation	51
A Single-Beam, Ponderomotive-Optical Trap for Energetic Free Electrons	60
A Diode-Pumped Nd:YLF Master Oscillator for the OMEGA Laser	68
Heat Transfer from Localized Absorbing Defects to the Host Coating Material in $\text{HfO}_2/\text{SiO}_2$ 351-nm High Reflectors	74
Signatures of Target Performance and Mixing in Titanium-Doped Target Implosions on OMEGA	82
Dephasing Time of an Electron Accelerated by a Laser Pulse	92
The Effect of Increased Irradiation Uniformity on Imprinting by 351-nm Laser Light	101
Single-Mode Rayleigh–Taylor Growth-Rate Measurements with the OMEGA Laser System	106
Fabrication of Polyimide Shells for Use as ICF Targets	113

Sideward Stimulated Raman Scattering of a Short Laser Pulse in a Plasma Channel	127
Picosecond Photoresponse in Polycrystalline Silicon	132
Thermal Distortions in Laser-Diode- and Flash-Lamp-Pumped Nd:YLF Laser Rods	137
The OMEGA Target-Positioning System	145
OMEGA Experimental Program and Recent Results	161
Imaging the Cold, Compressed Shell in Laser Implosions Using the $K\alpha$ Fluorescence of a Titanium Dopant	168
Near-Forward Stimulated Brillouin Scattering	177
Pulse Shaping on the OMEGA Laser System	184
Angular-Scattering Characteristics of Ferroelectric Liquid Crystal Electro-Optical Devices Operating in the TSM and ESM Modes	188
Magnetorheological Finishing of IR Materials	197
Integrated Circuit Tester Using Interferometric Imaging	209
LLE's Summer High School Research Program	214
Laser Facility Report	216
National Laser Users' Facility News	217
Publications and Conference Presentations	221

Executive Summary

The fiscal year ending September 1997 (FY97) concluded the fifth year of the cooperative agreement (DE-FC03-92SF19460) with the U. S. Department of Energy (DOE). This report summarizes research at the Laboratory for Laser Energetics (LLE) and is the final report for the first five years of the cooperative agreement. In September 1997, the cooperative agreement was renewed for an additional five years. We summarize our research during FY97, the operation of the National Laser Users' Facility (NLUF), and the education of high school, undergraduate, and graduate students in LLE programs.

Progress in Laser Fusion

A general introduction to LLE's experimental physics program and a report on recent results are found on pp. 161–167. This article includes a useful summary of the system's operational capabilities and system parameters after three years of operation. Direct-drive inertial confinement fusion requires precise drive uniformity, the control of hydrodynamic instabilities during the implosion of the fusion target, and accurate target fabrication and characterization. The article summarizes a wide variety of experiments relating to direct-drive laser fusion, from high-yield implosion experiments to planar and spherical Rayleigh–Taylor experiments, laser-imprinting experiments, and laser–plasma interaction experiments.

A detailed analysis of the equation of motion for an electron in a plane wave is presented beginning on p. 24. A guiding center model is postulated and compared to numerical simulation of the actual particle motion. The formula is also verified analytically using the method of multiple scales. Work continues on this formalism to study the effects of the ponderomotive force on laser–plasma interactions. A theoretical calculation of the dephasing time of an electron accelerated by a laser pulse is found on pp. 92–100. The trajectory of a charged particle, determined analytically for various pulse shapes, is then used to determine the dephasing time of an accelerated particle.

In other laser–plasma interaction studies, we present a theoretical investigation of the spatiotemporal evolution of

sideward stimulated Raman scattering in a plasma channel created by the ponderomotive expulsion of the plasma from the region of the laser-pulse axis (pp. 127–131). The partial reflection of Stokes light by the channel walls enhances the instability by allowing a spatial eigenmode to form in the laser pulse, which grows exponentially in time during the remainder of the pulse. The dependence of the growth rate of sideward SRS on the physical parameters is discussed. A study of the instabilities associated with near-forward stimulated Brillouin scattering can be found beginning on p. 177. It includes a calculation of the saturation times and steady-state gain exponent.

We summarize an experimental study of target performance and mixing in titanium-doped target implosions on OMEGA (pp. 82–91). Results from a recent experiment show the predicted absorption features from a thin Ti layer. These features were used to estimate the core temperature and ρR of the compressed target. In addition the EXAFS spectrum above the Ti *K* edge was observed for the first time in an implosion experiment that enabled the density of the Ti layer to be measured. These experimental techniques will be important in measuring improvements in target performance as OMEGA's uniformity improves.

The Rayleigh–Taylor instability is perhaps the most important hydrodynamic effect to control to reach ignition on the National Ignition Facility (NIF). A study of the growth of mass perturbations due to the Rayleigh–Taylor hydrodynamic instability at the ablation interface is presented on pp. 106–112. The observed growth of well-defined sinusoidal mass perturbations agreed well with data from the 2-D numerical simulation obtained using the hydrodynamic code *ORCHID*. The detailed analysis of these experiments showed that the determination of the Rayleigh–Taylor growth rate from the experimental x-ray radiograph data is problematic due to the evolution of the target density.

A technique for fabricating polyimide target shells, together with preliminary material property data for the shells, is discussed on pp. 113–126. This study represents the first

successful attempt to make shells from polyimide. Polyimide shells withstand greater buckling and bursting pressures—and are more permeable—than existing hydrocarbon shells. This material is a likely candidate to fabricate the cryogenic capsules required for experiments at Rochester and on the NIF.

Diagnostics Development

A method for measuring the areal density $\rho\Delta r$ of a compressed target shell based on the observation of absorption lines from a titanium-doped layer was tested using a simulated spectrum from a one-dimensional *LILAC* simulation (pp. 11–17). The predicted peak $\rho\Delta r$ of the compressed shell was within 17% of the value calculated directly from *LILAC*—LLE's one-dimensional laser-fusion simulation code.

Methods to measure the fuel compression in high-density implosions are described beginning on p. 46. The three ρR diagnostics being developed for OMEGA based on nuclear particles are reviewed.

A promising method to directly observe the cold compressed shell of an imploding target is described beginning on p. 168. Normally, an imploding shell is observed by backlighting. The proposal described here is to use a high- Z dopant that fluoresces under radiation from the hot core in the $K\alpha$ line. The $K\alpha$ emission signature should prove useful in tracking the improved target performance as the laser uniformity improves on OMEGA.

A detailed engineering description of the design and operation of the OMEGA target positioner can be found on pp. 145–159. This equipment allows a target to be precisely aligned using up to 4° of freedom: X , Y , Z , and ω ; a two-axis stage is available to provide θ and ϕ rotation. This positioner has been used to support and align a wide variety of targets in the experimental chamber including spherical implosion targets, multiple flat-foil instability targets, and hohlraums for indirect-drive experiments.

Laser and Optical Technology

We review (pp. 2–10) 2-D SSD as implemented on the OMEGA laser system. A summary of the detailed mathematical formalism is given together with the predicted level of uniformity achievable on OMEGA. The first experimental results on uniformity using narrow-band 2-D SSD are compared to theoretical calculations. Excellent agreement between experiment and theory is found, giving us confidence that broadband 2-D SSD with polarization wedges should achieve

an rms nonuniformity in the 1%–2% level necessary for cryogenic implosion experiments. In related uniformity of irradiation experiments, we used a new beam-smoothing device—distributed polarization rotators—in concert with existing techniques (distributed phase plates and 2-D smoothing by spectral dispersion) to improve the on-target uniformity of each beam (pp. 101–105). The result of this improved radiation uniformity was a substantive reduction in imprinting—the nonuniformity caused by the laser to drive a fusion target. A novel way to study the time dependence of this imprinting is also presented in this article

We describe modeling of the temporal-pulse-shape dynamics of the regenerative amplifier used on OMEGA (pp. 18–23). The laser rate equations are presented along with a discussion of the regenerative-amplifier dynamics. There is excellent agreement between the model's predicted output and the experimentally observed output. It is now possible to model the entire OMEGA laser system, enabling on-target pulse shapes to be specified.

A general description of OMEGA's pulse-shaping system begins on p. 184. The system, although continuously undergoing improvements, can be considered fully operational. For the variety of pulse shapes we have produced, we have demonstrated that we can accurately model the performance of the laser, conversion crystals, and other transfer functions involved in pulse-shape generation.

We summarize the new diode-pumped Nd:YLF master oscillator for the OMEGA laser (pp. 68–73). Special attention is paid to ensure long-pulse operation and high stability. Experimental results are presented that confirm the excellent amplitude stability, low timing jitter, and long-term frequency stability of the new master oscillator.

We carried out simulations of heat transfer from localized absorbing defects to the host coating material in $\text{HfO}_2/\text{SiO}_2$ 351-nm high reflectors (pp. 74–81). Atomic-force microscopy has shown laser-induced submicron cratering of UV multilayers, attributed to nanoscale, localized absorbers. The model used to simulate these absorbers is described. When thermal conduction is the only heat-transport mechanism, very high defect temperatures are required to explain the measured damage; consequently, other physical mechanisms must be considered.

The OMEGA laser driver line currently incorporates silicon switches as part of the electro-optic pulse-shaping system. An

experimental examination of the photoconductive impulse response in small-grain-polysilicon thin-film switches is described beginning on p. 132. The response was measured to have an upper limit of 3 ps and 36 ps at 0.8 μm and 1.55 μm , respectively. Understanding how to optimize this response time will improve the photoconductive switches used in the OMEGA electro-optic pulse-shaping system.

Advanced Technology

A subpicosecond electro-optic sampling system capable of imaging the voltage distribution over a rectangular region is described (pp. 36–45). The system is comparable to an ultrafast sampling oscilloscope with more than 180,000 channels and can be used to test microwave devices or digital (e.g., CMOS) circuits.

A single-beam, ponderomotive optical trap for energetic free electrons has been developed that can be used to study high-field atomic physics (pp. 60–67). Numerical results show that a phase mask can form a central intensity minimum, producing a three-dimensional ponderomotive trap. A novel segmented wave plate has been manufactured and used on the tabletop terawatt (T^3) laser to experimentally confirm the existence of this three-dimensional structure. Work continues to confirm the trapping properties of this unique optical device.

We describe the angular-scattering characteristics of ferroelectric liquid crystal electro-optical devices operating in transient and extended scattering modes (pp. 188–196). The possibility of applying these devices as modulators in practical IR imaging systems is evaluated.

A faster method of shaping and finishing IR materials by using magnetorheological fluids is explained beginning on p. 197. Magnetorheological fluid polishing of optical glasses is one of the most-promising modern techniques for achieving high-quality optical finishes. Extension of these techniques for use with IR materials is important for many laser applications. Detailed specifications and test results are included.

An integrated circuit tester based on interferometric imaging is described on pp. 209–213. This technique holds promise of ultrafast noninvasive testing of the voltage states of sections of microchips.

Laser Facility Report

The status of the optics on the OMEGA laser system after the first 18 months of operation is reviewed (pp. 51–59). A

vigorous program to monitor the performance of the optics has been followed since the inception of the OMEGA laser. The article presents results from these observations and defines the various types of possible damage to the optics. Many of the optics have not damaged, such as the frequency-conversion crystals, polarizers, calorimeters, and liquid crystal optics. The most significant damage has been sustained by the fused-silica spatial filter lenses. There has been no evidence of any propagation of damage downstream of damaged optics. In fact, after 1000 target shots the OMEGA laser has sustained remarkably little damage.

Beginning on p. 216, we summarize the laser facility operations for the fiscal year. The system was used for 722 target shots, with the balance of the 1946 shots for the year being used to effect system improvements for uniformity, diagnostic checks, and laser component testing.

National Laser Users' Facility (NLUF)

The FY97 Annual Report concludes with a description of the National Laser Users' Facility operations (pp. 217–219). During FY97 we advanced the proposal review schedule by three months, produced a new NLUF Users' Guide, and added an NLUF Web Page [<http://www.lle.rochester.edu/pr/nluf.html>] to existing LLE information on the Internet, and created a users' questionnaire to help improve operations and support for the users. The NLUF section summarizes current programs and lists the proposals approved for FY98.

Education at LLE

As the only university major participant in the National ICF Program, education continues to be a most important mission for the Laboratory. Graduate students continue to play a significant role in LLE's research activities and are participating in research using the world's most powerful ultraviolet laser for fusion research on OMEGA. Fourteen faculty from five departments collaborate with LLE's scientists and engineers. Presently 44 graduate students are pursuing Ph.D. degrees at the Laboratory. The research interests vary widely and include theoretical and experimental plasma physics, laser-matter interaction physics, high-energy-density physics, x-ray and atomic physics, nuclear fusion, ultrafast optoelectronics, high-power-laser development and applications, nonlinear optics, optical materials and optical fabrication technology, and target fabrication. Technological developments from ongoing Ph.D. research will continue to play an important role on OMEGA.

One hundred twenty-eight University of Rochester students have earned Ph.D. degrees at LLE since its founding. An additional 48 graduate students were funded by NLUF grants. The most recent University of Rochester Ph.D. graduates and their thesis titles are:

A. Lobad	<i>Femtosecond Spectroscopy of Trapping and Recombination Processes in Low-Temperature-Grown GaAs and InGaAs</i>
J. Mastrangelo	<i>Design and Synthesis of Vitriifiable Low-Molar-Mass Organic Materials</i>
L. Mu	<i>A Study of Photoconductive Switching Phenomena and an Application to Ultrafast Pulse Shaping</i>
C.-C. Wang	<i>Ultrafast Testing of Electronic/Optoelectronic Devices</i>
L. Zheng	<i>Third-Harmonic Generation of Intense Laser Pulses</i>
J. Zuegel	<i>Nonradiative Processes Affecting Laser Dynamics in Neodymium:Yttrium Lithium Fluoride.</i>

Approximately 40 University of Rochester undergraduate students participated in work or research projects at LLE this past year. Student projects include operational maintenance of the OMEGA laser system, work in the materials and optical-thin-film coating laboratories, programming, image processing, and diagnostic development. This is a unique opportunity for these students, many of whom will go on to pursue a higher degree in the area in which they have participated at the Laboratory.

LLE continues to run a Summer High School Student Research Program (pp. 214–215) where ten high school juniors spend eight weeks performing individual research projects. Each student is individually supervised by a staff scientist or an engineer. At the conclusion of the program, the students make final oral and written presentations on their work. The written reports are published as an LLE report. The 1997 program is described on pp. 214–215.

In 1997, LLE initiated an Inspirational Science Teacher Award. Alumni of our Summer High School Research Program were asked to nominate teachers who had a major role in exciting their interest in science, mathematics, and/or technology. The award, which includes a \$1000 cash prize, was presented at the High School Student Summer Research Symposium; the 1997 winner was Mr. Raymond Sherbinski from Brighton High School. Mr. Sherbinski, a mathematics teacher, was nominated by Daniel Nelson, a participant in our 1994 program. In nominating Mr. Sherbinski, Daniel wrote, "I left my year with Mr. Sherbinski with renewed vigor for not only math, but for all learning. Mr. Sherbinski creates a class atmosphere that removes a student's inhibitions toward risking vocal expression of ideas before his or her peers."

Robert L. McCrory
Director

Two-Dimensional SSD on OMEGA

The ultimate goal of the LLE uniformity program is to reduce the rms laser-irradiation nonuniformity to the 1%–2% level, which is required for cryogenic implosion experiments on OMEGA. The combination of distributed phase plates (DPP's), two-dimensional (2-D) smoothing by spectral dispersion (SSD), polarization wedges, and beam overlap should be sufficient to reach this goal. We present here a discussion of the mathematical formalism of 2-D SSD with numerical calculations illustrating the levels of uniformity that can be achieved on OMEGA. The initial implementation of 2-D SSD is described, and the initial experimental results for uniformity are compared with theory.

2-D SSD Concept

The level of uniformity that can be achieved with SSD is determined by two factors: bandwidth and spectral dispersion. The amount of bandwidth determines the rate of smoothing, and the amount of spectral dispersion determines the maximum reduction in nonuniformity that can be achieved (as well as the longest spatial wavelength of nonuniformity that can be smoothed). Frequency-tripled glass lasers place constraints on both bandwidth and spectral dispersion. Current techniques for

the high-efficiency frequency tripling of laser light limit the (full-width) bandwidth to 3 Å to 4 Å in the IR for OMEGA. Spatial-filter pinholes in the laser chain limit the spectral spread of the beam to five to ten times the beam's IR diffraction limit. With these constraints, the levels of uniformity required for OMEGA experiments can be achieved using SSD.

The starting point for a description of the uniformity that can be achieved by 2-D SSD is the speckle pattern produced by a phase plate. An example is shown in Fig. 69.1(a). It is characterized by a smooth, well-defined intensity envelope on target. However, superposed on the envelope is highly modulated intensity structure (known as speckle), which is produced by interference between light that has passed through different portions of the phase plate. SSD smoothes this speckle structure in time by progressing through a sequence of many copies of this speckle pattern, each shifted in space, so that peaks of some fill in the valleys of others at different times. When averaged in time, this effect is qualitatively similar to whole-beam deflection: 1-D SSD has the effect of sweeping the beam in only one direction, and 2-D SSD is similar to sweeping the beam in two dimensions.

Figure 69.1

The effect of overlapping a large number of statistically different speckle patterns. For a single pattern the rms nonuniformity $\sigma_{\text{rms}} = 100\%$. For N patterns, the rms nonuniformity is reduced by $1/\sqrt{N}$.

The reduction in rms nonuniformity is statistical in nature. Patterns shifted by more than about 1/2 of a speckle width act as entirely different random speckle patterns. The overlap of N random speckle patterns reduces the rms nonuniformity σ_{rms} by $1/\sqrt{N}$. Examples in Figs. 69.1(b) and 69.1(c) show how 100 and 1,000 overlapping speckle patterns reduce the phase plate σ_{rms} of 100% to values of 10% and 3%, respectively. OMEGA will require 2000 to 10,000 of such overlapping patterns, produced as follows: 2-D SSD will provide ~500 of the speckle patterns (for a 1-ns smoothing time); the polarization shifter will provide an additional factor of 2; and multiple-beam overlap will provide another factor of ~4, for a resulting nonuniformity of 1%–2%.

SSD generates these shifted speckle patterns in a two-step process.¹ The beam is passed through an electro-optic modulator, which imposes a small spread of frequencies (bandwidth) upon the laser light. The bandwidth is then spectrally dispersed by means of diffraction gratings. For 1-D SSD, one modulator is used. For 2-D SSD, two modulators (of different frequencies) are used, with diffraction gratings oriented such that each bandwidth is dispersed in a perpendicular direction. Because of the dispersion, each spectral component focuses onto the target in a slightly different position, producing the required shifted speckle patterns (Fig. 69.2).

Figure 69.2

Schematic illustrating the use of spectral dispersion to generate overlapping speckle patterns. The different spectral modes are spatially shifted in the target plane. Speckle patterns that are shifted by more than 1/2 of a speckle size are statistically independent.

The spectral dispersion does not reduce the speckle fluctuations instantaneously. It creates an entirely different speckle pattern, but one that changes in time. Interference between the electric fields from different sections of the beam will fluctuate in time because of their different frequencies, and the time-averaged interference approaches zero at a rate inversely proportional to the difference in frequency. As the contributions from interference become small, the time-averaged

uniformity “smoothes” to the result expected from superposing shifted intensity profiles, each corresponding to a different frequency.

The time-averaged uniformity approaches an asymptotic level that is determined by the number of statistically independent speckle patterns (which is generally smaller than the number of spectral components with different frequencies). This depends on the ratio between the maximum spatial shift (S_{max}) that can be produced by the laser and the smallest shift (S_{min}) that will produce statistical smoothing. The maximum spatial shift S_{max} is defined as the full-width distance in the target plane that rays in the laser beam are deflected by SSD dispersion. The smallest shift S_{min} is 1/2 of a speckle size and is given by

$$S_{\text{min}} = F\lambda/D, \quad (1)$$

where F is the focal length, λ is the UV wavelength, and D is the diameter of the focus lens. The maximum shift S_{max} is determined by the maximum angular spread of the light that can propagate through the spatial-filter pinholes of the laser (Fig. 69.3). This can be conveniently expressed as a multiple s of the whole beam diffraction limit (pinhole sizes are often expressed in terms of this parameter). Thus,

$$S_{\text{max}} = s \cdot (2.4 F\lambda/D). \quad (2)$$

Currently, OMEGA pinholes can accommodate $s = 15$, while $s = 30$ is planned for the future. (Note: In this article s is quoted in terms of UV wavelength; if expressed in terms of the IR wavelength, s should be decreased by a factor of 3.) In terms of these parameters, the number of statistically independent speckle patterns (N_{stat}) is

$$N_{\text{stat}} = (S_{\text{max}}/S_{\text{min}})^2 = (2.4s)^2, \quad (3)$$

where the ratio is squared because 2-D SSD allows spectral shifting in two directions. (This estimate for N_{stat} is somewhat simplified because it assumes that all 2-D SSD modes have different frequencies and are therefore independent; more accurate calculations are given in the next section.) The current and future values of N_{stat} are 1300 and 5000 for OMEGA, which, by itself, should reduce the rms speckle fluctuation to 3% and 1.5%, respectively, in the asymptotic limit. For 1-D SSD, $N_{\text{stat}} = 2.4s$, with the asymptotic nonuniformity about a factor of 5 larger.

Figure 69.3

Schematic illustrating the limitation imposed by the spatial-filter pinholes on the amount of spectral dispersion that can propagate through the laser chain. Since the minimum separation between modes in the target plane is $1/2$ the speckle size, the pinholes provide an upper limit on the number of modes available for smoothing.

For the current OMEGA pinhole size ($s=15$), $S_{\max}=76\text{ }\mu\text{m}$; thus, the bandwidth and grating dispersion are chosen such that rays from each phase-plate element are deflected $\pm 38\text{ }\mu\text{m}$ during an SSD modulation cycle. This is small in comparison with a typical target diameter ($\sim 1\text{ mm}$); the performance of SSD is thus limited not by the finite target size but by the selection of pinholes in the laser system.

In addition to the asymptotic level of uniformity, the rate of smoothing is of crucial importance. Smoothing must occur before the target can significantly respond to the laser nonuniformity. A rough estimate for the rate of smoothing provided by SSD can be obtained from the following argument. The rms nonuniformity, averaged over time T , will decrease with the number of noninterfering spectral components N_T as $1/\sqrt{N_T}$. (The time-averaged nonuniformity will continue to decrease until $N_T = N_{\text{stat}}$, and then it will asymptote.) The smallest frequency difference $\delta\nu$ for which interference is negligible is roughly $\delta\nu = 1/T$. For a bandwidth $\Delta\nu$, one has $N_T = \Delta\nu/\delta\nu = \Delta\nu T$. Thus, the rms nonuniformity decreases as $1/\sqrt{\Delta\nu T}$. As an example, with $\Delta\nu = 300\text{ GHz}$ (equivalent to an IR bandwidth of $3\text{ }\text{\AA}$) and $T = 1\text{ ns}$, one finds $N_T \approx 300$. This value of N_T is less than N_{stat} for a spectral separation of 15 times diffraction limit, so that smoothing will continue beyond 1 ns (but such large smoothing times might be too long to affect target performance).

This rate of smoothing is expected to be fast enough to perform the high-compression experiments planned for OMEGA. When combined with polarization shifters and beam overlap, the resulting nonuniformity will be in the range of 1% to 2% with a smoothing time of $\sim 500\text{ ps}$. Higher levels of uniformity could be achieved with the development of new technologies for tripling larger bandwidths. New tripling crystals with a larger bandwidth acceptance are under investigation,

as are improved tripling configurations using existing materials. One option under consideration is to vary the bandwidth in time. At early times, when the intensity is low and high irradiation uniformity is critical, the bandwidth would be large. At these intensities the bandwidth acceptance of the tripling crystals is larger, thus maintaining efficiency. Near the peak of the pulse where high tripling efficiency is crucial, but where laser uniformity can be relaxed because of the smoothing characteristics of the plasma atmosphere that has formed around the target, the bandwidth can be reduced.

2-D SSD Formalism and Results

The principal components of 2-D SSD are shown schematically in Fig. 69.4. With this configuration the bandwidth imposed by the two modulators will be dispersed in two perpendicular directions.

The effect of this configuration on the laser's electric field can be determined approximately from the following treatment. The electric field of the laser entering the first diffraction grating can be written as

$$E(t) = E_0(t)e^{i\omega t}, \quad (4)$$

where the spatial dependence has been suppressed, and the pulse-shape dependence is contained in $E_0(t)$. The first grating in Fig. 69.4 will introduce a time shear across the beam in the x direction, which is equivalent to spectral dispersion. (The directions of dispersion will be referred to as x or y and are perpendicular to the direction of propagation. The change in propagation direction produced by each grating is not shown in Fig. 69.4 and is not relevant to the present discussion.) With the time shear, the electric field becomes

$$E_1(t, x) = E_0(t - \beta x)e^{i\omega(t - \beta x)}, \quad (5)$$

where β is related to the grating dispersion ($\Delta\theta/\Delta\lambda$) by

$$\beta = \frac{2\pi}{\omega} \left(\frac{\Delta\theta}{\Delta\lambda} \right). \quad (6)$$

The quantity $D\beta$ (where D is the beam diameter) remains invariant throughout the laser chain as the beam diameter changes size; its value is the time delay across the beam. For parameters at the end of the IR portion of the laser ($\Delta\theta/\Delta\lambda = 31\text{ }\mu\text{rad}/\text{\AA}$, $D \approx 30\text{ cm}$, $\omega \approx 1.8 \times 10^{15}\text{ s}^{-1}$), the time delay $D\beta$ is about 300 ps.

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Figure 69.4

Schematic of 2-D SSD, illustrating how the bandwidths from the two modulators are dispersed in perpendicular directions.

The first electro-optic modulator introduces sinusoidal phase-modulated bandwidth to E_1 , with amplitude δ_1 and angular frequency ω_1 . The resulting electric field is

$$E_2(t, x) = E_0(t - \beta x) e^{i[\omega(t - \beta x) + \delta_1 \sin(\omega_1 t)]}. \quad (7)$$

The second grating reverses the time shear of the first and disperses the bandwidth in the x direction:

$$E_3(t, x) = E_2(t + \beta x, x). \quad (8)$$

The third grating introduces a time shear in the y direction, which adds spectral dispersion in that direction. The same time shear, given by β , is used in both x and y directions, although this is not a requirement:

$$E_4(t, x, y) = E_3(t - \beta y, x). \quad (9)$$

Gratings 2 and 3 can be replaced by a single grating oriented at 45° to the gratings shown. The second electro-optic modulator introduces additional bandwidth with parameters δ_2 and ω_2 . At this point the electric field is

$$E_5(t, x, y) = E_0(t - \beta y) \cdot \exp i[\omega(t - \beta y) + \delta_1 \sin \omega_1(t + \beta x - \beta y) + \delta_2 \sin(\omega_2 t)]. \quad (10)$$

(There can be an arbitrary phase difference between the two modulators, but this does not affect uniformity on target and has not been included here.) Finally, the fourth grating reverses the effect that the third grating had on the beam due to the bandwidth imposed by the first modulator, and introduces y dispersion (and a time shear) to the bandwidth from the second modulator. The resulting electric field from this configuration has the bandwidth from the first modulator dispersed only in the x direction, the bandwidth from the second modulator dispersed only in the y direction, and no time shear across the beam:

$$E = E_0(t) \cdot \exp i[\omega t + \delta_1 \sin \omega_1(t + \beta x) + \delta_2 \sin \omega_2(t + \beta y)]. \quad (11)$$

This spectrally dispersed light propagates through the laser chain, through the frequency-tripling crystals and phase plates, and through the focus lenses onto the target. For the bandwidths considered here, the main effect of frequency tripling is that the modulation amplitudes δ_1 and δ_2 are each tripled, as is the laser angular frequency ω . Frequencies at the extremes of the bandwidth will triple less efficiently than those near the center. For current frequency-tripling crystals, the difference in efficiency for bandwidths below $\sim 4 \text{ \AA}$ does not significantly effect the uniformity.

The electric field on target (in the focal plane) is

$$E_F = E_0(t) e^{i3\omega t} \sum_{mn} J_m(3\delta_1) J_n(3\delta_2) \times \text{sinc}(m\gamma_1 + q) \cdot \text{sinc}(n\gamma_2 + p) \times e^{i(m\omega_1 + n\omega_2)t} \times \sum_{KL} e^{-i[2(m\gamma_1 + q)L + 2(n\gamma_2 + p)K + \phi_{KL}]} \quad (12)$$

This expression includes the effect of the DPP, frequency tripling, and spectral dispersion. Here the notation of Ref. 1 has been used with a generalization to two-dimensional dispersion. For simplicity, the results are written in terms of a two-level phase plate rather than the more general DPP,² which is actually used on OMEGA. The main effect of the more general phase plate is to replace the "sinc" envelope shape by a more general function, but the speckle statistics are very similar. The variables are defined as follows: p and q are dimensionless variables related to y and x by the factor $k_3\Delta/2F$, where k_3 is the wave number of the frequency-tripled fundamental, Δ is the distance between phase-plate elements, F is the focal length, and $\gamma_{1,2} = \omega_{1,2}\beta\Delta/2$. The Bessel functions J_m and J_n are the amplitudes from a Fourier decomposition of the sinusoidal phase modulation. Formally the m - n sum extends to $\pm\infty$, but the contributions from $|m| > 3\delta_1$ and $|n| > 3\delta_2$ are very small. Thus, to a good approximation, the largest spectral modes of interest are given by $m = M \equiv 3\delta_1$ and $n = N \equiv 3\delta_2$. The K - L sum is the factor that describes the phase-plate speckle. The sum is over all phase-plate elements, each having a phase ϕ_{KL} , which is either 0 or π . Note that each spectral component (m, n) has exactly the same speckle structure but is shifted by $m\gamma_1$ in q and by $n\gamma_2$ in p .

These variables have the following physical significance. The phase-plate element size Δ is generally chosen so that the distance between the zeros of the sinc function is slightly larger than the largest target that will be irradiated (to assure good uniformity). Thus, $p, q = \pi$ is characteristic of the target radius. The total spectrally induced spreads in the two directions, as fractions of the target diameter, are approximately $M\gamma_1/\pi$ and $N\gamma_2/\pi$. The total bandwidths in the two directions are $\Delta\nu_1 = M\omega_1/\pi$ and $\Delta\nu_2 = N\omega_2/\pi$. The number of times that the phase modulation repeats across the beam is $\gamma_{1,2}K_{\max}/\pi$, where $K_{\max}$ is the number of phase-plate elements in one direction.

The laser intensity I in the focal plane is given by the square of $|E_F|$:

$$I = I_0(t) \sum_{\substack{mm' \\ nn'}} J_m(3\delta_1) J_{m'}(3\delta_1) J_n(3\delta_2) J_{n'}(3\delta_2) \cdot e^{i(m\omega_1 + n\omega_2 - m'\omega_1 - n'\omega_2)t} \times \text{sinc}(m\gamma + q) \text{sinc}(m'\gamma + q) \times \text{sinc}(n\gamma + p) \text{sinc}(n'\gamma + p) \times \sum_{\substack{KK' \\ LL'}} e^{i2(L'-L)q + i2(K'-K)p} \times \left\{ e^{i2\gamma(m'L' - mL + n'K' - nK) + i(\phi_{KL'} - \phi_{KL})} \right\} \quad (13)$$

The final summation has been written in the form of a Fourier decomposition of the speckle nonuniformity with the term in brackets being the Fourier coefficients.

The time-averaged uniformity at time t in the focal plane is

$$\langle I(t) \rangle_T = \frac{1}{T} \int_{t-T/2}^{t+T/2} I(t') dt' \quad (14)$$

for an averaging time T . To examine the smoothing effect of 2-D SSD, it is instructive to evaluate the asymptotic limit of $\langle I \rangle_T$ as $T \rightarrow \infty$. To simplify the result, consider the special case for which (1) $I_0(t)$ is constant; (2) the spectral shifts are sufficiently small that the sinc envelope is not modified; and (3) the modulation frequencies are incommensurate, i.e., $m\omega_1 \neq n\omega_2$ for all integers m and n . The resulting asymptotic uniformity can be written in the following form:

$$\langle I \rangle_\infty = I_{\text{env}}(p, q) \cdot \left[1 + \sum_{k\ell \neq 0} e^{i(2kp + 2\ell q)} C_{k\ell} \cdot R_{k\ell} \right], \quad (15)$$

where I_{env} is the diffraction-limited phase-plate envelope, C is the spatial autocorrelation function for the phase plate, R is the reduction in nonuniformity produced by SSD, and the summation is over the distance between phase-plate elements (i.e., $k = K' - K$ and $\ell = L' - L$). These terms are defined as follows:

$$I_{\text{env}} = I_0 \text{sinc}^2(p) \text{sinc}^2(q) N^2, \quad (16)$$

where N^2 is the number of phase-plate elements,

$$C_{k\ell} = \sum_{KL} e^{i(\phi_{K+k, L+\ell} - \phi_{KL})}, \quad (17)$$

and

$$R_{k\ell} = \sum_{mn} J_m^2(3\delta_1) J_n^2(3\delta_2) e^{i(2\gamma_1 m \ell + 2\gamma_2 n k)}. \quad (18)$$

Using the Bessel function identity

$$\sum_n J_n(3\delta) J_{n+v}(3\delta) e^{in\gamma\ell} = J_v(w) e^{iv\theta}, \quad (19)$$

where $w = 6\delta |\sin \gamma \ell|$ and $\sin \theta = (3\delta \sin \gamma \ell)/w$, the SSD reduction factor becomes

$$R_{k\ell} = J_0(6\delta_1 \sin \gamma_1 \ell) \cdot J_0(6\delta_2 \sin \gamma_2 k). \quad (20)$$

In the form of Eq. (15), the nonuniformity structure has been Fourier decomposed in terms of the variables p and q , with the difference between phase-plate elements k and ℓ acting as dimensionless wave numbers. The rms fluctuation is determined by the square of the Fourier coefficients:

$$\sigma_{rms} = \left[\sum_{k\ell} |C_{k\ell} \cdot R_{k\ell}|^2 \right]^{1/2}. \quad (21)$$

The phase-plate result (without SSD) is recovered for $R_{k\ell} = 1$ (for all k and ℓ), in which case the rms nonuniformity is 100%.

The SSD reduction factor depends on the spatial wavelength of the nonuniformity through k and ℓ : For $\gamma_1 \ell < (6\delta_1)^{-1}$ and $\gamma_2 k < (6\delta_2)^{-1}$ the factor $R_{k\ell}$ is approximately 1. When either k or ℓ is small, the factor reduces to the 1-D SSD result. This occurs when the wave vector for the nonuniformity points toward one of the directions of dispersion. Maximum smoothing from 2-D SSD occurs when both k and ℓ are large. Fig-

ure 69.5 shows the reduction factor $|R_{k\ell}|$ using the current OMEGA values ($\delta_1 = 5.1$, $\delta_2 = 4.6$) for the cases when one wave number is kept small ($\ell = 0$) and for the case when the wave vector is at 45° to the direction of dispersion ($\ell = k$). The improved smoothing in the second case, corresponding to two-dimensional smoothing of the nonuniformity in both directions, is apparent.

The condition that there is no significant smoothing for $6\delta_1 \gamma_1 \ell < 1$ has a simple physical interpretation. The factor $6\delta_1 \gamma_1$ is the total spectrally induced shift $S_{\max}$ in the target plane in units of q . The parameter ℓ is related to the wavelength of nonuniformity $\Delta\lambda$ by $\Delta\lambda = \pi/\ell$ (again in units of q). Thus, the largest wavelength of nonuniformity that will be smoothed by SSD is approximately $\pi S_{\max}$. In other words, the spectral shift must be a significant portion of the nonuniformity wavelength for smoothing to occur.

The amount that the speckle pattern is shifted can be doubled (in one direction) by means of a polarization shifter.^{3,4} This is illustrated in Fig. 69.6 for the shifter currently under investigation at LLE, a birefringent wedge of KDP placed after the frequency-tripling crystals. The birefringence of KDP separates the laser beam into two orthogonal polarizations, which are deflected by the wedge through slightly different angles. (Alternatively, a liquid-crystal wedge could be used. Depending on the type of liquid crystal, the orthogonal polarizations could be linear or circular.) The two speckle patterns produced on target are spatially displaced, thus doubling the number of independent speckle patterns. The most effective way to use this technique is to disperse the polarization over a distance larger than the spectral dispersion produced by SSD. In this way, the polarization shift smoothes out modes of nonuniformity that are not smoothed by SSD, and overlap between the two techniques is avoided.⁵ The reduction

Figure 69.5

The absolute value of the 2-D SSD reduction factor as a function of nonuniformity wave vector (or equivalently the separation k between phase-plate elements). (a) The wave vector is along one of the dispersion directions. The reduction is the same as for 1-D SSD. (b) The wave vector is at 45° , resulting in the maximum 2-D reduction.

Figure 69.6

Principle of the (birefringent) polarization wedge. The laser beam is split equally into two orthogonal polarizations ("o" and "e"), which are deflected by the wedge through slightly different angles. This results in two spatially displaced speckle patterns in the target plane with orthogonal polarizations, which add in intensity rather than electric field, providing an instantaneous reduction of $1/\sqrt{2}$ in the rms nonuniformity.

in nonuniformity resulting from the polarization shifter is instantaneous because there is no interference between the two polarizations.

The predicted improvement in uniformity produced by 2-D SSD on OMEGA is shown in Fig. 69.7 as a function of the smoothing time. The results are for multiple overlapping beams on a spherical target and thus include the smoothing effect of overlapping beams. Two cases are presented. One shows the results for the current implementation of 2-D SSD in which $1.5 \text{ \AA}$ of bandwidth is dispersed in each direction. The second case shows the result of doubling the bandwidth in one direction and adding polarization dispersion in the second. This case achieves the uniformity goal of OMEGA, namely reaching the 1%–2% rms level within a smoothing time of less than 500 ps.

Implementation of 2-D SSD

2-D SSD was implemented on OMEGA in January 1996. An aggressive optical design program showed that 2-D SSD

Figure 69.7

Reduction in rms nonuniformity as a function of smoothing time for multiple beam overlap on a spherical target, using the 60-beam OMEGA geometry. Spherical harmonic modes up to $\ell = 500$ have been considered with no additional smoothing assumed in the plasma atmosphere around the target. The lower curve corresponds to a higher bandwidth in one direction and the inclusion of polarization shifters.

could be implemented over roughly the same time interval that was originally planned for 1-D SSD. However, it was important to adopt a conservative strategy with regard to bandwidth during the time that the operating characteristics of the new laser and SSD were being studied, because spectrally dispersed bandwidth can introduce intensity fluctuations that can damage laser optical components. As such, the initial implementation of 2-D SSD used bandwidths of only 1.5 Å and 0.75 Å. The initial results presented in this section were generated with these bandwidths. The bandwidth has recently been increased to 1.5 Å in each direction. To accommodate the increased dispersion of the beam, it has been necessary to enlarge one spatial-filter pinhole in the driver and to enlarge the second SSD modulator crystal. Future plans call for an increase in pinhole diameters later in the laser chain to allow the propagation of up to 3 Å of bandwidth.

A far-field image of the beam, after the final set of gratings, is shown in Fig. 69.8. Although the individual spectral components can not be seen in this figure, the two-dimensional dispersion is clearly evident. The extremes of the spectrum are the most intense portions for sinusoidal phase modulation. This is seen in the corners of the figure. The two directions of dispersion are not exactly orthogonal due to a 6° misalignment of the periscope that takes the beam out of the pulse generation room (PGR), but this has no effect on the irradiation uniformity on target. The amount of dispersion in each direction is proportional to the bandwidth from each modulator.

Images of the final beam profile,⁶ in an equivalent target plane (ETP), are shown in Fig. 69.9. These images are for a single beam profile and therefore take no account of the additional smoothing achieved by beam overlap. The first

Figure 69.8

A far-field image of the beam (after the final diffraction gratings) showing that the two-dimensional dispersion of the beam is proportional to the bandwidth applied by each modulator.

image [Fig. 69.9(a)] shows the frequency-tripled beam without a phase plate or SSD. The intensity nonuniformity is the result of phase aberrations that have accumulated throughout the laser chain. Figure 69.9(b) shows the improvement produced by a phase plate. A well-defined intensity envelope has been established, but superposed on the envelope is highly modulated speckle. Figure 69.9(c) shows smoothing of the speckle by 1-D SSD, for which the bandwidth is turned on in one modulator and off in the other. Both combinations are shown. The "stripes" in the images show the direction of spectral dispersion. Nonuniformity perpendicular to the dispersion is not smoothed. Finally, the last image [Fig. 69.9(d)] shows the improved smoothing produced by two-dimensional dispersion. Note that the "stripes" have now been eliminated.

These images were time integrated over the laser pulse. The uniformity achieved is characteristic of an SSD smoothing time roughly equal to the pulse width, which was ~1 ns for this experiment. The rms fluctuations of the intensity around smooth envelopes are listed in Table 69.I, which gives both the measured and theoretically predicted values. A large part of the improved uniformity for 2-D SSD compared with 1-D SSD (a factor of ~2) is related to the smoothing duration: 1-D SSD reaches an asymptotic level of uniformity after ~300 ps; 2-D SSD continues to smooth throughout the entire time of the pulse (1 ns). There is an increased bandwidth for 2-D SSD (due to contributions from both modulators), but this has a much smaller effect on the improved uniformity than the increased smoothing duration, for this example.

Very recently a prototype KDP polarization wedge has been tested on OMEGA. One-dimensional lineouts through equivalent-target-plane (ETP) images (Fig. 69.10) show that the predicted $\sqrt{2}$ uniformity improvement is indeed obtained.

Table 69.I: rms nonuniformity for the single-beam images in Fig. 69.9, compared with the theoretically predicted values (when averaged over the 1-ns pulse).

Image	Bandwidth (Å)	Measured rms	Calculated rms
(b)	0	0.96	0.98
(c)	0.75, 0	0.27	0.28
(c)	0, 1.5	0.24	0.21
(d)	0.75, 1.5	0.12	0.11

Figure 69.9

Equivalent-target-plane images, integrated over the ~ 1 -ns pulse width, of a single OMEGA beam with four levels of smoothing: (a) unsmoothed, frequency-tripled; (b) phase plate, no bandwidth; (c) bandwidth in only one modulator; (d) bandwidth in both modulators.

Figure 69.10

One-dimensional lineouts of the ETP images on OMEGA produced (a) without a polarization wedge and (b) with a polarization wedge. The measured rms nonuniformities of $\sigma_{\text{rms}} = 98\%$ and 68% , respectively, demonstrate the predicted $\sqrt{2}$ uniformity improvement made possible by the wedge.

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Areal Density Measurement of Laser Targets Using Absorption Lines

Absorption lines due to $1s-2p$ transitions in titanium ions are predicted to be observed in the implosion of titanium-doped shells. The measured absorption of these lines can be used to determine the peak areal density $\rho\Delta r$ of the doped layer and hence of the total compressed shell. The absorption lines are studied by solving the radiation transport equation using opacity tables and hydrodynamic simulations. The absorption is a function not only of $\rho\Delta r$ but also of the density and temperature of the absorbing layer. However, it is shown that the areal density can be estimated with reasonable accuracy by using the measured intensity of absorption and its distribution over the various absorption lines. The considerations affecting the choice of doping parameters are discussed, as well as the effect of integrating the measured spectrum over time and target volume.

Recently published work¹ showed that a titanium-doped layer within a polymer shell of a laser-imploded target can yield information on core-shell mixing. We show here that the absorption lines predicted to be produced by such a doping provide a signature of shell areal density ($\rho\Delta r$) at peak compression. This areal density is an important parameter characterizing the implosion performance, which in turn is determined by target instability and mixing.

In the work on mixing diagnostics,¹ the doped-layer location within the shell was removed from the fuel-shell interface so that, in the absence of mixing, only titanium absorption lines would be observed, but mixing could cause titanium material to move into inner target regions and emit titanium lines. Even for the assumed level of mixing, *the absorption lines were essentially unchanged from the no-mixing case*. Thus, such a doped layer can provide information on the overall $\rho\Delta r$ of the compressed (or overdense) shell, even if its inner part undergoes mixing. The doped layer is placed in the target far from the interface but still close enough not to be ablated away.

Measurement of the absorption spectrum yields the areal density of only the doped layer.² This by itself is of interest because the comparison of measured and predicted $\rho\Delta r$ for

the doped layer is a measure of target performance. Additionally, the total areal density of the compressed shell can be deduced if the thickness of target layer that is ablated away is measured (e.g., by charge collectors).

The absorption lines are studied by solving the radiation transport equation using OPLIB³ opacity tables and one-dimensional LILAC⁴ hydrodynamic simulations. The absorption is a function not only of $\rho\Delta r$ but also of the density and temperature of the absorbing layer. However, it is shown that the areal density can be estimated with reasonable accuracy by using the measured intensity of absorption and its distribution over the various absorption lines. The initial $\rho\Delta r$ of the titanium doping must be small enough, or the effect of local emission in the absorption region will lead to underestimating the $\rho\Delta r$. It is further shown that integrating the measured spectrum over the target volume causes only a small error, while integrating over time causes the $\rho\Delta r$ to be underestimated by about a factor of 2.

Modeling a Test Case

We calculate the expected x-ray spectrum of a particular simulated target implosion on the OMEGA laser. For this test case, LILAC results were used for the expected temperature (T) and density (ρ) profiles, and a post-processor⁵ code was used to calculate the emission and radiation transport through the target. To simulate the emergent absorption and emission spectrum, multigroup opacity tables were generated using the OPLIB opacity library.³

The target is a polymer shell of 940- μm diam and 30- μm thickness, filled with 80 atm DT gas. A layer doped with 1% titanium (by atom number) is embedded in a CH polymer shell; the doped layer is placed 2.4 μm from the interface and its thickness is varied. As explained in earlier work,¹ such a doped layer shows significant absorption at the wavelengths of titanium lines, while showing little effect on overall target behavior. The laser pulse is trapezoidal, rising linearly over a 0.1-ns period to 13.5 TW, then remaining constant for 2.2 ns, before dropping linearly over a 0.1-ns period. LILAC hydrody-

namic simulations show¹ that the shell compresses to a mean radius of $\sim 50 \mu\text{m}$ and thickness of $\sim 30 \mu\text{m}$, with a density in the range of ~ 10 to 50 g/cm^3 , corresponding to a $\rho\Delta r$ value of $\sim 90 \text{ mg/cm}^2$. The electron temperature in the shell ranges from ~ 800 (at the shell-fuel interface) to $\sim 80 \text{ eV}$ (at the peak of the shell density); the line absorption occurs within the colder, outer part of this compressed shell. The profiles at maximum compression are shown in Fig. 69.11.

E6738

Figure 69.11

The density and temperature profiles predicted by the *LILAC* hydrodynamics simulation code at peak compression.

Figure 69.12 shows examples of the predicted spectrum at two instances during the implosion, showing titanium absorption lines. The time 2.9 ns corresponds to peak compression. These lines correspond to transitions of the type $1s-2p$ in titanium ions of an incomplete L shell: Ti^{+13} to Ti^{+20} . The designation Li in Fig. 69.12 stands for lithium-like titanium ion, and likewise for the other designations. Each peak contains several lines that are unresolved mostly because of broadening due to the finite source size. For example, the peak marked C (carbon-like) consists of 35 transitions of the type $1s^2 2s^2 2p^2 - 1s 2s^2 2p^3$. On approaching peak compression, the intensity of the continuum radiation is seen to rise sharply and the absorption-line manifold is seen to shift to higher ionization states. This shift, caused by the increase in shell temperature as the shell becomes more compressed, is discussed further below.

The absorption lines are formed when radiation emitted by the compressed core traverses colder shell layers. By the definition of the opacity k , the intensity within an absorption

E8120

Figure 69.12

The predicted spectrum emitted by the test case (1%-titanium-doped layer) at two times during the implosion; $E_0 = 7 \times 10^{15} \text{ keV}/(\text{keV ns } \Omega)$. The time 2.9 ns corresponds to peak compression. The designation Li stands for the lithium-like titanium ion, and likewise for the other designations. The emission is integrated over the target volume.

line at an energy E is given by

$$I(E) = I_0(E) \exp\left[-\int k(E)\rho dr\right], \quad (1)$$

where I_0 is the core-emitted intensity and $k(E)$ is the opacity per unit areal density. Equation (1) assumes that the local emission within the absorption region is neglected; this point is discussed further below. From Eq. (1) it follows that the $\rho\Delta r$ of the absorbing layer can be deduced from the measured spectrum through the relationship $\rho\Delta r = \ln[I_0(E)/I(E)]/k(E)$. However, $k(E)$ depends also on the (unknown) temperature and density in the absorption layer. Note that the density dependence of the opacity is in addition to the explicit dependence of the attenuation on $\rho\Delta r$ [Eq. (1)]. Next we examine these dependencies and show how Eq. (1) can be used to estimate the $\rho\Delta r$ of the absorbing layer, even without an exact knowledge of the temperature and the density.

Areal Density Determination Using Absorption Lines

Figures 69.13 and 69.14 show examples from OPLIB³ of the opacity spectrum for 1% titanium in CH, as a function of temperature and density. As seen, 1%-titanium doping provides an ample contrast ratio between line absorption by titanium ions and the nonresonant absorption by both titanium

and CH. As the temperature increases (Fig. 69.13), the absorption-line manifold shifts to higher photon energies, or higher ionizations. The shift seen in Fig. 69.12 is similarly due to an increase in shell temperature on approaching peak compression. On the other hand, Fig. 69.14 shows that for the same temperature, a higher density causes the manifold to shift to lower ionizations. This is because the three-body (or collisional) recombination increases with increasing density faster than the two-body collision processes. The three-body recombination is important only at or near LTE (local thermodynamic equilibrium),⁶ which is the atomic model used to generate the OPLIB tables. However, it can be shown⁷ that for the ions and temperatures under consideration, the LTE approximation for the distribution of ion populations holds for densities (of the mostly polymer material) higher than $\sim 6 \text{ g/cm}^3$; much higher shell densities than this are predicted for the test case. From Fig. 69.13 a doubling of the temperature leads to a 0.05-keV shift in the peak absorption. A similar shift in peak absorption requires, from Fig. 69.14, a factor-of-10 increase in density. Thus the peak positions are more sensitive to changes in temperature than in density. Figures 69.13 and 69.14 clearly show that the absorption at any given photon energy depends on both the temperature and the density; consequently, Eq. (1) cannot be used in a simple way to derive $\rho\Delta r$. In addition it should be noted that (1) changes in temperature and density cause primarily an energy shift in the absorption lines and

(2) the total absorption-line manifold has only a weak dependence on temperature and density. Thus we choose the integral $\int \ln[I_0(E)/I(E)] dE$ over the whole absorption-line manifold as the experimental signature. We refer to the integral $\int \ln(I_0/I) dE$ as the area within the absorption lines [if film density is proportional to $\ln(I)$, this indeed is the area under the absorption lines on film]. From Eq. (1) we obtain

$$\int \ln[I_0(E)/I(E)] dE = \rho \Delta r \int k(E) dE. \quad (2)$$

The opacity integral on the right can be obtained from opacity spectra such as those in Figs. 69.13 and 69.14. It turns out that even after integrating the opacity over the absorption-line manifold, the opacity still depends appreciably on temperature and density. For example, the opacity varies from 35 to 85 $\text{cm}^2 \text{ keV/g}$ when the temperature is varied from 0.3 to 10 keV and the density is varied from 1 to 50 g/cm^3 . However, in addition to the integral over the spectrum we can also make use of the measured distribution among the absorption peaks. More specifically, the peak of strongest absorption provides an additional signature that also depends on both temperature and density and helps to narrow the range of opacity values.

To make use of this additional information, we calculated the integrated opacity for temperatures in the range of 0.3 to 1.0 keV and densities in the range of 1 to 50 g/cm^3 . Fig-

Figure 69.13

The opacity of 1%-titanium-doped CH at a density of 3.2 g/cm^3 (from OPLIB³), showing a shift to higher photon energies (higher ionizations) with increasing temperature. The absorption lines are due to $1s-2p$ transitions in titanium ions with increasing number of L -shell vacancies.

Figure 69.14

The opacity of 1%-titanium-doped CH at a temperature of 300 eV, showing a shift to lower photon energies (lower ionizations) with increasing density.

E8187

Figure 69.15

Area under the opacity-spectrum peaks, such as those in Figs. 69.13 and 69.14. For temperatures in the range of 300 to 1000 eV and densities in the range of 1 to 50 g/cm³, the points are grouped according to the strongest absorption peak in the opacity spectrum (Li stands for cases where the lithium-like peak is the strongest, and likewise for the other designations).

Figure 69.15 shows the results, arranged according to the strongest opacity peak (marked Li for cases where the lithium-like peak is the strongest, etc.). The four charge states (C-like to Li-like) cover the range likely to be encountered in target implosions. Figure 69.15 shows that knowing the peak of maximum absorption (without knowing the temperature or the density) narrows down the range of integrated opacity from $\pm 42\%$ to within $\pm 15\%$.

To apply Fig. 69.15 to an experimental result, we first note the peak of maximum absorption in the spectrum. For example, if the Be-like peak dominates, we obtain from Fig. 69.15 the value 52 for the integrated opacity $\int k(E)dE$, without having to know the temperature or the density. Substituting this value in Eq. (2) produces a relation between the $\rho\Delta r$ of the doped layer and the measured area within the absorption lines (the integral on the left side of the equation). The error in determining $\rho\Delta r$ based on Fig. 69.15 is $\pm 15\%$; other sources of error are discussed below.

Using the area under the absorption lines rather than the intensity profiles obviates the need to know the line profiles (e.g., due to Doppler and Stark effects). However, the measured and calculated intensity profiles should actually be quite similar because the absorption profile for each ion specie is comprised of many closely spaced components. In the experiment these components are smeared by the source-size

broadening, whereas in the OPLIB³ calculations they are smeared by the energy bin width, chosen to be similar to the source-size broadening. The final absorption profile for each ion specie is the envelope over these components and depends mostly on their relative intensity and separation.

Application of the Method to a Simulated Implosion

To test the method of determining $\rho\Delta r$ by the spectrum of absorption lines we apply it to a simulated spectrum from the implosion of the test case and compare the resulting $\rho\Delta r$ to the actual $\rho\Delta r$ of the profiles used in the calculation of the spectrum. This procedure simulates the application of the method to an experimentally observed spectrum.

Figure 69.16 shows two examples of the calculated emergent spectrum for the test case. For very high opacities, the attenuation in the emergent lines cannot be arbitrarily high because local emission within the absorption region fills in the absorption dips. Therefore, the doped layer must be sufficiently thin, or the $\rho\Delta r$ will be underestimated. Similarly, the depth of the observed absorption lines is limited by the finite dynamic range of the film (or other detector). To demonstrate the effect of self-absorption, we show in Fig. 69.16 a comparison of the spectra emitted from two identical titanium-doped targets except for the thickness of the doped layer, which was

E8123

Figure 69.16

Comparison of the spectrum emitted from two identical 1%-titanium-doped targets except for the thickness of the doped layer, which was 0.4 and 1.2 μm , respectively. Both spectra refer to the time of peak compression and are integrated over the target volume; $E_0 = 7 \times 10^{15}$ keV/(keV ns Ω). The near coincidence of the two spectra is caused by the local emission in the absorption region.

0.4 and 1.2 μm , respectively. We refer to these two target simulations as the "thin case" and the "thick case." The separation of the outer surface of the doped layer from the shell-fill interface was 2.4 μm in both cases (the thick case's inner boundary was closer to the shell-fill interface than the thin case's). Both spectra refer to the time of peak compression and are integrated over the target volume. The similarity of the two spectra is clear evidence of local emission in the absorption region. Thus, at the wavelengths of highest absorption, only radiation emitted on the outer surface of the doped layer can emerge, and its intensity is independent of the thickness of the doped layer. On the other hand, for the weak absorption peaks [where the attenuation factor $\ln(I_0/I)$ is smaller than ~ 1] we would expect the absorption depths to be proportional to the doped-layer thickness. The reason for the weaker changes evident in Fig. 69.16 is that the thin doped layer happens to sample the highest-absorbing part (higher ρ , lower T) of the doped region. Thus, a too-thick doped layer gives rise to difficulties due to both local emission and gradients in plasma parameters.

We now apply the method of determining the areal density to the two spectra in Fig. 69.16. The area $\int \ln(I_0/I) dE$ in Fig. 69.16 is equal to 0.17 keV for the thin case and 0.21 keV for the thick case. The strongest absorption peak is at 4.65 keV and corresponds to Be-like titanium. From Fig. 69.15 we use the average integrated opacity for the Be-like titanium, which has the value 52 $\text{cm}^2 \text{ keV/g}$ (with an error of $\pm 15\%$). Substituting these values into Eq. (2) we finally obtain $\rho\Delta r = 3.2 \text{ mg/cm}^2$ for the thin case and $\rho\Delta r = 4.0 \text{ mg/cm}^2$ for the thick case. This is compared with *LILAC*'s $\rho\Delta r$ values at peak compression: $\rho\Delta r = 3.1 \text{ mg/cm}^2$ for the thin case and $\rho\Delta r = 13.5 \text{ mg/cm}^2$ for the thick case. Thus, the method works well for the thin case but significantly underestimates the $\rho\Delta r$ in the thick case because of local emission within the absorption region.

To obtain accurate $\rho\Delta r$ measurements from this technique the target doping must be such that $k\rho\Delta r < 1$, which can be achieved in two ways: The choice of target doping can be based on simulations, or one can start with high- $\rho\Delta r$ doping and repeat the experiment with successively lower doping levels until the measured absorption is seen to decrease with decreasing doping. The latter procedure is also desirable for the following reason: if the achieved $\rho\Delta r$ is smaller than predicted, the absorption lines may not be observed at all. Thus, starting with a thicker doped layer ensures the observation of absorption lines even in such a case. In the test case

studied here, a 0.4- μm -thick layer doped at 1% was shown to be an appropriate final choice for doping.

The emergent spectrum from the simulated target has been shown in earlier publications.¹ However, that spectrum was computed for a spatially and temporally resolved measurement (axial view at peak compression). We show here the effect of integrating the calculated emergent spectrum over space and time. First, Fig. 69.17 shows a comparison of an axial-view spectrum (per unit area) with the space-integrated spectrum. The spectral intensity refers to the flux per unit area in units of $I_0 = 7 \times 10^{20} \text{ keV}/(\text{keV ns cm}^2 \Omega)$, whereas the emission refers to the flux from the whole target in units of $E_0 = 7 \times 10^{15} \text{ keV}/(\text{keV ns } \Omega)$. The ratio between the two curves is essentially equal to the inverse of the cross section of the emitting core (giving in Fig. 69.17 a core diameter of $\sim 80 \mu\text{m}$). The two spectra have very similar shapes, and the spatially integrated spectrum has an area $\int \ln(I_0/I) dE$ that is only $\sim 15\%$ smaller than that of the spatially resolved spectrum. This is not surprising since any ray from the core in the direction of observation traverses essentially the same shell thickness.

Figure 69.18 shows a comparison between the emergent spectrum at the time of peak compression (2.9 ns) and the time-integrated spectrum; both are integrated over the target volume. Note that the peak in absorption for the time-integrated spectra is not as clear cut as the peak for the time-resolved spectra. As

Figure 69.17

Comparison of simulated axial-view spectrum (per unit area) and space-integrated spectrum from an imploded, 1%-titanium-doped target. $I_0 = 7 \times 10^{20} \text{ keV}/(\text{keV ns cm}^2 \Omega)$; $E_0 = 7 \times 10^{15} \text{ keV}/(\text{keV ns } \Omega)$. The ratio between the two is essentially equal to the inverse of the cross section of the emitting core.

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Figure 69.18

Comparison of simulated peak-compression and time-integrated spectra from an imploded, 1%-titanium-doped target (both are integrated over the target volume). $E_0 = 7 \times 10^{15}$ keV/(keV ns Ω); $F_0 = 2.3 \times 10^{15}$ keV/(keV Ω). The ratio between the two curves is essentially equal to the inverse of the core emission duration in nanoseconds.

the implosion progresses, the titanium dopant experiences an increase in temperature and density, and consequently there is a shift in the absorption peak with time. The time-integrated measurement averages over all these peaks. The emission, as in Fig. 69.17, is in units of $E_0 = 7 \times 10^{15}$ keV/(keV ns Ω) and the fluence is in units of $F_0 = 2.3 \times 10^{15}$ keV/(keV Ω). The ratio between the two curves is essentially equal to the inverse of the core emission duration in nanoseconds (Fig. 69.18 gives a duration of ~ 0.6 ns). The area $\int \ln(I_0/I) dE$ for the time-integrated spectrum is about half that of the peak-emission spectrum and using it would result in an underestimate of the $\rho\Delta r$ by the same factor.

As mentioned earlier, the total $\rho\Delta r$ of the compressed shell, $(\rho\Delta r)_{\text{total}}$, can be deduced from that of the doped layer, $(\rho\Delta r)_{\text{doped}}$, if the total ablated mass is measured (e.g., by charge collectors). A simple geometrical consideration shows that for a *uniformly compressed shell* the ratio

$$(\rho\Delta r)_{\text{doped}} / (\rho\Delta r)_{\text{total}}$$

is proportional to its initial value through a proportionality constant C , which depends on the spherical convergence and the location of the doped layer within the shell; C is smaller than, but close to 1. In the simulated test case the density of the compressed shell was constant to within $\pm 30\%$. For that case, with an initial 30- μm -thick shell and a 0.4- μm -thick doped

layer (the thin case), we deduced above a $\rho\Delta r$ for the doped layer of 3.2 mg/cm². The simulations by *LILAC* show that 16 μm of the shell thickness is unablated (this in an experiment will be deduced from charge collectors), so the total shell areal density at peak compression is estimated as $\rho\Delta r = 3.2 \times (16/0.4) = 128$ mg/cm². Here we assumed that the convergence is unknown and used $C = 1$. This result compares with *LILAC*'s value for peak compression of $\rho\Delta r = 110$ mg/cm². If the convergence is known (for example, from imaging), a better value of C can be used and the discrepancy reduced.

It might be thought that a thicker layer doped at a lower concentration could reduce the problem of relating the measured $\rho\Delta r$ to that of the total shell. However, this is undesirable because enlarging the doped layer results in averaging over the steep temperature gradient in the shell. This introduces an uncertainty in the choice of peak to be used in Fig. 69.15 and thus an additional error in the derived areal density.

Conclusion

An effective method has been described for measuring the areal density of the compressed (or overdense) shell of laser-imploded targets, based on the observation of absorption lines from a titanium-doped layer. Four factors have been shown to affect the precision of such measurement:

- If the initial $\rho\Delta r$ of the doped layer is too high, the problem of local emission in the absorption region will lead to an underestimate of the $\rho\Delta r$ at peak compression. To address this problem, the experiment can be repeated with progressively thinner doped layers (or lower doping concentrations), until the absorption is seen to decrease with decreasing $\rho\Delta r$ of the initial doped layer. This procedure is also desirable to ensure the observation of absorption lines even if the actual compression is smaller than predicted.
- If the spectrum is not streaked in time, the $\rho\Delta r$ will be grossly underestimated (typically by a factor of 2).
- If the spectrum is not spatially resolved, the $\rho\Delta r$ will be slightly underestimated (typically by $\sim 15\%$). Thus, it is more important to resolve the spectrum in time than in space.
- The total ablated mass must first be determined (for example, by employing charge collectors) to be able to relate the measured $\rho\Delta r$ of the doped layer to the total $\rho\Delta r$ of the overdense shell.

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Modeling the Temporal-Pulse-Shape Dynamics of an Actively Stabilized Regenerative Amplifier for OMEGA Pulse-Shaping Applications

Advances in laser-fusion technology indicate that the temporal profile of the laser pulse applied to laser-fusion targets is important for improving the performance of these targets.¹ The OMEGA laser is a 60-beam laser-fusion system capable of producing a total of 30 kJ of ultraviolet (351-nm) energy on target, where the temporal profile of the optical pulse applied to a laser-fusion target can be specified in advance. This is accomplished by a pulse-shaping system that produces an optical pulse with a specific temporal pulse shape at the nanojoule energy level.² This pulse seeds an actively stabilized Nd:YLF regenerative amplifier³ (regen) followed by and wavelength matched (1053 nm) to a series of Nd:glass amplifiers. The beams are then frequency tripled to the third harmonic using KDP nonlinear crystals. To achieve the desired on-target optical pulse shape, the temporal dynamics of the entire OMEGA laser system must be accurately modeled to determine the specific temporal profile of the seed pulse required from the pulse-shaping system at the beginning of the laser. The temporal profile of this low-energy seed pulse, when amplified and frequency tripled by the laser system, will then compensate for the temporal distortions caused by gain saturation in the regen and amplifiers and by the tripling process, and will produce the desired pulse shape on target.

To determine the required temporal profile of the optical pulse at the beginning of the system, all sources of temporal distortions in the system must be understood and compensated for. The temporal distortion due to the frequency-tripling process is modeled with a time-dependent simulation of the appropriate nonlinear equations for this process.⁴ The temporal-pulse distortions in the system's Nd:glass amplifiers are easily modeled with our beam code *RAINBOW*. Modeling the actively stabilized Nd:YLF regen at the beginning of the system is the topic of this article. With the regen model described here, OMEGA's temporal dynamics can now be completely modeled. Pulse distortions in the system can be easily compensated for by proper choice of the seed temporal profile determined from the overall model.

The gain of the OMEGA system from the pulse-shaping modulator to the target is approximately 10^{14} ; a gain of 10^7 in the actively stabilized regen is included in this overall gain. Modeling the actively stabilized regen is complicated by many factors. The regen must be treated as a multipass amplifier, the last few passes of which experience significant gain saturation in the Nd:YLF laser rod. The lifetime of the lower-laser-level manifold in Nd:YLF has been measured to be 21 ns,⁵ and the round-trip time in the cavity is 26 ns. A Frantz-Nodvik-type solution⁶ for the gain in the Nd:YLF medium cannot account for this finite lower-laser-level lifetime and, hence, is inappropriate; the rate equations must be used to describe the single-pass gain in the Nd:YLF medium. Finally, the regen incorporates a feedback mechanism that measures the circulating pulse energy each round-trip.³ When the circulating pulse energy exceeds a threshold ($\sim 25 \mu\text{J}$), a feedback mechanism is activated. The feedback mechanism introduces appropriate losses into the cavity each round-trip in order to stabilize the circulating pulse energy to a fixed but low-energy level. During this prelude stabilization phase the regen is operating with a net gain (round-trip gain/loss) approximately equal to unity and establishes a constant and stable pulse-energy output from the regen. After this prelude stabilization is achieved, the laser can be *Q*-switched by eliminating the feedback losses from the cavity. The regen will then emit a *Q*-switched envelope of pulses. The pulse at the peak of the *Q*-switched envelope is switched out and sent to the OMEGA amplifiers. This stabilized regen produces pulses with a long-term shot-to-shot energy stability of approximately 0.2%, despite the fluctuations introduced by the flash-lamp pumping, and is insensitive to injected-pulse energy variations of more than two orders of magnitude.³

The regen is modeled by numerical integration of the rate equations and careful consideration of the regen dynamics. The regen model described here includes gain saturation in the Nd:YLF laser rod, intracavity losses, lower-laser-level lifetimes, and the active losses introduced by the stabilizer-feed-

back circuit. Careful measurements of the input and output shaped optical pulses from the regen have been made and will be discussed. The calculations on this regen agree well with the measured output of the regen and serve as a model for this important OMEGA component. With this regen model, the entire temporal-pulse-shaping dynamics of OMEGA can now be modeled from the pulse-shaping system to the final on-target pulse shape. This modeling provides us with the capability to accurately produce any desired temporally shaped optical pulse on target for laser-fusion experiments.

Rate Equations

OMEGA's pulse-shaping system produces a shaped optical pulse that is injected into the actively stabilized regenerative amplifier. The output-pulse shape of the regen is determined by gain saturation in the active medium and by the cavity dynamics. The regen model consists of injecting a temporally shaped pulse into the cavity and calculating the new shape after every pass through the cavity. The effect on the pulse shape due to each component is treated separately in the calculation. In this section we discuss the temporal distortion due to a single pass through the gain medium and in the next section incorporate this into the calculation of the overall regen dynamics.

Gain saturation due to a single pass through a gain medium is calculated by solving the laser rate equations⁷

$$\frac{\partial \phi(z,t)}{\partial t} + \frac{c}{n} \frac{\partial \phi(z,t)}{\partial z} = \frac{c \sigma \phi(z,t)}{n} [f_2 N_2(z,t) - f_1 N_1(z,t)], \quad (1a)$$

$$\frac{\partial N_1(z,t)}{\partial t} = \frac{c \sigma \phi(z,t)}{n} [f_2 N_2(z,t) - f_1 N_1(z,t)] + \frac{N_2(z,t)}{\tau_{21}} - \frac{N_1(z,t)}{\tau_{10}}, \quad (1b)$$

$$\frac{\partial N_2(z,t)}{\partial t} = -\frac{c \sigma \phi(z,t)}{n} \times [f_2 N_2(z,t) - f_1 N_1(z,t)] - \frac{N_2(z,t)}{\tau_{21}}, \quad (1c)$$

which describe the evolution of the cavity photon flux ϕ , the atomic population N_1 of the lower laser manifold, and the atomic population N_2 of the upper laser manifold. These rate equations explicitly account for the lifetimes of these upper

and lower laser manifolds. In these equations c is the speed of light in vacuum; n is the index of refraction of the gain medium; σ is the stimulated emission cross section ($1.2 \times 10^{-19} \text{ cm}^2$); and $\tau_{i,j}$ is the relaxation time of the transition from manifold i to manifold j (here, level 0 represents the ground state). The upper laser manifold N_2 consists of two sublevels labeled with energies $E_{2,n}$ ($n = 1, 2$), and the lower laser manifold N_1 consists of six sublevels labeled with energies $E_{1,m}$ ($m = 1$ to 6), two of which are degenerate in energy (sublevels 2 and 3) as shown in Fig. 69.19. The stimulated emission terms in the rate equations [Eqs. (1)] (first terms on the right side in each equation) involve transitions between the sublevels E_{21} and E_{12} as shown in Fig. 69.19; hence, the thermal occupation $f_i N_i$ of these laser-active sublevels is used in these terms. The thermal occupation of these sublevels is calculated by

$$f_1 N_1 = \frac{e^{-E_{12}/kT} N_1}{\sum_{m=1}^{m=6} e^{-E_{1m}/kT}} = 0.207 N_1 \quad (2a)$$

and

$$f_2 N_2 = \frac{e^{-E_{21}/kT} N_2}{e^{-E_{21}/kT} + e^{-E_{22}/kT}} = 0.570 N_2, \quad (2b)$$

where $E_{i,j}$ is the energy of level i sublevel j relative to the lowest

Figure 69.19

The energy levels involved in the 1.053- μm -Nd:YLF laser transition. The arrow shows the laser transition between sublevels within the manifolds shown.

energy level in the manifold, k is the Boltzman constant, and T is the temperature (assumed to be room temperature).

The rate equations [Eqs. (1)] can be solved numerically. We transform these equations along their characteristics in the time-distance plane with the transformation equations

$$x \rightarrow z + \frac{ct}{n} \quad \tau \rightarrow t - \frac{nz}{c}. \quad (3)$$

If we use the chain rule and the substitution

$$\Delta N = \frac{\sigma\phi(x)}{2} [f_2 N_2(z, t) - f_1 N_1(z, t)] dx \quad (4)$$

with $cdt/n = dz = dx/\sqrt{2}$, we get the set of finite-difference equations

$$d\phi = \Delta N, \quad (5a)$$

$$dN_1 = \sqrt{2} \Delta N + \frac{N_2}{\tau_{21}} dt - \frac{N_1}{\tau_{10}} dt, \quad (5b)$$

$$dN_2 = -\sqrt{2} \Delta N - \frac{N_2}{\tau_{21}} dt \quad (5c)$$

for the transformed rate equations. Here we have transformed Eq. (1a) using the transformation equations [Eqs. (3)], and we have left Eqs. (1b) and (1c) untransformed since the photon flux evolves in both space and time, whereas the populations evolve in time only. These equations can be solved numerically given appropriate boundary conditions.

In the model, the photon flux is specified at the entrance face of the laser rod and is given by the temporal profile of the pulse entering the rod. The initial upper-laser-level population is determined from measurements of the laser rod small-signal gain, and for simplicity the initial lower-laser-level population is assumed to be zero. With these boundary conditions, Eqs. (5) can be numerically integrated to yield the photon flux at any time and for any position in the laser rod. Of interest for our calculations is the output-pulse shape specified at the

output face of the laser rod. These equations with these boundary conditions, along with the regen dynamics discussed below, have been solved numerically, and the results are presented below.

Regenerative Amplifier Model

Modeling the regen consists of injecting a pulse with a given pulse shape and energy into the regen and calculating the new pulse shape and energy after each round-trip through the cavity. A single round-trip through the regenerative amplifier is depicted in Fig. 69.20. The pulse first experiences gain through the laser rod followed by propagation to the output-coupling mirror and back. The pulse then experiences gain again followed by propagation to the end mirror and back. Losses due to the output-coupling mirror and the feedback stabilizer (discussed below) are included in the calculation. During propagation of the pulse in the cavity, the upper- and lower-laser-level manifolds are allowed to decay with their respective lifetimes. This calculation for a single pass through the cavity gives the output-pulse shape and energy, given the input-pulse shape and energy for the pass. The output pulse for each pass is used as the input pulse for the next pass through the cavity, and the procedure is repeated for a given number of round-trips through the cavity.

The loss due to the feedback stabilizer depends on several factors. The cavity incorporates two Pockels cells, one of which is feedback controlled. Specific voltages are applied to all four electrodes of the two Pockels cells at specific times.³ During the beginning of the flash-lamp cycle, high losses are introduced into the cavity to allow the gain to build up in the rod. At the peak of the gain, a pulse is injected into the cavity at time t_1 , and all losses are removed from the cavity (with the exception of the static losses here assumed to be 55% in our laser, which includes the 50% output coupler loss) allowing the circulating-pulse energy to increase. The applied voltages after time t_1 are shown schematically in Fig. 69.21(a) (however, not to scale). When the circulating-pulse energy reaches a threshold value (adjusted to $\sim 25 \mu\text{J}$), the feedback stabilizer is activated. At this time (t_2) a dc voltage V_{dc} is applied to one electrode of the first Pockels cell, which introduces a dc loss into the cavity. Simultaneously, a modulated feedback-con-

Figure 69.20

The regen model calculates the pulse shape and energy after a round-trip in the cavity, then iterates for many round-trips. The calculation includes the effects of gain saturation, propagation, and static and feedback losses.

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Figure 69.21

Regen temporal dynamics showing (a) feedback-controlled Pockels cell voltages (not to scale), and (b) measured regen-output envelope filtered to remove individual pulses in the train.

trolled voltage $V(t)$ is applied to an electrode of the feedback-controlled Pockels cell, which introduces a feedback-controlled modulated loss in the cavity. The function of the feedback-controlled modulated loss is to stabilize the circulating pulse energy to a specified constant low value. If the pulse energy falls below (above) this energy, loss is removed (added) to maintain the specified output-pulse energy. Specially designed circuitry for this modulated feedback-controlled voltage³ eliminates pulse-shape distortions caused by fast feedback-voltage changes during pulse propagation through the Pockels cell, as in the previous design.⁸ This ensures that pulse-shape distortions in the regen are due mainly to gain saturation. Finally, after the output-pulse energy is stabilized by the feedback mechanism during this prelude phase, the laser is Q -switched, at which time (t_3) all feedback loss is removed and a Q -switched pulse envelope is allowed to build up. (During this time an adjustable low-level dc loss is left in the cavity to control the final output-pulse energy; however, this loss is not included in the model.) The measured output-pulse train envelope from the regen is shown in Fig. 69.21(b).

The voltage applied to the feedback-controlled Pockels cell during the prelude stabilization is modulated every round-trip so that the Pockels cell transmission is given by

$$T = \cos^2\left(\frac{\pi}{2} \frac{V(t)}{V_{\pi}}\right), \quad (6)$$

where $V(t)$ is the instantaneous value of the modulated voltage difference between the electrodes and V_{π} is the quarter-wave voltage of the Pockels cell. The modulated voltage for a particular pass when the feedback circuitry is active is modeled by

$$V_{i+1} = [V_i + \Delta V_i] e^{-\tau_{rt}/\tau_{fb}}, \quad (7)$$

where V_i is the value of the modulated voltage at the beginning of the pass, ΔV_i is the increase in voltage due to the feedback circuitry, and V_{i+1} is the value of the modulated voltage after the pass. The change in voltage ΔV_i is given by

$$\Delta V_i = \text{pulse energy (J)} \times \text{feedback gain (V/J)}, \quad (8)$$

where the feedback gain is determined by the feedback circuitry. In Eq. (7), the final voltage is allowed to decay every round-trip (round-trip time $\tau_{rt} = 26$ ns) with the exponentially decreasing feedback decay time $\tau_{fb} = 35$ ns. When the laser is Q -switched at time t_3 , all feedback loss is removed from the cavity allowing the free buildup of the Q -switched pulse train.

The above model describes how to calculate the output-pulse shape from the regen given the input-pulse shape. Often it is necessary to calculate the inverse, that is, calculate the required input-pulse shape to the regen that will produce a desired output-pulse shape. A good approximation for this input-pulse shape can be gotten from the output-pulse shape with a simple procedure. A transfer function for the regen can be calculated by using the desired regen-output-pulse shape $[I^{\text{out}}(t)]$ as input to the calculation to obtain a new output-pulse shape [e.g., $I^{\text{new}}(t)$]. The transfer function $T(t)$ for the regen is obtained by dividing these two functions to get

$$T(t) = \frac{I^{\text{new}}(t)}{I^{\text{out}}(t)}. \quad (9)$$

The required input-pulse shape $[I^{\text{in}}(t)]$ can now be calculated with this transfer function and is given by

$$I^{\text{in}}(t) = \frac{I^{\text{out}}(t)}{T(t)}. \quad (10)$$

This simple procedure is used to obtain the required regen-input-pulse shape that will produce the desired regen-output-pulse shape. More importantly, this procedure is useful in producing the desired OMEGA on-target pulse shape.

Experiments

The regen in OMEGA uses a Nd:YLF laser rod pumped to a single-pass, small-signal gain of approximately 2.9. The laser uses a 50% reflecting output coupler, the cavity-round-trip time is 26 ns, and the laser operates at 5 Hz. Typical output energies of the pulse switched out at the peak of the Q -switched envelope are approximately 1.0 mJ.

The measured output-pulse train from the regen is shown in Fig. 69.21(b). The output has been filtered to show only the envelope of the pulse-train output from the regen. It can be seen that the feedback is activated at time t_2 approximately 600 ns after the pulse is injected into the cavity at time $t_1 = 0$. At $t_3 = 2.9 \mu\text{s}$, the laser is Q -switched and a pulse train builds up and decays as the gain is depleted.

Figure 69.22 shows the calculated-output-pulse train from the regen for the above case. Individual pulses within the train are shown. The calculation is based on the model described above with typical values for the regen parameters. Note the good agreement between the measured-output-pulse train in Fig. 69.21(b) and the predictions shown in Fig. 69.22.

E8271

Figure 69.22
Calculated regen-output envelope corresponding to the case measured in Fig. 69.21(b). Individual pulses are shown.

Figure 69.23 shows regen input/output-pulse shapes for a square pulse injected into the regen. The output-pulse shape is the pulse that is switched out at the peak of the Q -switched envelope. The input square pulse (curve plotted with long dashed lines) and measured-regen-output pulse (curve plotted with short dashed lines) are shown in Fig. 69.23, along with the calculated-output-pulse shape (curve plotted with solid line) obtained with the above numerical method using the measured-square-pulse shape as input to the calculation. The regen parameters used in the calculation correspond to the measured regen parameters with slight adjustments to obtain good agreement with the data. By adjusting the regen parameters in this way, the model is calibrated to the data. Once this calibration procedure is performed, the parameters in the model are left unchanged and other shaped pulses can be calculated and compared to measurements.

Figure 69.24 shows the same information as Fig. 69.23, but for a shaped optical pulse injected into the regen. The regen parameters were identical to those used for the calculation in Fig. 69.23. This pulse shape, when injected into OMEGA, will produce a square pulse shape at 351-nm wavelength at the output of OMEGA.

In summary, we have modeled the temporal evolution of a shaped optical pulse injected into our feedback-stabilized regen to a high degree of accuracy. We have solved the rate

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Figure 69.23
Square-pulse distortion from the regen showing the input-pulse shape (long dashed lines), the measured-output-pulse shape (short dashed lines), and the calculated-output-pulse shape (solid line).

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Figure 69.24

Shaped pulse from the regen showing the input-pulse shape (long dashed lines), the measured-output-pulse shape (short dashed lines), and the calculated-output-pulse shape (solid line).

equations including upper- and lower-laser-level lifetimes explicitly. We provide a prescription for determining the injection-pulse shape required to produce a given output-pulse shape from this regen. Finally, with this model of the regen, the entire OMEGA laser system can be modeled, and on-target pulse shapes can be specified in advance by OMEGA users.

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Multiple Scale Derivation of the Relativistic Ponderomotive Force

The ponderomotive force associated with a light wave of variable amplitude¹⁻¹⁰ drives many phenomena that occur in inertial confinement fusion¹¹ and particle acceleration¹² experiments. The existing formula for the ponderomotive force was derived under the assumption that the quiver speed of electrons oscillating in the applied electric field is much less than the speed of light. With the advent of intense laser pulses,¹³ it is important to extend this formula to electron quiver speeds that are comparable to the speed of light.

As an introduction to this subject, we review the derivation of the ponderomotive term in the electron-fluid momentum equation. The standard form of this equation is

$$(\partial_t + \mathbf{v} \cdot \nabla)(\gamma \mathbf{v}) = -(\mathbf{E} + \mathbf{v} \times \mathbf{B}), \quad (1)$$

where

$$\gamma = (1 - v^2)^{-1/2} \quad (2)$$

is the Lorentz factor associated with the fluid velocity and

$$\mathbf{E} = -\partial_t \mathbf{A}, \quad \mathbf{B} = \nabla \times \mathbf{A} \quad (3)$$

in the radiation gauge. These differ from the usual equations in that $\omega t \rightarrow t$, $k\mathbf{x} \rightarrow \mathbf{x}$, $\mathbf{v}/c \rightarrow \mathbf{v}$, $e\mathbf{E}/m\omega c \rightarrow \mathbf{E}$, $e\mathbf{B}/m\omega c \rightarrow \mathbf{B}$, and $e\mathbf{A}/mc^2 \rightarrow \mathbf{A}$.

By using the vector identity¹⁴

$$(\mathbf{v} \cdot \nabla)(\gamma \mathbf{v}) = \nabla \gamma - \mathbf{v} \times [\nabla \times (\gamma \mathbf{v})], \quad (4)$$

one can rewrite the momentum equation as

$$\partial_t(\gamma \mathbf{v} - \mathbf{A}) = \mathbf{v} \times [\nabla \times (\gamma \mathbf{v} - \mathbf{A})] - \nabla \gamma, \quad (5)$$

from which follows the relativistic vorticity equation

$$\partial_t [\nabla \times (\gamma \mathbf{v} - \mathbf{A})] = \nabla \times \{ \mathbf{v} \times [\nabla \times (\gamma \mathbf{v} - \mathbf{A})] \}. \quad (6)$$

For a plasma that is at rest before the laser pulse arrives, $\nabla \times (\gamma \mathbf{v} - \mathbf{A}) = 0$ initially. Equation (6) ensures that $\nabla \times (\gamma \mathbf{v} - \mathbf{A}) = 0$ for all time. Thus, the momentum equation can be rewritten as¹⁴

$$\partial_t(\mathbf{u} - \mathbf{A}) = -\nabla \gamma, \quad (7)$$

where the fluid momentum $\mathbf{u} = \gamma \mathbf{v}$. It follows from this definition that $\gamma = (1 + u^2)^{1/2}$.

The ponderomotive term on the right side of Eq. (7) is valid for arbitrary laser intensity. Together with the continuity and Maxwell equations, it allows one to analyze the interaction of a laser pulse with an electron fluid. However, there is a tradition in plasma physics of looking at the same phenomenon from different viewpoints. By doing so, one often gains physical insight into the phenomenon under study. The ponderomotive term in Eq. (7) is not the force on a Lagrangian fluid element or a single electron. Consequently, it cannot be used as the foundation of a single-particle or kinetic analysis of the interaction of a laser pulse with a plasma.

In the following sections we present (1) an analytical study of the motion of an electron in a light wave of constant amplitude; (2) using the results of this study, a heuristic derivation of the formula for the ponderomotive force associated with a light wave of variable amplitude; (3) numerical and analytical verification of this formula; and, finally, (4) a summary of the results.

Particle Motion in a Plane Wave

The motion of a charged particle, of charge q and mass m , in an electromagnetic field is governed by the equation¹⁵

$$d_\tau(u_\mu + a_\mu) = u^\nu \partial_\mu a_\nu, \quad (8)$$

where τ is the proper time of the particle multiplied by c , u^μ is the four-velocity of the particle divided by c , a^μ is the four-potential of the field multiplied by q/mc^2 and $\partial_\mu = \partial/\partial x^\mu$. For an elliptically polarized field

$$a^\mu = (0, 0, e_y \cos \phi, e_z \sin \phi), \quad (9)$$

where $e_y = e\delta$, $e_z = e(1 - \delta^2)^{1/2}$, and $\phi = t - x$.

The motion of a charged particle in a plane wave is well known.¹⁶⁻¹⁹ We present an analysis of this motion here because it is the foundation of analyses presented later in this article. Since the four-potential does not depend on y or z , it follows from Eq. (8) that

$$d_\tau(\mathbf{u}_\perp + \mathbf{a}_\perp) = 0. \quad (10)$$

Transverse canonical momentum is conserved. It follows from Eq. (10) that

$$\mathbf{u}_\perp(\tau) = \mathbf{u}_\perp(0) + \mathbf{a}_\perp(0) - \mathbf{a}_\perp(\tau). \quad (11)$$

The t and x components of Eq. (8) are

$$d_\tau \gamma = \frac{1}{2} \partial_t u_\perp^2, \quad d_\tau u_x = -\frac{1}{2} \partial_x u_\perp^2. \quad (12)$$

Since the four-potential is a function of $t - x$, it follows from Eqs. (12) that

$$d_\tau(\gamma - u_\parallel) = 0. \quad (13)$$

Because the particle gains energy and momentum at the expense of the field, the ratio of particle momentum to particle kinetic energy is identical to the ratio of field momentum to field energy, which is 1 in the units of Eq. (8). By combining Eq. (13) with the definition of γ , one can show that

$$u_\parallel(\tau) = u_\parallel(0) + \frac{u_\perp^2(\tau) - u_\perp^2(0)}{2[\gamma(0) - u_\parallel(0)]}. \quad (14)$$

The corresponding equation for $\gamma(t)$ follows from Eqs. (13) and (14). Because the transverse potential $\mathbf{a}_\perp$ is a function of ϕ

rather than τ , Eqs. (11) and (13) describe the particle momentum implicitly. One can make this description explicit and determine the particle trajectory $x^\mu(\tau)$ by using the result

$$d_\tau \phi = \gamma(0) - u_\parallel(0). \quad (15)$$

The proper frequency of the wave is constant.

It is clear from Eqs. (11), (14), and (15) that the particle motion is a superposition of sinusoidal oscillations in τ and steady drifts in τ . It follows from Eq. (11) that the transverse drifts are given by

$$\begin{aligned} \langle u_y \rangle &= u_y(0) + e_y \cos(-x_0), \\ \langle u_z \rangle &= u_z(0) + e_z \sin(-x_0), \end{aligned} \quad (16)$$

where $\langle \cdot \rangle$ denotes the τ -average $\int_0^{2\pi} \cdot d\tau/2\pi$ and $(x_0, 0, 0)$ is the initial position of the particle. By decomposing the longitudinal momentum into its oscillatory component

$$u_\parallel(\tau) - \langle u_\parallel \rangle = \frac{u_\perp^2(\tau) - \langle u_\perp^2 \rangle}{2[\gamma(0) - u_\parallel(0)]} \quad (17)$$

and its drift component

$$\langle u_\parallel \rangle = u_\parallel(0) + \frac{\langle u_\perp^2 \rangle - u_\perp^2(0)}{2[\gamma(0) - u_\parallel(0)]}, \quad (18)$$

and combining Eqs. (11) and (18), one can show that the longitudinal drift is given by

$$\begin{aligned} \langle u_x \rangle &= u_x(0) + \left[4\langle u_y \rangle e_y \cos(-x_0) \right. \\ &\quad + 4\langle u_z \rangle e_z \sin(-x_0) - e_y^2 \cos(-2x_0) \\ &\quad \left. + e_z^2 \cos(-2x_0) \right] / 4[\gamma(0) - u_x(0)]. \end{aligned} \quad (19)$$

For linear polarization Eq. (19) reduces to

$$\begin{aligned} \langle u_x \rangle &= u_x(0) + \left[4\langle u_y \rangle e \cos(-x_0) \right. \\ &\quad \left. - e^2 \cos(-2x_0) \right] / 4[\gamma(0) - u_x(0)], \end{aligned} \quad (20)$$

whereas for circular polarization it reduces to

$$\begin{aligned} \langle u_x \rangle &= u_x(0) + e \left[\langle u_y \rangle \cos(-x_0) \right. \\ &\quad \left. + \langle u_z \rangle \sin(-x_0) \right] / \sqrt{2} [\gamma(0) - u_x(0)]. \end{aligned} \quad (21)$$

The corresponding equations for $\langle \gamma \rangle$ follow from Eq. (13) and Eqs. (19)–(21).

For completeness, a covariant analysis of the particle motion is given in **Appendix A**.

Heuristic Derivation of the Ponderomotive Force

The method used to solve Eq. (8) for a plane wave of constant amplitude can also be used when the wave amplitude e is a function of $t-x$. In fact, Eqs. (11), (14), and (15) are still valid. When the wave amplitude varies slowly compared to the wave phase, the particle motion consists of an oscillation about a guiding center and a guiding-center drift that varies slowly. As the guiding center drifts, the oscillation amplitude follows the wave amplitude at the guiding center adiabatically.

To describe this motion quantitatively, let ξ^μ be the position four-vector of the guiding center and $v^\mu = d_\tau \xi^\mu$ be the associated four-momentum. The ponderomotive four-force is the proper rate of change of the guiding-center four-momentum. One might expect this four-force to also be the average rate of change of the particle four-momentum. However, by averaging the transverse particle motion, one finds that

$$\begin{aligned} \langle d_\tau u_y \rangle &\approx [d_{\tau_0} e_y(\tau_0)] \cos(\tau_0), \\ \langle d_\tau u_z \rangle &\approx [d_{\tau_0} e_z(\tau_0)] \sin(\tau_0), \end{aligned} \quad (22)$$

where τ_0 is the initial phase with respect to which the average is taken. Because the oscillation amplitude changes during each oscillation, the transverse components of the momentum change by amounts that depend on the initial phase. However, it follows from Eq. (11) that the transverse components of the guiding-center momentum are constant. Thus, if one is to determine the ponderomotive four-force by averaging, one must discount terms that depend on the initial phase. With this caveat added to the definition of $\langle \cdot \rangle$, one can write

$$d_\tau v_y = \langle d_\tau u_y \rangle \approx 0, \quad d_\tau v_z = \langle d_\tau u_z \rangle \approx 0 \quad (23)$$

and show that

$$d_\tau v_x = \langle d_\tau u_x \rangle \approx d_\tau (e_y^2 + e_z^2) / 4 [\gamma(0) - u_{||}(0)]. \quad (24)$$

By using the relationship between x and ϕ , and Eq. (15), one can show that $d_\tau = -[\gamma(0) - u_{||}(0)] \partial_x$. It follows from this result and Eq. (24) that

$$d_\tau v_x \approx -\partial_x (e^2/4). \quad (25)$$

In a similar way, one can show that

$$d_\tau v_t \approx \partial_t (e^2/4). \quad (26)$$

By using the facts that $e^2/2 = \langle a_\perp^2 \rangle$ and $a_\perp^2 = -a_v a^v$, one can rewrite Eqs. (23), (25), and (26) as

$$d_\tau v_\mu \approx -\partial_\mu \langle a_v a^v / 2 \rangle. \quad (27)$$

The second term in this relation is the ponderomotive four-force.

The guiding-center Eq. (27) was derived for the special case in which e is a function of $t-x$. However, the principle of Lorentz covariance suggests that it is valid for the general case in which e is a function of t , x , y , and z . Consequently, we postulate that²⁰

$$d_{\tau\tau}^2 \xi_\mu = -\partial_\mu \langle a_v a^v / 2 \rangle |_{\xi_\mu} \quad (28)$$

and the initial guiding-center momentum in a wave of variable amplitude is identical to the particle drift momentum in a wave of constant amplitude, which is given by Eqs. (16) and (19). For future reference, Eq. (28) has associated with it the conservation equation

$$d_\tau (v_\mu v^\mu / 2 + \langle a_v a^v / 2 \rangle) = 0. \quad (29)$$

Numerical Study of the Particle Motion

To test the guiding-center model described in the previous section, we studied three representative examples numerically. The first example concerns a particle that moves in front of a laser pulse. We considered a wide, circularly polarized pulse,

with $e = 3 \sin^2[0.05(t-x)]$, and chose $u_x(0) = 1$, $u_y(0) = 1$, and $u_z(0) = 1$. Because the pulse propagates at the speed of light, it overtakes the particle. The resulting particle motion is illustrated in Figs. 69.25 and 69.26, in which the solid lines denote the particle trajectory, determined numerically from Eq. (8) and the initial conditions, and the dashed lines denote the guiding-center trajectory, determined numerically from Eqs. (28), (16), and (19). As the pulse overtakes the particle, the amplitudes of the transverse components of the oscillation increase and decrease in proportion to the pulse intensity. However, there is no change in the transverse components of the average momentum, and the particle exits the pulse with $u_y = 1$ and $u_z = 1$. The amplitude of the longitudinal component of the oscillation also increases and decreases in proportion to the pulse intensity. However, because Eq. (14), which describes the relation between the longitudinal and transverse components of the momentum, is nonlinear, the longitudinal component of the average momentum changes. This change

can be analyzed quantitatively. It follows from the t and x components of Eq. (28), and the assumed dependence of e on $t-x$, that

$$d_\tau(v_t - v_x) = 0. \quad (30)$$

Since v_y and v_z are constant, Eq. (29) reduces to

$$d_\tau[(v_t^2 - v_x^2)/2 - e^2/4] = 0. \quad (31)$$

By combining Eqs. (30) and (31) with the initial conditions, one can show that $v_t = 2 + e^2/4$ and $v_x = 1 + e^2/4$. At the peak of the pulse $v_x = 13/4$, in agreement with Fig. 69.25(a). Because the x component of the ponderomotive force is positive in the front of the pulse and negative in the back of the pulse, the guiding center is accelerated and decelerated by equal amounts. In this example the correspondence between the guiding-center motion and the particle motion is excellent.

Figure 69.25

Particle motion (solid line) and guiding-center motion (dashed line) caused by a circularly polarized pulse with amplitude $e = 3 \sin^2[0.05(t-x)]$. Initially, $u_x = 1$, $u_y = 1$, and $u_z = 1$. (a) The x component of the momentum. (b) The x component of the displacement caused by the pulse. The initial drift upon which this displacement is superimposed is not shown.

Figure 69.26

Particle motion (solid line) and guiding-center motion (dashed line) caused by a circularly polarized pulse with amplitude $e = 3 \sin^2[0.05(t-x)]$. Initially, $u_x = 1$, $u_y = 1$, and $u_z = 1$. (a) The y component of the momentum. (b) The z component of the momentum.

The second example concerns a particle that is born inside a laser pulse by high-field ionization.²¹ We considered a long pulse that is linearly polarized in the y direction, with $e = \cos^2(0.05z)$, and chose $u_x(0) = 0$, $u_y(0) = 0$, and $u_z(0) = 0$. The resulting particle motion is illustrated in Figs. 69.27 and 69.28. The particle is born near the propagation axis of the pulse and is pushed outward by the z component of the ponderomotive force. As the particle moves outward, the amplitudes of the longitudinal and transverse components of the oscillation decrease in proportion to the pulse intensity. This transverse expulsion can be analyzed quantitatively. Since v_r , v_x , and v_y are all constant, Eq. (29) reduces to

$$d_\tau(v_z^2/2 + e^2/4) = 0, \quad (32)$$

in which $v_z^2/2$ plays the role of kinetic energy and $e^2/4$ plays the role of potential energy. It follows from Eq. (32) and the initial conditions that $v_z^2 \approx (1 - e^2)/2$. As the guiding center exits the pulse, $v_z \approx 1/\sqrt{2}$, in agreement with Fig. 69.27(a).

Figure 69.27

Particle motion (solid line) and guiding-center motion (dashed line) caused by a linearly polarized pulse with amplitude $e_y = \cos^2(0.05z)$. Initially, $u_x = 0$, $u_y = 0$, and $u_z = 0$. (a) The z component of the momentum. (b) The z component of the displacement.

Although the particle is born at rest, it exits the pulse with $u_x \approx 3/4$ and $u_y \approx 1$. This behavior is consistent with Eqs. (16) and (19). In this example the correspondence between the guiding-center motion and the particle motion is excellent.

The third example concerns a particle that is injected into a laser pulse from the side. We considered a long pulse that is linearly polarized in the y direction, with $e = \sin^2(0.05y)$, and chose $u_x(0) = 0.0$, $u_y(0) = 0.7$, and $u_z(0) = 0.0$. The resulting particle motion is illustrated in Figs. 69.29 and 69.30. As the particle moves inward, the amplitudes of the longitudinal and transverse components of the oscillation increase in proportion to the pulse intensity. However, the y component of the ponderomotive force opposes the inward motion, and the particle is repelled just before it reaches the propagation axis of the pulse. As the particle moves outward, the amplitudes of the longitudinal and transverse components of the oscillation decrease in proportion to the pulse intensity. This transverse repulsion can be analyzed quantitatively. Since v_r , v_x , and v_z are all constant, Eq. (29) reduces to

Figure 69.28

Particle motion (solid line) and guiding-center motion (dashed line) caused by a linearly polarized pulse with amplitude $e_y = \cos^2(0.05z)$. Initially, $u_x = 0$, $u_y = 0$, and $u_z = 0$. (a) The x component of the momentum. (b) The y component of the momentum.

Figure 69.29

Particle motion (solid line) and guiding-center motion (dashed line) caused by a linearly polarized pulse with amplitude $e_y = \sin^2(0.05y)$. Initially, $u_x = 0.0$, $u_y = 0.7$, and $u_z = 0.0$. (a) The y component of the momentum. (b) The y component of the displacement.

Figure 69.30

Particle motion (solid line) and guiding-center motion (dashed line) caused by a linearly polarized pulse with amplitude $e_y = \sin^2(0.05y)$. Initially, $u_x = 0.0$, $u_y = 0.7$, and $u_z = 0.0$. (a) The z component of the momentum. (b) The x component of the momentum.

$$d_\tau(v_y^2/2 + e^2/4) = 0. \quad (33)$$

It follows from Eq. (33) and the initial conditions that $v_y^2 \approx (1 - e^2)/2$. The outward guiding-center trajectory is the inverse of the inward trajectory. In this example the correspondence between the guiding-center motion and the particle motion is good. We found the correspondence to be even better for gentler gradients in pulse intensity.

In Figs. 69.25–69.30 the particle and guiding-center positions were plotted as functions of the proper time. We verified numerically that plotting the spatial components of the guiding-center position as functions of the temporal component of the guiding-center position produces the correct guiding-center motion in the laboratory frame.

Multiple Scale Analysis of the Particle Motion

In this section we verify Eq. (28) analytically. Because the fast variation of the four-potential depends on the phase rather

than the proper time, it is advantageous to change the independent variable in Eq. (8) from τ to ϕ . The result is

$$d_\phi(d_\tau \phi d_\phi x_\mu + a_\mu) = d_\phi x^\nu \partial_\mu a_\nu, \quad (34)$$

where $d_\tau \phi = (d_\phi x^\nu d_\phi x_\nu)^{-1/2}$. The resolution of Eq. (34) into longitudinal and transverse components is facilitated by the introduction of the four-vector k^μ , which is defined by the equation $\phi = k^\nu x_\nu$, and the four-vector l^μ , which is defined by the equations $l^\nu l_\nu = 0$, $k^\nu l_\nu = 2$, and $a^\nu l_\nu = 0$, where a^μ is the transverse four-potential of a plane wave. In the laboratory frame $k^\mu = (1, 1, 0, 0)$ and $l^\mu = (1, -1, 0, 0)$. By using these four-vectors one can write

$$x^\mu = y^\mu + \theta k^\mu / 2 + \phi l^\mu / 2, \quad (35)$$

where $\theta = l^\nu x_\nu$. The transverse position four-vector satisfies the equations $k^\nu y_\nu = 0$ and $l^\nu y_\nu = 0$. In a similar way, one can write

$$a^\mu = b^\mu + qk^\mu/2 + pl^\mu/2, \quad (36)$$

where the transverse four-potential satisfies the equations $k^\nu b_\nu = 0$ and $l^\nu b_\nu = 0$. By substituting the decompositions (35) and (36) into Eq. (34) and collecting like terms, one can show

$$\frac{d}{d\phi} \left(\frac{1}{\sigma} \frac{dy_\mu}{d\phi} + b_\mu \right) = \frac{\partial b^\nu}{\partial y^\mu} \frac{dy_\nu}{d\phi} + \frac{1}{2} \left(\frac{\partial p}{\partial y^\mu} \frac{d\theta}{d\phi} + \frac{\partial q}{\partial y^\mu} \right), \quad (37)$$

$$\frac{d}{d\phi} \left(\frac{1}{\sigma} + p \right) = 2 \frac{\partial b^\nu}{\partial \theta} \frac{dy_\nu}{d\phi} + \frac{\partial p}{\partial \theta} \frac{d\theta}{d\phi} + \frac{\partial q}{\partial \theta}, \quad (38)$$

$$\frac{d}{d\phi} \left(\frac{1}{\sigma} \frac{d\theta}{d\phi} + q \right) = 2 \frac{\partial b^\nu}{\partial \phi} \frac{dy_\nu}{d\phi} + \frac{\partial p}{\partial \phi} \frac{d\theta}{d\phi} + \frac{\partial q}{\partial \phi}, \quad (39)$$

where

$$\sigma = \left(d_\phi y^\nu d_\phi y_\nu + d_\phi \theta \right)^{1/2}. \quad (40)$$

Equation (39) can be derived from Eqs. (37) and (38), as shown in **Appendix B**, and need not be considered further.

One can solve Eqs. (37) and (38) by using multiple scale analysis. Let ε be a measure of the rate at which the wave amplitude varies relative to the rate at which the phase varies. We introduce the scales

$$\phi_0 = \phi, \quad \phi_1 = \varepsilon \phi \quad (41)$$

to resolve the fast oscillation and the slow change in the guiding-center drift, respectively. It follows that

$$\frac{d}{d\phi} = \frac{d}{d\phi_0} + \varepsilon \frac{d}{d\phi_1}. \quad (42)$$

We used the notation $d/d\phi_0$ and $d/d\phi_1$ in Eq. (42) to distinguish these convective derivatives from the partial derivatives of the four-potential. We assume that the dependent variables can be written as

$$y_\mu \approx \varepsilon^{-1} y_\mu^{(-1)}(\phi_1) + y_\mu^{(0)}(\phi_0, \phi_1) + \varepsilon y_\mu^{(1)}(\phi_0, \phi_1), \quad (43)$$

$$\theta \approx \varepsilon^{-1} \theta^{(-1)}(\phi_1) + \theta^{(0)}(\phi_0, \phi_1) + \varepsilon \theta^{(1)}(\phi_0, \phi_1).$$

The variables $y_\mu^{(-1)}$ and $\theta^{(-1)}$ describe the guiding-center drift, which changes on the slow scale ϕ_1 . The variables $y_\mu^{(0)}$ and $\theta^{(0)}$ describe the fast oscillation of the particle about the guiding center, the amplitude of which changes on the slow scale.

The four-potential satisfies Maxwell's wave equation¹⁴

$$\left(\partial^\lambda \partial_\lambda g_\nu^\mu - \partial^\mu \partial_\nu \right) a^\nu = 0, \quad (44)$$

where $g_\nu^\mu = \text{diag}(1, -1, -1, -1)$ is the metric four-tensor. For a wave of constant amplitude, $a_\mu(x^\nu) = b_\mu^{(0)}(\phi_0)$. For a wave of variable amplitude we assume that

$$a_\mu(x^\nu) \approx a_\mu^{(0)}(\phi_0 | \varepsilon x^\nu) + \varepsilon a_\mu^{(1)}(\phi_0 | \varepsilon x^\nu). \quad (45)$$

Each contribution to the four-potential and its derivatives can be written approximately as

$$\begin{aligned} a(\phi_0 | \varepsilon y_\nu, \varepsilon \theta, \varepsilon \phi) \approx & a \left(\phi_0 \left| y_\nu^{(-1)}, \theta^{(-1)}, \phi_1 \right. \right) \\ & + \varepsilon y^{(0)\nu} \partial_\nu a \left(\phi_0 \left| y_\nu^{(-1)}, \theta^{(-1)}, \phi_1 \right. \right) \\ & + \varepsilon \theta^{(0)} \partial_\theta a \left(\phi_0 \left| y_\nu^{(-1)}, \theta^{(-1)}, \phi_1 \right. \right). \end{aligned} \quad (46)$$

The first term on the right side of Eq. (46) is the contribution evaluated at the guiding center, and the second and third terms are the deviations from this average contribution that are felt by the particle as it oscillates about the guiding center. The corresponding approximation for the convective derivative of the four-potential is discussed in **Appendix C**. Henceforth, we will use $\bar{a}$ to denote the guiding-center contribution

$$a \left(\phi_0 \left| y_\nu^{(-1)}, \theta^{(-1)}, \phi_1 \right. \right).$$

To proceed further one substitutes Eqs. (42), (43), and (46) in Eqs. (37) and (38) and collects terms of like order. The order ε^{-1} equations are satisfied identically by Ansatz (43).

The order 1 equations are

$$\frac{d}{d\phi_0} \left[\frac{1}{\sigma^{(0)}} \frac{dy_\mu^{(0)}}{d\phi_0} + \bar{b}_\mu^{(0)} \right] = 0, \quad (47)$$

$$\frac{d}{d\phi_0} \left[\frac{1}{\sigma^{(0)}} \right] = 0, \quad (48)$$

where

$$\sigma^{(0)} = \left\{ \left[d_0 y^{(0)\nu} + d_1 y^{(-1)\nu} \right] \times \left[d_0 y_V^{(0)} + d_1 y_V^{(-1)} \right] + d_0 \theta^{(0)} + d_1 \theta^{(-1)} \right\}^{1/2} \quad (49)$$

and $d_n = d/d\phi_n$.

Equation (47) is the analog of Eq. (10). It follows from the former equation that

$$d_0 y_\mu^{(0)} = -\sigma^{(0)} \bar{b}_\mu^{(0)}. \quad (50)$$

The arbitrary function of ϕ_1 that results from the ϕ_0 integration can be neglected because $y_\mu^{(-1)}$ already accounts for the slowly varying drift with which this function is associated. Equations (48) and (49) do not resemble any of the equations in the section **Particle Motion in a Plane Wave**. However, different forms of the latter equations are discussed in **Appendix A**, from which it is clear that Eqs. (48) and (49) comprise the analog of Eq. (A9). It follows from Eq. (48) that $\sigma^{(0)}$ is a function of ϕ_1 alone. This result is the analog of Eq. (15) and facilitates the integration of Eq. (50). By combining Eqs. (49) and (50), and equating the oscillatory and slowly varying terms that result, one can show that

$$d_0 \theta^{(0)} = 2\sigma^{(0)} d_1 y^{(-1)\nu} \bar{b}_\nu^{(0)} + [\sigma^{(0)}]^2 \left[\langle \bar{b}^{(0)\nu} \bar{b}_\nu^{(0)} \rangle - \bar{b}^{(0)\nu} \bar{b}_\nu^{(0)} \right] \quad (51)$$

and

$$d_1 \theta^{(-1)} + d_1 y^{(-1)\nu} d_1 y_V^{(-1)} = [\sigma^{(0)}]^2 \left[1 - \langle \bar{b}^{(0)\nu} \bar{b}_\nu^{(0)} \rangle \right]. \quad (52)$$

Equation (51) is the analog of Eq. (17) and the oscillatory part of Eq. (A7), and is easy to integrate.

Now consider the initial condition on the order 1 four-momentum. Consistent with Eq. (35), one can write the initial four-momentum as

$$u^\mu(0) = v^\mu(0) + l^\nu u_\nu(0) k^\mu / 2 + k^\nu u_\nu(0) l^\mu / 2. \quad (53)$$

It follows immediately that

$$d_1 y_\mu^{(-1)}(0) = \sigma^{(0)} v_\mu(0) - d_0 y_\mu^{(0)}(0), \quad (54)$$

$$d_1 \theta^{(-1)}(0) = \sigma^{(0)} l^\nu u_\nu(0) - d_0 \theta^{(0)}(0). \quad (55)$$

Equation (54) is the analog of Eqs. (16), and Eq. (55) is consistent with Eqs. (17) and (18).

The order ε equations are

$$\begin{aligned} & \frac{d}{d\phi_1} \left\{ \frac{1}{\sigma^{(0)}} \left[\frac{dy_\mu^{(0)}}{d\phi_0} + \frac{dy_\mu^{(-1)}}{d\phi_1} \right] + \bar{b}_\mu^{(0)} \right\} \\ & + \frac{d}{d\phi_0} \left\{ \frac{1}{\sigma^{(0)}} \left[\frac{dy_\mu^{(1)}}{d\phi_0} + \frac{dy_\mu^{(0)}}{d\phi_1} \right] + \bar{b}_\mu^{(1)} \right\} \\ & - \frac{d}{d\phi_0} \left\{ \frac{\sigma^{(1)}}{2[\sigma^{(0)}]^3} \left[\frac{dy_\mu^{(0)}}{d\phi_0} + \frac{dy_\mu^{(-1)}}{d\phi_1} \right] \right\} \\ & = \left[\frac{dy^{(0)\nu}}{d\phi_0} + \frac{dy^{(-1)\nu}}{d\phi_1} \right] \frac{\partial \bar{b}_\nu^{(0)}}{\partial y^\mu} \end{aligned} \quad (56)$$

and

$$\begin{aligned} & \frac{d}{d\phi_1} \left[\frac{1}{\sigma^{(0)}} \right] - \frac{d}{d\phi_0} \left\{ \frac{\sigma^{(1)}}{2[\sigma^{(0)}]^3} + \bar{p}^{(1)} \right\} \\ & = 2 \left[\frac{dy^{(0)\nu}}{d\phi_0} + \frac{dy^{(-1)\nu}}{d\phi_1} \right] \frac{\partial \bar{b}_\nu^{(0)}}{\partial \theta}, \end{aligned} \quad (57)$$

where

$$\sigma^{(1)} = \frac{d\theta^{(1)}}{d\phi_0} + \frac{d\theta^{(0)}}{d\phi_1} + 2 \left[\frac{dy^{(0)\nu}}{d\phi_0} + \frac{dy^{(-1)\nu}}{d\phi_1} \right] \left[\frac{dy^{(1)\nu}}{d\phi_0} + \frac{dy^{(0)\nu}}{d\phi_1} \right], \quad (58)$$

and $\bar{b}_\mu^{(1)}$ and $\bar{p}^{(1)}$ represent the sum of the order ε four-potential and the order ε corrections to the order-1 four-potential caused by the oscillation of the particle about the guiding center.

Although Eqs. (56)–(58) are lengthy, they do not need to be solved in their entirety. By equating the slowly varying terms in Eqs. (56) and (57), one can show that

$$\frac{1}{\sigma^{(0)}} \frac{d}{d\phi_1} \left[\frac{1}{\sigma^{(0)}} \frac{dy_\mu^{(-1)}}{d\phi_1} \right] = -\frac{1}{2} \frac{\partial \langle \bar{b}^{(0)\nu} \bar{b}_\nu^{(0)} \rangle}{\partial y^\mu} \quad (59)$$

and

$$\frac{1}{\sigma^{(0)}} \frac{d}{d\phi_1} \left[\frac{1}{\sigma^{(0)}} \right] = -\frac{\partial \langle \bar{b}^{(0)\nu} \bar{b}_\nu^{(0)} \rangle}{\partial \theta}. \quad (60)$$

It follows from Eq. (52) that

$$\begin{aligned} & \frac{1}{\sigma^{(0)}} \frac{d}{d\phi_1} \left[\frac{1}{\sigma^{(0)}} \frac{d\theta^{(-1)}}{d\phi_1} \right] \\ &= \frac{1 - \langle \bar{b}^{(0)\nu} \bar{b}_\nu^{(0)} \rangle}{\sigma^{(0)}} \frac{d\sigma^{(0)}}{d\phi_1} - \frac{d \langle \bar{b}^{(0)\nu} \bar{b}_\nu^{(0)} \rangle}{d\phi_1} \\ & \quad - \frac{1}{\sigma^{(0)}} \frac{d}{d\phi_1} \left\{ \sigma^{(0)} \left[\frac{1}{\sigma^{(0)}} \frac{dy^{(-1)\nu}}{d\phi_1} \right] \left[\frac{1}{\sigma^{(0)}} \frac{dy_\nu^{(-1)}}{d\phi_1} \right] \right\}. \quad (61) \end{aligned}$$

When applied to any guiding-center quantity, the operator

$$\frac{d}{d\phi_1} = \frac{dy^{(-1)\nu}}{d\phi_1} \frac{\partial}{\partial y^\nu} + \frac{\partial}{\partial \phi_1} + \frac{d\theta^{(-1)}}{d\phi_1} \frac{\partial}{\partial \theta}. \quad (62)$$

By combining Eq. (61) with Eqs. (59), (60), and (62), one can show that

$$\frac{1}{\sigma^{(0)}} \frac{d}{d\phi_1} \left[\frac{1}{\sigma^{(0)}} \frac{d\theta^{(-1)}}{d\phi_1} \right] = -\frac{\partial \langle \bar{b}^{(0)\nu} \bar{b}_\nu^{(0)} \rangle}{\partial \phi_1}. \quad (63)$$

Recall that the preceding derivation of Eq. (63) is based on Eq. (38). Had we analyzed Eq. (39) instead, we would have needed to determine $b_\nu^{(1)}$, $p^{(1)}$, $q^{(1)}$, and $y_\nu^{(1)}$ explicitly.

In the notation of this section, Eq. (28) can be rewritten as

$$\frac{d^2 x_\mu^{(-1)}}{d\tau_1^2} = -\frac{1}{2} \frac{\partial \langle \bar{b}^{(0)\nu} \bar{b}_\nu^{(0)} \rangle}{\partial x_1^\mu}, \quad (64)$$

where $\tau_1 = \varepsilon \tau$ and $x_1^\mu = \varepsilon x^\mu$. Since $d\phi_1/d\tau_1 \approx 1/\sigma^{(0)}$, Eq. (59) is the transverse part of Eq. (64). By contracting Eq. (64) with k^μ and l^μ , and using the identities $k^\mu \partial_\mu = 2\partial_\theta$ and $l^\mu \partial_\mu = 2\partial_\phi$, and the fact that $\phi \approx k^\mu x_\mu^{(-1)}$, one can show that Eqs. (60) and (63) are equivalent to the longitudinal part of Eq. (64). Thus, Eq. (28) is correct.

Finally, notice that Eq. (64) for the guiding-center drift is written in terms of the proper time, which includes the effects of the oscillation about the guiding center. Although this fact does not affect the utility of Eq. (64), it calls into question the aesthetic qualities of the equation. Just as the proper time is defined by the equation $d\tau = (dx^\nu dx_\nu)^{1/2}$, one can define the drift time by the equation $ds = [dx^{(-1)\nu} dx_\nu^{(-1)}]^{1/2}$.

It follows from this definition, Eq. (52), and the discussion of the preceding paragraph that

$$\frac{ds_1}{d\tau_1} = \left[1 - \langle \bar{b}^{(0)\nu} \bar{b}_\nu^{(0)} \rangle \right]^{1/2}. \quad (65)$$

Equation (65) can be used to write Eq. (28) in terms of the drift time.

Summary

In this article we solved the equation of motion for an electron in a plane wave. We used this solution and the principle of Lorentz covariance to deduce a formula for the ponderomotive force exerted by an intense laser pulse on an electron. We verified this formula numerically, for three cases of current interest, and analytically, using the method of multiple scales.

The aforementioned formula can be used to study the effects of the radial ponderomotive force on laser-plasma interactions. For particle accelerators, these effects include the divergence of an electron bunch that is accelerated by a laser pulse,²² the relativistic focusing of the pulse, and electron cavitation and magnetic field generation in the wake of the pulse.

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Appendix A: Covariant Analysis of the Particle Motion in a Plane Wave

The motion of a charged particle in an electromagnetic field is governed by Eq. (8). For a plane wave the four-potential a^μ is a function of the phase $\phi = k^\nu x_\nu$. It follows that $\partial_\mu a_\nu = k_\mu a'_\nu$, where $' = d/d\phi$, and, hence, that

$$d_\tau(u_\mu + a_\mu) = u^\nu k_\mu a'_\nu. \quad (A1)$$

By substituting the decomposition

$$u_\mu(\tau) = v_\mu(\tau) + k^\nu u_\nu(\tau) l_\mu / 2 + l^\nu u_\nu(\tau) k_\mu / 2 \quad (A2)$$

into Eq. (A1), where l^ν was defined after Eq. (8), and v_μ satisfies the equations $k^\nu v_\nu = 0$ and $l^\nu v_\nu = 0$, one can show that

$$d_\tau(v_\mu + a_\mu) = 0, \quad d_\tau(k^\nu u_\nu) = 0, \quad d_\tau(l^\nu u_\nu) = 2v^\nu a'_\nu. \quad (A3)$$

It follows from the first of Eqs. (A3) that

$$v_\mu(\tau) = v_\mu(0) + a_\mu(0) - a_\mu(\tau). \quad (A4)$$

It follows from the second of Eqs. (A3) that

$$k^\nu u_\nu(\tau) = k^\nu u_\nu(0) \quad (A5)$$

and, hence, that

$$\phi = k^\nu u_\nu(0) \tau. \quad (A6)$$

Equations (A4) and (A6) determine $v_\mu(\tau)$ explicitly. There are at least three ways to obtain an expression for $l^\nu u_\nu$. In the first approach, one uses Eq. (A4) to rewrite the right side of the third of Eqs. (A3) in terms of a^μ . It follows from this equation and Eq. (A6) that

$$\begin{aligned} l^\nu u_\nu(\tau) &= l^\nu u_\nu(0) + 2[v^\nu(0) + a^\nu(0)] \\ &\quad \times [a_\nu(\tau) - a_\nu(0)] / k^\nu u_\nu(0) \\ &\quad + [a^\nu(0)a_\nu(0) - a^\nu(\tau)a_\nu(\tau)] / k^\nu u_\nu(0). \end{aligned} \quad (A7)$$

In the second approach, one uses Eq. (A4) to rewrite the right side of the third of Eqs. (A3) in terms of v^μ . It follows from this equation and Eq. (A6) that

$$\begin{aligned} l^\nu u_\nu(\tau) &= l^\nu u_\nu(0) + [v^\nu(0)v_\nu(0) \\ &\quad - v^\nu(\tau)v_\nu(\tau)] / k^\nu u_\nu(0). \end{aligned} \quad (A8)$$

In the third approach one uses decomposition (A2) to rewrite the identity $u^\nu u_\nu = 1$ as

$$(k^\nu u_\nu)(l^\nu u_\nu) + v^\nu v_\nu = 1. \quad (A9)$$

Since $k^\nu u_\nu$ and $v^\nu v_\nu$ are known quantities, Eq. (A9) provides a third expression for $l^\nu u_\nu$. By rewriting the 1 on the right side of Eq. (A9) in terms of the initial values of the quantities on the left side, one can rewrite Eq. (A9) in the form of Eq. (A8). All three approaches have their uses. Equation (A4) is the covariant version of Eq. (11), and Eqs. (A5) and (A8) are the covariant versions of Eq. (14) for $u_{||}$ and its analog for γ .

Appendix B: Covariant Lagrangian for the Particle Motion

For a particle in an electromagnetic field the normalized motion¹⁹

$$S = -\int \left[\left(dx^\nu dx_\nu \right)^{1/2} + a^\nu dx_\nu \right]. \quad (\text{B1})$$

Traditionally, one parameterizes the particle motion in terms of the proper time τ , which is a Lorentz invariant. In this case

$$S = -\int \left[\left(d_\tau x^\nu d_\tau x_\nu \right)^{1/2} + a^\nu d_\tau x_\nu \right] d\tau. \quad (\text{B2})$$

By applying the Euler-Lagrange equations to the integrand of Eq. (B2), one finds that

$$d_\tau (d_\tau x_\mu + a_\mu) = d_\tau x^\nu \partial_\nu a_\mu, \quad (\text{B3})$$

in agreement with Eq. (8). Alternately, one can parameterize the particle motion by the phase $\phi = k^\nu x_\nu$, which is also a Lorentz invariant. In this case

$$S = -\int \left[\left(d_\phi x^\nu d_\phi x_\nu \right)^{1/2} + a^\nu d_\phi x_\nu \right] d\phi. \quad (\text{B4})$$

By using the decompositions (35) and (36) one can rewrite Eq. (B4) as

$$S = -\int \left[\left(d_\phi y^\nu d_\phi y_\nu + d_\phi \theta \right)^{1/2} + b^\nu d_\phi y_\nu + p d_\phi \theta / 2 + q / 2 \right] d\phi. \quad (\text{B5})$$

By applying the Euler-Lagrange equations to the integrand of Eq. (B5), one can show that

$$\begin{aligned} & \frac{d}{d\phi} \left[\frac{1}{\left(d_\phi y^\nu d_\phi y_\nu + d_\phi \theta \right)^{1/2}} \frac{dy_\mu}{d\phi} + b_\mu \right] \\ &= \frac{\partial b^\nu}{\partial y^\mu} \frac{dy_\nu}{d\phi} + \frac{1}{2} \left(\frac{\partial p}{\partial y^\mu} \frac{d\theta}{d\phi} + \frac{\partial q}{\partial y^\mu} \right), \end{aligned} \quad (\text{B6})$$

$$\begin{aligned} & \frac{d}{d\phi} \left[\frac{1}{\left(d_\phi y^\nu d_\phi y_\nu + d_\phi \theta \right)^{1/2}} + q \right] \\ &= 2 \frac{\partial b^\nu}{\partial \theta} \frac{dy_\nu}{d\phi} + \frac{\partial p}{\partial \theta} \frac{d\theta}{d\phi} + \frac{\partial q}{\partial \theta}, \end{aligned} \quad (\text{B7})$$

in agreement with Eqs. (37) and (38). One can reproduce Eq. (39) by multiplying Eq. (B6) by $-2d_\phi y^\mu$ and Eq. (B7) by $-d_\phi \theta$, and adding the resulting equations.

Appendix C: Evaluation of the Four-Potential

The left side of Eq. (34) contains the term $da_\mu/d\phi$, which must be evaluated at the position of the particle. In the section **Multiple Scale Analysis of the Particle Motion** we used Eqs. (42), (43), and (46) to make a guiding-center expansion of a_μ before we took the convective derivative. Specifically, we wrote

$$d_\phi a_\mu \approx [d_0 + \varepsilon d_1] \left[a_\mu^{(0)} + \varepsilon a_\mu^{(1)} \right], \quad (\text{C1})$$

where

$$a_\mu^{(0)} = \bar{a}_\mu \quad (\text{C2})$$

is the four-potential evaluated at the guiding center and

$$\bar{a}_\mu^{(1)} = y^{(0)\nu} \partial_\nu \bar{a}_\mu + \theta^{(0)} \partial_\theta \bar{a}_\mu \quad (\text{C3})$$

is the correction to the four-potential caused by the oscillation of the particle about the guiding center. Since the guiding-center coordinates $y^{(-1)}$ and $\theta^{(-1)}$ are functions of ϕ_1 by construction,

$$\frac{d\bar{a}_\mu}{d\phi_0} = \frac{\partial \bar{a}_\mu}{\partial \phi_0} \quad (\text{C4})$$

and

$$\frac{d\bar{a}_\mu}{d\phi_1} = \frac{dy^{(-1)\nu}}{d\phi_1} \frac{\partial \bar{a}_\mu}{\partial y^\nu} + \frac{\partial \bar{a}_\mu}{\partial \phi_1} + \frac{d\theta^{(-1)}}{d\phi_1} \frac{\partial \bar{a}_\mu}{\partial \theta}. \quad (\text{C5})$$

It follows from Eqs. (C1), (C4), and (C5) that

$$\left[\frac{da_\mu}{d\phi} \right]^{(0)} = \frac{\partial \bar{a}_\mu}{\partial \phi_0} \quad (C6)$$

and

$$\begin{aligned} \left[\frac{da_\mu}{d\phi} \right]^{(1)} &= \frac{dy^{(-1)v}}{d\phi_1} \frac{\partial \bar{a}_\mu}{\partial y^v} + \frac{\partial \bar{a}_\mu}{\partial \phi_1} + \frac{d\theta^{(-1)}}{d\phi_1} \frac{\partial \bar{a}_\mu}{\partial \theta} \\ &+ \frac{dy^{(0)v}}{d\phi_0} \frac{\partial \bar{a}_\mu}{\partial y^v} + y^{(0)v} \frac{\partial^2 \bar{a}_\mu}{\partial \phi_0 \partial y^v} \\ &+ \frac{d\theta^{(0)}}{d\phi_0} \frac{\partial \bar{a}_\mu}{\partial \theta} + \theta^{(0)} \frac{\partial^2 \bar{a}_\mu}{\partial \phi_0 \partial \theta}. \end{aligned} \quad (C7)$$

Alternately, one can write

$$\frac{da_\mu}{d\phi} = \frac{dy^v}{d\phi} \frac{\partial a_\mu}{\partial y^v} + \frac{\partial a_\mu}{\partial \phi} + \frac{d\theta}{d\phi} \frac{\partial a_\mu}{\partial \theta}, \quad (C8)$$

in which the guiding-center expansion is made *after* the partial derivatives are taken. Since the variation of a_μ with the position variables y^v and θ is slow,

$$\begin{aligned} \frac{da_\mu}{d\phi} &\approx \varepsilon \left[\frac{dy^{(0)v}}{d\phi_0} + \frac{dy^{(-1)v}}{d\phi_1} \right] \frac{\partial a_\mu}{\partial y^v} + \frac{\partial a_\mu}{\partial \phi_0} \\ &+ \varepsilon \frac{\partial a_\mu}{\partial \phi_1} + \varepsilon \left[\frac{d\theta^{(0)}}{d\phi_0} + \frac{d\theta^{(-1)}}{d\phi_1} \right] \frac{\partial a_\mu}{\partial \theta}. \end{aligned} \quad (C9)$$

The derivatives of the four-potentials appearing in the order ε terms can be approximated by their guiding-center values. The remaining term

$$\frac{\partial a_\mu}{\partial \phi_0} \approx \frac{\partial \bar{a}_\mu}{\partial \phi_0} + \varepsilon y^{(0)v} \frac{\partial^2 \bar{a}_\mu}{\partial y^v \partial \phi_0} + \varepsilon \theta^{(0)} \frac{\partial^2 \bar{a}_\mu}{\partial \theta \partial \phi_0}. \quad (C10)$$

Equations (C9) and (C10) are equivalent to Eqs. (C4) and (C5). This result shows that the guiding-center expansion discussed previously was made consistently. The expansion based on

Eq. (C1) is better because it facilitates the identification of combinations of terms that are oscillatory and, hence, do not affect the guiding-center motion.

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Subpicosecond Imaging System Based on Electro-Optic Effect

Ultrafast electro-optic (EO) sampling was first demonstrated in 1982¹ and has since become a valuable tool for testing optoelectronic and electronic devices and materials.² Conventional EO sampling of weak electric fields employs a tightly focused, pulsed-laser probe beam to measure electric-field-induced birefringence in an EO crystal; hence, it is referred to as "point" sampling.

Densely packed analog and digital devices make it necessary to probe many nodes simultaneously. Meyer and Mourou³ first demonstrated electric field mapping by scanning an area using a point sampler. Martin⁴ reviews the development of two-dimensional field measurement technologies including an automated scanning point sampler. Two groups studying photoconductive switches^{5,6} pioneered the use of EO imaging, by mapping the field strength with a detector array. Their work differs from the present in that their devices exhibited high fields and were adequately described with 200-ps temporal resolution.

An EO sampling system capable of imaging the voltage distribution over a rectangular region is described. It is comparable to an ultrafast sampling oscilloscope having more than 180,000 channels. This analysis focuses on techniques that take advantage of the speed and convenience of a charge-coupled-device (CCD) sensor while overcoming its limited dynamic range.

System Descriptions

EO sampling requires a pulsed (or gated) laser source to probe the response of the device to the applied transient. Our lab uses a mode-locked Coherent Mira 900 Ti:sapphire laser. It produces a 76-MHz train of linearly polarized, ≈ 100 -fs FWHM pulses, tuned to ≈ 800 nm. Devices tested in our lab generally include a photoconductive switch that is excited with a fraction of the pulsed beam, thus triggering the measurement and eliminating electrical jitter.

In a point sampler,² the EO crystal may be either the device substrate (e.g., GaAs devices) or on an external probe. A linearly polarized optical probe pulse enters the crystal through the first surface. In transmissive sampling, the probe is transmitted at the second surface after a single pass, whereas in reflective sampling, it is reflected, passing through the crystal a second time. The beam exits the crystal and is passed through a compensator or wave plate to introduce a static polarization bias. The bias is adjusted so that in the absence of an electric field, the probe is circularly polarized at the input of an analyzer, thus giving maximum sensitivity and linearity when a field is applied. The analyzer separates the beam into orthogonal polarization components, which are measured by a pair of detectors connected to a differential lock-in amplifier. Signal-to-noise improvements are obtained when the signal is modulated at frequencies approaching the laser $1/f$ noise floor.

Figure 69.31 depicts the imaging system hardware. Reflective sampling was chosen because it doubles system sensitivity, although transmissive sampling is also possible. The laser source is directed through a high-speed modulator followed by a variable-intensity beam splitter consisting of a half-wave plate and polarizing beam splitter. The horizontally polarized "probe" beam is directed back through the polarizer, then into a spatial filter and beam expander. The vertically polarized "excitation" beam passes through a variable-length optical delay and into a fiber coupler.

The probe beam is split into two beams in a small, rigid interferometer. The device-under-test (DUT) is mounted in the device "leg" of the interferometer, and a mirror is installed in the reference "leg." The beams pass through a polarizing filter and relay lens to create an interference pattern at the camera. The beam splitter in the interferometer is an uncoated, ≈ 3 -mm-thick, BK-7 wedged window. The first surface of the window is aligned at Brewster's angle to eliminate multiple reflections and maximize transmitted intensity. The reference

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Figure 69.31

Imaging system hardware: (A) Ti:sapphire laser, (B) high-frequency modulator, (C) half-wave plate, (D) polarizing beam splitter, (E) optical delay stage, (F) excitation beam fiber coupler, (G) spatial filter and probe beam expander, (H) wedged beam splitter, (I) reference mirror on piezoelectric actuator, (J) EO crystal on DUT, (K) polarizing filter, (L) relay lens and aperture, and (M) CCD camera.

mirror is mounted on a piezoelectric actuator, which is used to modulate the length of the reference leg. The DUT is mounted on a stationary structure. Each leg has adjustments for static alignment.

As in point sampling, the electric fields on the DUT are measured by using the linear EO, or Pockels, effect. A propagating electrical transient is launched on the DUT when an optical excitation pulse is applied to a biased, photoconductive switch. An EO crystal having a high-reflectivity (HR) coating on one side covers the region of interest with the coating in intimate contact with the DUT. "Fringing" E -fields caused by the propagating transient couple into the crystal to produce a temporally and spatially variant refractive index.

The EO-induced index perturbation in the crystal alters the phase of the linearly polarized optical probe as it traverses the device leg of the interferometer. When recombined with an unperturbed reference beam, an intensity pattern results that corresponds to phase differences between the two legs of the interferometer. If the reference beam is static, then changes in

intensity at each point can be attributed to spatial phase variations in the crystal induced by the EO effect.

A video camera (DVC Corp., DVC-0A) having a low-noise, frame-transfer charge-coupled device (CCD) (Texas Instruments, TC-245) records the intensity pattern created by the interferometer. The analog camera output is digitized by a frame grabber (Matrox Corp., Pulsar) and stored on a personal computer (Pentium 133-MHz, PCI bus). Timing control for modulation uses custom-built electronics (see **Modulation** section).

Spatial resolution of the system is determined by the active image area and number of discrete pixels in the image sensor. Image (de)magnification can be adjusted by altering the position of the relay lens and camera.

The CCD has 755 ($8.5\text{-}\mu\text{m}$) pixels horizontally (H) and 242 ($19.75\text{-}\mu\text{m}$) pixels vertically (V) for an active area of 6.4 mm (H) $\times$ 4.8 mm (V). Typical magnification is 4:1, giving a measurement area of 1.6 mm (H) $\times$ 1.2 mm (V). The resulting spatial resolution is $2.13\text{ }\mu\text{m}$ (H) $\times$ $4.9\text{ }\mu\text{m}$ (V), which is comparable to point sampling. If desired, cylindrical lenses or prisms could be used to correct the pixel aspect ratio.

It is possible to increase optical magnification to 8:1, then digitally average 2×2 -pixel cells to obtain 4:1 effective magnification. This would reduce noise by $1/2$; however, it may prove disadvantageous since more photons from the excitation source will be collected by the sensor [see also **Interferometer Operation** section].

Important distinctions exist between the imager and scanning point samplers. The imaged nodes must lie within a finite rectangular region, whereas a scanning system can probe random points over an extended area. Furthermore, the imager measures all nodes simultaneously, whereas a scanning sampler probes one node at a time.

Electro-Optic Interferometer

We present the reasons for choosing an interferometer and discuss its operation. We begin by mathematically describing the EO effect, and the relationship between the voltages present on the DUT, fringing fields coupled into the crystal, and resulting phase delay experienced by the optical probe. We then use this information to estimate the temporal resolution of the system. Following this discussion, we analyze the design in Fig. 69.31 to estimate the expected system sensitivity.

The refractive index in an EO crystal is altered in the presence of an electric field. The perturbed index n' is dependent on the field-free index n , field strength E , and Pockels coefficients r . By applying the techniques of Ref. 7 to x -cut LiTaO₃ (<3 m> point group), a material commonly used for EO sampling, we find (neglecting terms quadratic in field strength E_y)

$$n'_y = n_y - \frac{n_y^3}{2} (r_{22}E_y + r_{13}E_z) = n_y + \Delta n_y, \quad (1a)$$

$$n'_z = n_z - \frac{n_z^3}{2} (r_{33}E_z) = n_z + \Delta n_z. \quad (1b)$$

Numeric subscripts are indices of the tensor elements, and y, z subscripts are direction vectors in crystalline coordinates; z is parallel to the optic axis. These equations show that the refractive index along y is influenced by the electric fringing fields directed along both y and z , whereas the index along z is influenced only by fringing fields along z . It is also evident that the optical probe polarization must be aligned to measure the desired refractive index perturbation, while the optic axis of the crystal must be aligned on the DUT such that the fringing fields of interest maximize the index perturbation.

If we substitute values for LiTaO₃⁷ into Eqs. (1a) and (1b), we find that $\Delta n_z \approx 4.4 \Delta n_y$, and the contribution from E_y is negligible. In point sampling, it is common (and convenient) to measure the induced birefringence, which is the difference in index perturbation along z and y , or

$$\Delta n_{zy} = \Delta n_z - \Delta n_y \approx \Delta n_z / 1.3.$$

Since the refractive index change along z is greater than that along y and greater than the induced birefringence, system sensitivity will be maximized by measuring Δn_z . An interferometer was chosen for this purpose. [Note: EO materials from other point groups (e.g., ZnTe, <43 m>) have greater sensitivity when the induced birefringence is measured.]

Having determined that we wish to measure the refractive index perturbation using an interferometer, we must consider how it will be used. An interferometer is sensitive to phase delays imposed on a propagating optical wavefront, which in our case is the probe beam. As an optical beam traverses a dielectric material, it suffers a phase delay $\Delta\Gamma$, determined by the refractive index n , wavelength λ , and material thickness X :

$$\Delta\Gamma = \int_0^X \frac{2\pi}{\lambda} n(x) dx. \quad (2)$$

We showed above that the refractive index was dependent upon the electric fringing field, but we must also consider that the fringing field is not uniform throughout the thickness of the material. As a result, the refractive index is a function of depth x , determined by the penetration depth of the fringing field into the crystal.

Substituting Eq. (1b) into Eq. (2), we obtain a static phase delay component Γ_0 (independent of E fields):

$$\Gamma_0 = \frac{2\pi}{\lambda} n_z X, \quad (3a)$$

and a dynamic phase delay attributed to the EO effect $\Delta\Gamma_{EO}$. The interferometer measures $\Delta\Gamma_{EO}$, given by

$$\Delta\Gamma_{EO} = \frac{\pi}{\lambda} n_z^3 r_{33} \int_0^X E_z(x) dx. \quad (3b)$$

The E -field distribution within the crystal depends upon the test structure. For this example, consider a coplanar waveguide on which we wish to probe the E -field at the center of the gap g . In general, if a superstrate having the same relative dielectric constant as the substrate (ϵ_{sub}) is placed on a coplanar structure, we would expect the fringing fields in the superstrate to be confined to a depth comparable to the gap separating device features. When the superstrate is the EO crystal (dielectric = ϵ_{EO}), the depth of the fringing field, g' , is dependent upon the ratio of the two dielectric constants; the confinement depth becomes $g' \cong g \epsilon_{\text{sub}} / \epsilon_{EO}$. The field strength decreases rapidly inside the crystal, so we approximate the integral with the product $E_z^{\text{surface}} g'$, where E_z^{surface} is the transverse E -field magnitude at the surface of the crystal. We then obtain

$$\Delta\Gamma_{EO} \cong \frac{\pi}{\lambda} n_z^3 r_{33} \left(\frac{\epsilon_{\text{sub}} g E_z^{\text{surface}}}{\epsilon_{EO}} \right), \quad (3c)$$

or in words, the measured phase change at any point is proportional to the E -field at that point. The voltage on the gap V^{gap} is the product of the gap and the E -field:

$$V^{\text{gap}} = g E_z^{\text{surface}}. \quad (3d)$$

Note that for a given value of $\Delta\Gamma_{EO}$, the voltage is independent of the gap, whereas the field depends upon the gap. This can be understood by considering that a device having a larger gap has deeper fringing-field penetration in the crystal. The fields have a longer interaction length with the probe; hence, the field required to produce a given phase change is reduced.

Temporal resolution of the system is determined by the largest of (1) response time of the EO material, or (2) probe-pulse duration convolved with the fringing fields profile; this convolution is approximately equal to the sum of the pulse FWHM and the time of flight t_{fl} of an infinitely short pulse through the fringing fields. The EO response is limited by phonon resonance and for LiTaO₃ is of the order of 10^{-14} s.⁸ The probe pulse is ≈ 100 -fs FWHM and can be reduced to ≈ 50 -fs FWHM using a pulse compressor. The optical path length through the fringing fields is $pl = (2ng\epsilon_{sub}/\epsilon_{EO})$. Time of flight $t_{fl} = pl/c$, with c = speed of light in vacuum. For a coplanar waveguide fabricated on silicon ($\epsilon_{sub} = 11.9$), having $g = 10$ μ m, and LiTaO₃ ($n \approx 2.2$, $\epsilon_{EO} = 43$), we find $t_{fl} \approx 40$ fs. From these values, we expect (temporal resolution) $\approx$ (pulse FWHM + t_{fl}) = 140 fs, well below 1 ps.

Interferometer Operation

Now that we have described how Pockels effect alters the phase of an optical probe beam, we discuss the interferometer in detail. We begin with its intensity transfer function and discuss the ideal case. We then consider factors that cause deviations from ideal that reduce system sensitivity, and estimate their magnitude. Finally, we consider how to optimize system sensitivity given these constraints.

The normalized intensity measured by the detector, $I_d = I_{out}/I_{ref}$, is the ratio of the output intensity from the interferometer to the intensity present in the reference leg:

$$I_d \equiv \frac{I_{out}}{I_{ref}} = (1 + \alpha) + 2\sqrt{\alpha} \cos(\delta) + b. \quad (4)$$

I_d depends on $\alpha = I_{DUT}/I_{ref}$, the normalized intensity in the device leg, the phase difference δ between the E -field of the optical probe in each leg, and normalized background illumination b . Using Eqs. (3a) and (3c) to expand δ , we get

$$\delta = 2(\Gamma_0 + \Delta\Gamma_{EO}) \equiv 2\Gamma, \quad (5)$$

where Γ_0 was redefined to include both the static phase difference governed by the differing lengths of the interferom-

eter legs, as well as the static phase delay of the EO crystal. The factor of 2 results from using reflective sampling. The probe passes through the fringing field two times, accumulating twice the phase delay.

In an ideal interferometer $\alpha = 1$ and $b = 0$, and Eq. (4) reduces to

$$I_d \propto \cos^2(\Gamma), \quad (6)$$

which is also the intensity transfer function used to describe point sampling. As a result, all modulation and detection principles described herein apply equally to a system such as that in Ref. 2, wherein a variable retarder is used in place of the quarter-wave plate or optical compensator. The variable retarder would take on the modulation function of the piezoelectric actuator, as discussed in the section entitled **Modulation**.

CCD's have a finite electron well-capacity, and consequently, sensitivity will be greatest when the ratio q/Q is maximized, where q = number of electrons attributed to the EO signal and Q = total number of electrons. Assuming that the number of electrons in each pixel is linearly proportional to the incident radiant flux,

$$Q \propto I_d, \quad (7a)$$

and

$$q \propto \Delta I_{EO} \equiv I_d - I_d \Big|_{\Delta\Gamma_{EO}=0}, \quad (7b)$$

where ΔI_{EO} is the intensity contribution from the EO effect alone. Combining Eqs. (4), (5), (7a), and (7b) yields

$$\frac{q}{Q} = \frac{\cos 2(\Gamma_0 + \Delta\Gamma_{EO}) - \cos 2(\Gamma_0)}{p + \cos 2(\Gamma_0 + \Delta\Gamma_{EO})}, \quad (8)$$

where

$$p \equiv \frac{1 + \alpha + b}{2\sqrt{\alpha}}. \quad (9)$$

In the small-signal limit, Eq. (8) becomes

$$\frac{q}{Q} \Big|_{\text{Lim } \Delta\Gamma_{EO} \rightarrow 0} = f(p, \Gamma_0) \Delta\Gamma_{EO}, \quad (10a)$$

where

$$f(p, \Gamma_0) \equiv \left[\frac{2 \sin(2\Gamma_0)}{p + \cos(2\Gamma_0)} \right]. \quad (10b)$$

Equation (10a) describes the fraction of electrons in each pixel attributed to the EO effect. We now consider how to use this information to optimize the system sensitivity.

Figure 69.32(a) presents a plot of Eq. (9), and Fig. 69.32(b) shows $f(p, \Gamma_0)$ defined in Eq. (10b) for the nonideal case $p = 1.05$. From Fig. 69.32(b), we see that $f(p, \Gamma_0)$ has two points for which the amplitude is a maximum. We wish to find Γ_0^{opt} —bias points for which this function is optimized. To do so, we take the derivative of $f(p, \Gamma_0)$ with respect to Γ_0 , equate to zero, and solve

$$\Gamma_0^{\text{opt}}(p) = \frac{1}{2} \cos^{-1}(-p^{-1}). \quad (11)$$

Equation (11), plotted in Fig. 69.32(c), shows a distinct Γ_0 that maximizes q/Q . Therefore, we wish to optically bias the interferometer at this point, about which the small EO signal is superimposed. We note that for an ideal interferometer $p = 1$, giving $\Gamma_0^{\text{opt}} = \pi/2$ where the derivative of Eq. (6) is zero. This conclusion is very different from wide-bandwidth detectors used in point sampling that achieve maximum sensitivity when $\Gamma_0 = \pi/4$, where the derivative of Eq. (6) is maximized, as explained in Ref. 2. In the general case of a nonideal interfer-

ometer, two solutions exist [as originally expected from Fig. 69.32(b)], one on either side of $\pi/2$.

Substituting Eq. (11) into Eq. (10b), we find

$$f(p, \Gamma_0) \Big|_{\Gamma_0 = \Gamma_0^{\text{opt}}} = \frac{2}{p\sqrt{1-p^{-2}}} \equiv f^{\text{opt}}(p). \quad (12)$$

The factor $f^{\text{opt}}(p)$ is also plotted in Fig. 69.32(c). It has the greatest value for an ideal interferometer and decreases as we depart from ideal.

Until now, we have ignored sources of optical phase-front distortions to the probe beam. Each optical component has a finite surface accuracy and refractive index inhomogeneities. These inaccuracies are stationary in time, and spatially random, so the cumulative error is the rms combination of all components. These errors will make it impossible to achieve optimum system sensitivity at every pixel simultaneously, but two observations can be made:

1. As the region-of-interest (ROI) is decreased (within diffraction limits), the magnitude of phase distortions will decrease, and
2. As a system, sensitivity will be maximized when the average optical bias point in the ROI corresponds to the optimum bias conditions.

Figure 69.32

(a) Parameter $p(\alpha, b) = 1$ for an ideal interferometer and increases as (α, b) depart from ideal. (b) Sensitivity factor $f(p, \Gamma_0)$ for an arbitrary nonideal interferometer ($p = 1.05$) has two peaks, corresponding to optimum operating bias points Γ_0^{opt} . (c) Optimum optical bias points Γ_0^{opt} and sensitivity factor $f^{\text{opt}}(p)$ can be determined knowing p .

We have shown that there exists an optimum bias point about which we must modulate our signal. To determine this bias point, we must understand the origins of α and b and estimate their magnitudes.

In an ideal interferometer, the beam splitter would be infinitely thin, so that reflections occur only at one surface. Pellicle beam splitters are thin but are subject to acoustic and mechanical vibrations, making them unsuitable for this application. A thick beam splitter is more stable, but reflections from the second surface must be eliminated. Coated optics are an option, but we chose to eliminate unwanted reflections by using a wedged window. The beam is incident at Brewster's angle at the first surface such that reflectivity of p -polarized radiation is zero. The beam splitter then behaves ideally, i.e., $\alpha = 1.0$.

Background illumination, factor b in Eq. (4), is radiation collected by the detector that does not contribute to the desired signal. Fresnel reflections occur at each dielectric interface (window or lens), as in Fig. 69.33(a). Each transmitted beam is the superposition of many reflections. Beams that experience multiple reflections will be delayed more than the duration of the probe pulse, so will not interfere. To estimate b_{lens} , we compare the intensity of transmitted light delayed by more than $n_1 t$ to that delayed by exactly $n_1 t$. For a system of M windows and lenses,

$$b_{\text{lens}} = (1 - R^2)^{-M} - 1. \quad (13)$$

Since system sensitivity decreases with increasing b , it is advantageous to minimize the reflection coefficient R at each optical element by using coated optics. A conservative estimate for the system shown in Fig. 69.31 (not all optics shown) having seven uncoated BK-7 windows ($R = 0.04$) gives $b_{\text{lens}} \approx 1.2\%$.

Fresnel reflections occur also at the surface of the crystal. Most EO materials used for sampling have a large refractive index, giving large reflections. The following expression for background contributions from the crystal, b_{EO} , is evident from Fig. 69.33(b):

$$b_{\text{EO}} = 1 - (1 - R_{\text{EO}})^{-2}. \quad (14)$$

Reflections from the top surface of the crystal are potentially the most detrimental to system performance. Uncoated

LiTaO₃ has $R_{\text{EO}} \approx 14\%$, making $b_{\text{EO}} \approx 26\%$; anti-reflection (AR) coatings, which make $R_{\text{EO}} \approx 0.03$, yield $b_{\text{EO}} \approx 6\%$. R_{EO} also effectively reduces α to α' by

$$\alpha' = \frac{I_{\text{DUT}}}{I_{\text{ref}}} (1 - R_{\text{EO}})^2 = \alpha (1 - R_{\text{EO}})^2, \quad (15)$$

since only a fraction of the incident pulse makes exactly one round-trip through the crystal.

Uncoated LiTaO₃ makes $\alpha' = 0.74 \alpha$, whereas coated LiTaO₃ produces $\alpha' = 0.94 \alpha$. To minimize these detrimental effects, the crystal requires a nearly perfect AR coating on the first surface ($R_{\text{EO}} = 0$) and a perfect HR coating ($R = 1$) on the second surface.

The final source of background is the light reflected by the DUT from the fiber-coupled beam used to trigger the photoconductive switch. A conservative estimate assumes that the fiber is positioned at the DUT and pointed directly toward the interferometer beam splitter [Fig. 69.33(c)]. The results of this analysis will be at least an order of magnitude too large because the estimate neglects the following facts:

- (1) the fiber is directed toward the DUT and will shadow reflected light;
- (2) the DUT will absorb incident photons;
- (3) the polarizing filter will attenuate reflected (scattered) light that is depolarized; and
- (4) the photoconductive switch may be located outside the image area.

From the above argument, the results of the following simplified analysis will be reduced by a factor of 10.

The divergence angle of the beam is determined by the fiber diameter and wavelength. A fraction of the light is reflected off the beam splitter toward the camera. The distance between the relay lens and DUT is determined by the desired magnification m and lens focal length f . The lens has a finite aperture and collects only a fraction of the diverging beam from the fiber, b_{fiber} . Assuming a gaussian beam from the fiber tip, this simplified approach yields

$$b_{\text{fiber}} \approx \frac{I_{\text{ex}} R_{\text{bs}}}{I_{\text{ref}}} \text{erf} \left(\frac{D}{D'} \right), \quad (16a)$$

Figure 69.33

Factors that degrade interferometer performance: (a) multiple reflections from windows and lenses [R = intensity reflection coefficient, T = intensity transmission coefficient = $(1-R)$], (b) Fresnel reflections at the surface of the EO crystal, and (c) light escaping from excitation fiber.

where

$$\frac{D}{D'} = \frac{\lambda D m}{2df(m+1)} \quad (16b)$$

and where I_{ex} = intensity of the excitation pulse (≈ 1 mW), $I_{ref} \approx 10$ μ W (for pixel saturation), R_{bs} = beam-splitter reflectivity, $erf()$ is the error function, D = lens-aperture diameter, $D' = 1/e$ beam diameter at the lens, λ = wavelength, and d = fiber-core diameter.

Clearly, R_{bs} and m should be minimized, and f should be large. From Eqs. (16a) and (16b), typical operating condi-

tions give $b_{fiber} \approx 13\%$, which we reduce to 1.3%, as discussed above.

We have considered several factors that contribute to the nonideal terms α and b in the interferometer transfer function. It is essential to use precision optics and minimize front-surface reflections from the EO crystal to prevent system degradation. Proper adjustment of the excitation beam intensity and fiber placement will limit background contributions from the excitation source. Finally, coated optics will reduce multiple reflections from other system optics. For a well-designed system having an AR-coated crystal, we obtain $b = b_{EO} + b_{fiber} + b_{lens} \approx 0.06 + 0.012 + 0.013 = 0.085$, and $\alpha = 0.94$, thus making $p \approx 1.044$. This value for p will be used in the remaining discussion.

System Linearity and Sensitivity

Linearity of the measurement system can be derived from the ratio of Eqs. (10a) and (8), where Γ_0 is replaced with Γ_0^{opt} , and $\Delta\Gamma_{\text{EO}}$ is a small-signal perturbation about Γ_0^{opt} . Evaluating linearity at $p = 1.044$, one can show that the measured response is linear within $\pm 5\%$ for $|\Delta\Gamma_{\text{EO}}| < 0.015$ rad; this is more than adequate for expected signals.

We have obtained an expression for the optimum sensitivity factor $f^{\text{opt}}(p)$ and numerical estimates of the parameter p . The next step is to determine the measurement resolution of the system given this information. First, we determine the system dynamic range (DR) and minimum resolvable phase change, then the voltage and E -field needed to produce this phase change.

From Eq. (10), we can determine $\text{DR}_{\text{signal}}$ if we assume that the pixel is nearly saturated so that $Q \approx Q_{\text{well}}$. The CCD has an electronic noise-equivalent signal $q_{\text{eq}} = 30$ electrons, shot noise $q_{\text{shot}} = 40$ electrons, and well capacity $Q_{\text{well}} = 80 \times 10^3$ electrons.⁹ Setting $q_{\text{noise}} = (q_{\text{eq}}^2 + q_{\text{shot}}^2)^{1/2} = 50$ electrons, we find

$$\text{DR}_{\text{signal}} = 20 \cdot \log \left(\frac{Q_{\text{well}}}{q_{\text{noise}}} \frac{q}{Q} \right) \bigg|_{\frac{q}{Q} = f^{\text{opt}}(p) \Delta\Gamma_{\text{EO}}} \quad (17)$$

For $\Delta\Gamma_{\text{EO}} = \pm 0.015$ rad (the limit of "linear" range) and $p = 1.044$, we find $f^{\text{opt}}(p) = 6.25$, and $\text{DR}_{\text{signal}} = 43$ dB.

The minimum detectable signal $\Delta\Gamma_{\text{EO}}^{\text{min}}$ is that which makes $q/Q = q_{\text{noise}}/Q_{\text{well}}$:

$$\Delta\Gamma_{\text{EO}}^{\text{min}} = \left(\frac{q_{\text{noise}}}{Q_{\text{well}} f^{\text{opt}}(p)} \right), \quad (18)$$

which gives $\Delta\Gamma_{\text{EO}}^{\text{min}} = 100 \mu\text{rad}$, corresponding to $\lambda/6 \times 10^4$ resolution.

We relate $\Delta\Gamma_{\text{EO}}^{\text{min}}$ to the voltage necessary to produce it, using Eqs. (3c) and (3d):

$$\Delta\Gamma_{\text{EO}}^{\text{min}} \cong \frac{\pi}{\lambda} n_z^3 r_{33} \left(\frac{\epsilon_{\text{sub}}}{\epsilon_{\text{EO}}} \right) (V_{\text{gap}, \text{min}}). \quad (19)$$

The minimum detectable voltage $V_{\text{gap}, \text{min}}$ is constant for any (coplanar) gap geometry. $E_z^{\text{surface}, \text{min}}$ is the minimum field,

which, if present at the surface of the crystal, could be resolved by the system:

$$E_z^{\text{surface}, \text{min}} = \frac{V_{\text{gap}, \text{min}}}{g}. \quad (20)$$

When testing a device fabricated on silicon ($\epsilon_{\text{sub}} = 11.9$) using LiTaO_3 ($\epsilon_{\text{EO}} = 43$, $r_{33} = 33$ pm/V, $n_z = n_e = 2.180$),⁷ and $\lambda = 800$ nm, we find $\Delta\Gamma_{\text{EO}}^{\text{min}} = 3.7 \times 10^{-4} V_{\text{gap}, \text{min}}$. By equating $\Delta\Gamma_{\text{EO}}^{\text{min}}$ to $100 \mu\text{rad}$, $V_{\text{gap}, \text{min}} = 270$ mV, which corresponds to 27 kV/m on a $10\text{-}\mu\text{m}$ gap.

This sensitivity is well suited to measurement of microwave devices and complex transmission line structures. Several enhancements can be made to improve suitability for digital applications. A nonlinear organic salt known as DAST has $\epsilon_{\text{EO}} = 7.0$, Pockels coefficient $r_{11} = 160$ pm/V, and $n = 2.460$.¹⁰ From Eq. (19), this would increase sensitivity by a factor of 43. Cooling the sensor reduces shot noise so that $q_{\text{noise}} \cong q_{\text{eq}}$, thus by Eq. (18), increasing sensitivity by a factor of 1.6. In combination, these produce $V_{\text{gap}, \text{min}} \cong 4$ mV, and $E_z^{\text{surface}, \text{min}} \cong 400$ V/m.

Modulation

Having discussed the attributes of an integrating detector, we now show how the signal is modulated about the desired operating bias point. We first discuss how the signal is modulated in each image and the timing required. We then discuss the characteristics of a frame transfer sensor and how to use these characteristics to our advantage.

Figure 69.34(a) is an expanded view of Eq. (6), about $\pi/2$. In the absence of an electric field, points A and B have equal intensity when the optical bias is adjusted to $\pm\Gamma_0^{\text{opt}}$ by displacing the reference mirror. When an E -field is present on the DUT, the resulting EO phase shift is added to the optical bias. This causes the intensity at point A to increase to C, while that at point B decreases to D. Analysis of the data is achieved by subtracting field D from B, and C from A.

The data-acquisition system is synchronized to the camera's pixel and field clocks and acquires images with the timing shown in Fig. 69.34(b). The electrical bias on the device's photoconductive switch is synchronously modulated at the 30-Hz camera field clock frequency, thus decreasing $1/f$ noise (both laser and mechanical vibrations of the interferometer). The reference mirror position is modulated at $1/2$ the bias frequency. A trigger pulse generated on the mutual rising edge of bias and actuator signals triggers the digitization of

Figure 69.34

Modulation: (a) expansion of Eq. (6) about an interferometric null showing EO modulation about $\pm\Gamma_0^{opt}$ bias points, (b) modulation and image capture timing diagram.

four consecutive fields, corresponding to points C, A, D, B in Fig. 69.34(a).

The frame transfer CCD has two discrete sensor regions: an active-pixel site and a storage site of equal size. Each field is acquired over a 1/60-s integration period. An advantage of using a detector with a 1/60-s integration period is that 60-Hz electrical noise will average to zero. During integration, the active pixels integrate charge proportional to photon flux, while electrons in the storage site are clocked to the output amplifiers. During frame transfer, charges in the active pixels are transferred vertically via "bucket brigade" into the storage site. Charge transfer causes slight smearing due to transfer inefficiency, and distortion occurs for charge packets that are transferred through brightly illuminated pixels.

A high-speed modulator "gates" the laser "on" immediately before and after alternate frame transfer cycles, and "off" at all other times. This eliminates charge smearing during frame transfer and reduces the effective laser- and vibration-noise bandwidth significantly. The limiting speed for this modulation is governed by the frame transfer period (1.27 ms for our camera). If the laser is gated "on" for ~100 pulses, the effective modulation frequency would be ~750 Hz.

Summary

We have described and analyzed an ultrafast EO imaging system that uses an interferometer and CCD detector to map 2-D electric fields on an optoelectronic device. It is comparable to an ultrafast sampling oscilloscope having more than 180,000 channels. Limitations caused by using an integrating detector are reviewed, and optimum operating conditions are identified. Techniques are presented that allow modulation of the signals at 750 Hz, which will reduce sensitivity to laser and mechanical $1/f$ noise. System sensitivity in the absence of laser noise is estimated to be 270 mV, corresponding to 27 kV/m for a 10- μ m coplanar structure. These values make the system well suited for testing microwave devices. Sensor cooling and the use of alternative EO materials should improve sensitivity by factors of 1.6 and 43, respectively, making the minimum resolvable voltage 4 mV. The system would then be easily capable of digital (e.g., CMOS) circuit evaluation.

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Nuclear Diagnostics for High-Density Implosions

The measurement of fuel compression, specifically the density-radius product (ρR), is of fundamental importance in analyzing ICF implosions. To probe the large values of ρR anticipated for OMEGA (and NIF) experiments, nuclear diagnostics must be used.

Three ρR diagnostics are being developed for OMEGA. They rely on detecting the following nuclear particles: (1) energetic D and T ions (knock-ons) produced by colliding 14-MeV neutrons;¹ (2) elastically scattered DT neutrons;^{2,3} and (3) tertiary DT neutrons.^{2,4} The starting point for generating each of these particles is the 14-MeV neutron produced in DT fusion, i.e.,

A small percentage of these neutrons (typically less than 0.01% for ICF conditions) will scatter elastically from D or T ions in the fuel (a prime indicates a scattered particle):

where D' will be produced in a continuous spectrum from 0 to 12.5 MeV and T' in the range of 0 to 10.6 MeV, with the remainder of the 14.1-MeV energy going to the scattered neutron. The number of such scattering events per DT-fusion neutron is directly proportional to the ρR of the fuel. This forms the basis for two of the nuclear diagnostics under development—knock-ons and elastically scattered neutrons: (1) the techniques for detecting the D and T knock-ons that were developed on the 24-beam OMEGA system will be extended to accommodate the higher values of ρR expected in future, near-term experiments (total target $\rho R \leq 200 \text{ mg/cm}^2$); and (2) techniques for detecting the elastically scattered neutrons and separating that signal from the expected neutron background will be investigated.

The third diagnostic involves reactions with the knock-on D and T ions. As these ions pass through the fuel, there is a

small probability that they will undergo an in-flight fusion reaction with one of the thermal fuel ions:

and similarly for the knock-on tritons, where n'' indicates a neutron scattered by an already scattered D. Because the D and T knock-ons are charged particles, they will be slowing down as they fuse. This can complicate the interpretation of the diagnostic signal, as the rate of slowing down depends on the temperature in the fuel as well as the density. The number of these high-energy neutrons from the tertiary reaction actually measures the product (ρR) ($\rho R'$), where R' is either the radius of the target or the range of the D or T knock-on ions, whichever is smaller. The first factor of ρR comes from the production of the knock-ons. Thus, for small ρR (OMEGA), the tertiary yield varies as $(\rho R)^2$, and for large ρR (NIF), it is proportional to ρR .

This article presents more details about these three diagnostics, together with comments on where further development is necessary. Two other nuclear diagnostics are under consideration but will not be discussed here, because they involve the additional complication of adding He^3 to the target. These two diagnostics use the D- He^3 proton: the energy loss of the 14.7-MeV proton from thermal D- He^3 fusion can measure the target ρR up to several hundred mg/cm^2 ,¹ and tertiary protons from in-flight D- He^3 fusion, using knock-on D or He^3 ions, can measure the ρR of the He^3 in the target.⁵ The detection of these protons from different directions can give information about gross asymmetries of the compressed core.

Knock-Ons

The knock-on diagnostic was developed at LLE for the 24-beam OMEGA laser and used to diagnose implosions with a fuel ρR up to $\sim 30 \text{ mg/cm}^2$ and with a comparable $\rho \Delta R$ for the glass shell around the fuel.⁶ For values of fuel-plus-shell ρR in excess of 100 mg/cm^2 , the knock-on spectrum becomes

substantially distorted due to significant slowing down of the D and T ions as they pass through the target; this can introduce some uncertainty in interpretation of the diagnostic signal. It is important to spectrally resolve the high-energy portion of this charged-particle signal, to separate the knock-ons from other sources of energetic ions, and to determine how the knock-on spectrum has been modified by slowing down in the target. Gross spectral resolution was obtained on the 24-beam OMEGA laser using stacked track detectors.^{1,6} Higher resolution will be obtained in future OMEGA experiments using a charged-particle spectrometer. This will extend the applicability of the knock-on diagnostic to meet the needs of the OMEGA experimental program for the first few years of operation.

To illustrate how the diagnostic can be used, Fig. 69.35 shows the calculated spectrum of knock-on deuterons for the simple model of a hot DT core surrounded by a "cold" (0.5-keV), denser CH shell. A characteristic feature of the spectrum is a peak at high energies produced by the forward-peaked cross section for scattering with 14-MeV neutrons. (This peak contains ~16% of the deuteron knock-ons.) Figure 69.35 shows how the peak is shifted to lower energies due to increased slowing down as the $\rho\Delta R$ of the plastic shell

increases. [The energy loss in the fuel is relatively small in this example because the fuel temperature is high and the ρR is small (40 mg/cm²).] The peak changes position and shape but remains well defined. By spectrally resolving this peak, it is possible to obtain simultaneously two important pieces of information about the compressed core: the position of the peak determines the $\rho\Delta R$ of the plastic shell, and the number of knock-on deuterons in the peak determines the ρR of the fuel. The fuel ρR is related to the number of deuterons in the peak (N_D) and the DT neutron yield (Y) by the following relation:

$$\rho R = 83 N_D / Y \quad \text{g/cm}^2 \quad (4)$$

for equimolar DT. A similar relation is available for knock-on tritons, but the more energetic deuterons can be used to diagnose higher values of ρR .

To facilitate measurement of the knock-on spectrum, a charged-particle spectrometer⁷ is being developed at MIT for OMEGA experiments. The spectrometer uses a 7.5-kG magnet to momentum select the particles and deflect them from the straight-line path followed by neutrons and x rays. The charged-particle paths are determined from trajectory calculations. The magnet has recently been tested at MIT using protons with energies from 1 to 15 MeV. As depicted in Fig. 69.36, the particles are deflected and then impinge onto a detector plane where they are intercepted by a combination of charged-coupled devices (CCD's) and CR-39 plastic track detectors. As demonstrated in recent experimental studies at MIT, the CCD's act as high-resolution energy detectors. Through the combination of magnetic momentum selection and the energy determination of the detectors, either CCD or track, the energy and identity of each particle will be uniquely specified. From these data, the spectra of all charged particles will be constructed, and vital information about the core conditions and dynamics will be measured.

This spectrometer also forms the basis of a joint proposal between MIT, LLE, and LLNL for a NIF diagnostic that will measure both the implosion symmetry and the core ρR .⁵ This makes use of very energetic tertiary protons (~31 MeV) that can easily penetrate all plasmas envisaged for the NIF. For example, the range of these protons is ~3 g/cm², whereas the ρR expected for a typical NIF capsule is ~1 g/cm². Even for some implosion scenarios simulated for OMEGA,⁵ situations have been encountered where these tertiary protons could prove particularly useful for determining core conditions, beyond the range of applicability of the knock-on diagnostic.

Figure 69.35

Calculated spectrum of knock-on deuterons escaping from compressed targets, with constant fuel ρR (40 mg/cm²) and variable $\rho\Delta R$ (1 to 200 mg/cm²) for a CH shell (of temperature 0.5 keV) surrounding the fuel. The energy shift of the high-energy peak measures the $\rho\Delta R$ of the shell, and the number of deuterons in the peak determines the ρR of the fuel.

Figure 69.36

Measurement of the spectrum of charged particles escaping from the target. A spectrometer uses a magnet to momentum select the particles and CCD's or track detectors to detect the particles and provide energy resolution.

Elastically Scattered Neutrons

The second reaction product from (n,D) and (n,T) scattering, namely the elastically scattered neutron, could be used to diagnose all values of ρR of interest to ICF. Like the charged-particle knock-ons, the number of these is directly proportional to the ρR of the fuel but, since they are neutral, they do not slow down in the target, and therefore this diagnostic is not limited to small values of ρR .

To determine if the elastically scattered neutrons can be separated from other sources of neutrons produced in the target, Monte Carlo calculations have been performed to calculate the spectrum of all neutrons emerging from the target. The spectra for two values of ρR are shown in Fig. 69.37. The peak at 14 MeV is, of course, the primary source from DT fusion. The neutrons above 14 MeV are from tertiary reactions, discussed in the next section. In the range of 8 to 10 MeV, there are contributions from many sources including $(n,2n)$ reactions with deuterium and $T(T,2n)\alpha$ reactions; it is not possible to separate out the contribution from elastic scattering in this region. However, in the range of 10 to 13 MeV, elastically scattered neutrons completely dominate the spectrum. The number of neutrons in this range is large enough to provide a model-independent determination of ρR for all experiments planned for OMEGA and NIF. (It should be noted that the two curves shown in Fig. 69.37 are distinguishable because they are both normalized to the primary yield.)

The main effort in neutron diagnostic development is to devise a method for shielding against the effects of the primary DT neutrons. These 14-MeV neutrons can lose energy as they scatter from structural material around the target chamber or

within the diagnostic instrument itself, and they can enter the 10- to 13-MeV window that is being scanned for the elastically scattered neutrons from the target.

Tertiary Neutrons

A significant portion of the tertiary neutron spectrum lies above the 14-MeV primary source (Fig. 69.37). There are no other reactions that can produce neutrons in the range of ~15 to 30 MeV. However, because this is a tertiary reaction, the number of neutrons produced is several orders of magnitude lower than for elastic scattering. For the NIF, yields are sufficiently high that this should not be a problem. However, for OMEGA experiments, it might be difficult to collect a statistically useful number of tertiaries for targets with low neutron yields. Time-of-flight detectors with a solid angle of $\sim 10^{-5}$ should be adequate for the NIF. On OMEGA, it might be necessary to use carbon activation foils that could increase the detection solid angle by an order of magnitude. The carbon foils would detect all neutrons with energy above ~18 MeV, without spectral resolution. However, since there are no other sources of neutrons in this range, spectral separation is not necessary.

For OMEGA implosions, the "cold" part of the fuel will have temperatures below ~1 keV throughout the implosion. In this range, the slowing down of the high-energy knock-on deuterons is relatively temperature independent, for the intermediate step in tertiary-neutron production. Figure 69.38 shows the expected ratio of tertiary neutrons (detected by a carbon foil) to the primary DT fusion yield as a function of ρR , assuming temperatures of 0.5 keV and 1 keV. The neutron ratio determines ρR to within $\pm 5\%$ for these conditions. Superim-

TC3973

Figure 69.37

Calculated spectrum of neutrons escaping from a compressed target for two values of ρR . In the energy range between 10 MeV and 13 MeV, elastically scattered neutrons dominate; above 15 MeV, only tertiary neutrons contribute.

TC4309

Figure 69.38

Ratio of tertiary neutrons to the primary DT fusion yield as a function of fuel ρR for two values of fuel temperature.

posed on Fig. 69.38 are results from a full hydrodynamic simulation of an OMEGA implosion showing how the neutron ratio varies with the neutron-weighted ρR at different times during the implosion. The neutron ratio varies roughly as $(\rho R)^2$, as the range of the knock-on deuterons and tritons is larger than or comparable to the radius of the target.

For a burning NIF target, the temperatures will be considerably higher and the temperature dependence of the diagnostic will be much larger. Uncertainties of interpretation due to this temperature sensitivity can be reduced by a detailed analysis of the tertiary spectrum using a neutron time-of-flight detector.

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Status of Optics on the OMEGA Laser after Eighteen Months of Operation

The OMEGA laser has sustained approximately 1000 target shots over its first 18 months of operation without significant damage to the optics. Both rod and disk amplifiers are used in the OMEGA laser chain; after each amplification stage at 1ω the beams are spatially filtered to remove high-spatial-frequency noise in the beam and to ensure correct image relaying. Optical damage thresholds dictated the minimum beam diameter in each stage and fixed the magnification between stages. The final 28-cm beam propagates through thin-film polarizers before reaching the frequency-conversion cells.¹⁻³ Frequency conversion to the third harmonic (351 nm) is carried out using the polarization-mismatch method developed at LLE.^{4,5} A set of mirrors transports the 3ω beams through a final focus lens, and the last optic in the beam path serves as the vacuum barrier for the target chamber.

Fundamental to LLE's OMEGA operating philosophy is a maintenance plan to ensure (1) prevention of costly damage to system components and (2) reliable and efficient operation of the system.⁶ One part of this preventative maintenance plan involves an estimated 3000 optics on the OMEGA laser system, not including laser glass, driver-line optics, and diagnostic components. The Optical Manufacturing Group is responsible for coordinating the inspection of these optics and tracking this information. On a periodic basis, the optical components are inspected for laser damage, coating degradation, and particulate contamination. During planned shutdown periods, extensive maintenance (cleaning, recoating) and replacement of optics occur. To date, the majority of optics on the OMEGA laser system have not been damaged, and there is no propagating bulk or surface damage to components downstream of damaged optics. While the current optics replacement rate is below the anticipated rate of 10% per year, this rate is expected to increase as the laser ages. In the sections that follow, a review of the types of optical component degradation is presented.

Laser-Damage Classification

In discussing laser damage, a distinction is made between operational and accidental system damage. An optic damaged when the laser operates within normal design limits is classi-

fied under operational damage; whereas, accidental damage occurs when the laser operates outside the design limits. This classification is useful because accidental damage can result in significant damage to an optic, but it is not an indication of continuing problems with that component. Tracking optics damaged due to routine laser operation results in statistical information about the longevity of a particular optic; this information is incorporated into decisions regarding an appropriate spare optic inventory. Observations of an optic's mean lifetime are necessary to determine an effective and economical spares program.

1. Operational Laser Damage

As the laser continues to operate, it is not surprising to observe operational damage in optics located at high-fluence positions of the laser. Figure 70.1 shows a staging diagram that plots the peak design fluence—average fluence times a 1.8 intensity modulation factor—at each stage of a single beamline on OMEGA.⁷ The bold lines indicate areas where laser damage is occurring at present. Specific damage morphologies are presented on p. 52.

2. Accidental Laser Damage

While accidental laser damage occurs infrequently, it can result in significant optics damage. During the activation stage of OMEGA, six optics were damaged in a single shot when light back-propagated after striking the edge of an improperly positioned spatial filter pinhole assembly. The ring of damage is located at the edge of the beam diameter and was caused by the radial edge gain from the rod amplifiers. This beam imprint was observed on several mirrors and the liquid crystal component shown in Fig. 70.2.

Another instance of accidental damage was the development of a small damage site in one-third of the 3ω final focus lenses. A design analysis that locates the focus of unwanted surface reflections eliminated the possibility of a ghost reflection as the cause of this damage site.⁸ The single, bulk damage site is less than 1 mm in diameter and is located at the optic's center near the exit surface. The damage has not propagated to

Figure 70.1

Peak design fluence plotted at each stage of the OMEGA laser. Operational laser damage is occurring at the high fluence positions.

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Figure 70.2

A malfunctioning spatial filter pinhole assembly caused the accidental damage to this liquid crystal polarizer component (135-mm diameter).

the blastshield optic located about 75 mm in front of the lens or to opposing beam port optics. Figure 70.3 shows a timeline of the discovery of damaged focus lenses. At the end of OMEGA activation, a large number of focus lenses were suspected to have damage, with only a slight increase since the initial observation. We believe this damage does not result from the normal operation of the laser because the plot of energy does not correlate with the number of damaged lenses

reported. However, continued laser operation does contribute to the growth of the existing accidentally damaged site. During activation, a train of short pulses propagated down the system and were frequency converted to 2ω , instead of 3ω , which may have caused the damage. Since the sol-gel coating was designed for 3ω , a 2%–3% surface reflection occurs at 2ω . Analysis of multiple beam reflections inside the lens results in the formation of a 2ω caustic at the lens' center position 2 mm from the exit surface.⁹

Damage Morphologies

The damage observed on an optic can be classified as either stable or unstable. Given that the optical components will damage, it becomes important to distinguish between damage that is catastrophic to laser operation and damage that is tolerable. Stable damage does not affect the laser's performance and does not cause propagation of damage to other downstream optics. In addition, this type of damage does not appear to grow significantly on successive laser shots. Unstable damage does grow on subsequent laser shots and eventually prevents the safe operation of the laser. Unstable damage may result in catastrophic failure of a component; additionally, this type of damage could cause propagation of damage to other optics in a beamline. For these reasons, optics with unstable damage are monitored frequently for changes in damage size, and replacement occurs before the damage reaches a critical value.

Figure 70.3

Observation timeline of OMEGA final focus lenses. The left axis plots the number of damaged focus lenses observed over the lifetime of the laser. The right axis plots the 3ω cumulative energy per beam over that same time frame.

1. Stable Damage

The occurrence of stable damage on OMEGA has been documented since activation of the laser system. Two forms of damage have been observed: (1) self-focusing damage in spatial filter lenses and (2) plasma scalds on transport mirrors.

Self-focusing damage is seen infrequently in OMEGA lenses. Currently, nine fused silica lenses have evidence of “angel-hair” tracking, which occurs in the bulk material. These tracks appear as a linear sequence of microscopic damage in the bulk of the optic aligned with the propagation axis. Table 70.I provides a list of the damaged components and the peak fluences they are subject to.

Figure 70.4 shows a photograph of self-focusing damage observed on an OMEGA lens. The large white spot is in the exit surface of the lens, and the typical diameter of this damage area is 1 mm. This large area contains a number of smaller, individual damage sites formed when the self-focusing tracks intercepted the output surface of the lens. These sites do not seem to grow on successive laser shots.

Plasma scalds are frequently seen on OMEGA’s 3ω transport mirrors. As seen in Fig. 70.5, scalds appear as a

Figure 70.4

“Angel-hair” tracking damage observed in the bulk glass of an OMEGA stage-F input lens. The large bright area is the exit surface of the lens.

Table 70.I: Summary of the number of optics with self-focusing damage.

Component	Damaged Lenses/ Total Lenses	Design Peak Fluence (J/cm ²) with 1.1-ns pulse
Stage-E input spatial filter lens (1ω)	1/60	5.0
Stage-F input spatial filter lens (1ω)	5/60	6.6
Primary pickoff lens (3ω)	3/60	2.9

Figure 70.5
Plasma scalds are frequently seen on OMEGA's 3ω transport mirrors.

discoloration on the coated surface of the mirror. Scalds are initiated where particulates are on the mirror surface during a laser shot and are caused by an increase of temperature on the surface during plasma formation.¹⁰ OMEGA transport mirrors are *not* contained in enclosed structures but are mounted in open, upward/downward-facing structures. This open architecture, which allows personnel access to target chamber equipment, results in an increased level of contamination on target bay optics as opposed to laser bay optics, which are vertically mounted with protecting covers on the structures. These plasma-scald defects do not grow in size on subsequent

laser shots, but the number of scalds per surface may increase. At this time, no reflectance data is available to determine how plasma scalds affect performance of the mirror coating.

2. Unstable Damage

OMEGA optics that exhibit unstable damage include (1) tantala/silica-coated beam splitters, (2) BK-7 spatial filter lenses, and (3) fused-silica, input spatial filter lenses. Damage to these optics continues to grow on subsequent laser shots. The rate of growth and size of damage are dependent upon the type of damage observed.

All tantala/silica **beam splitters** located in OMEGA's C-split region (which is the region with the highest fluence of any multilayer coating) have significant damage after 18 months of laser operation. From Fig. 70.6, it is apparent that this damage was not accidental but operational laser damage. As the beam splitters continued to see an increase in 1ω cumulative energy, the observed number of damaged beam splitters increased. Microscopic evaluation of a beam splitter revealed coating-damage sites that were initially 50 to 100 μm in size and 7 μm deep. This depth indicates that the tantala/silica coating is damaged at the substrate interface. The size of these damage sites continues to grow on successive laser shots. Figure 70.7 shows a photograph of a beam splitter with the damage sites illuminated within the beam diameter.

During the OMEGA Upgrade, two types of beam splitters were produced using either the tantala/silica design or the hafnia(metal)/silica design. Hafnia(metal)/silica-coated beam splitters located in the same area as the tantala/silica beam splitters show no signs of coating damage, even though they have the same test damage threshold.¹¹ These test damage thresholds (reported in Fig. 70.8) were taken on a limited number of sites, and, as a result of this limited statistical process,

Figure 70.6
Observation timeline of OMEGA's stage-C beam splitters. The left axis plots the number of damaged beam splitters observed over the lifetime of the laser. The right axis plots the 1ω cumulative energy per beam over that same time frame.

Figure 70.7

Coating damage observed on a 5.5-in. $\times$ 7.5-in. OMEGA stage-C, tantala/silica beam splitter. The highest fluence seen by a multilayer coating is in the C-split area.

defect sites may have been missed. The results on OMEGA are equivalent to full-aperture testing, and this type of testing has revealed significant differences in damage thresholds between the two materials. The highest fluence seen by a multilayer coating is in the C-split area (which is 4.1 J/cm^2). Beam splitters in other areas see a lower fluence (typically $\sim 1 \text{ J/cm}^2$), and neither design has demonstrated any sign of significant damage in these low fluence areas.

Isolated bulk damage sites have been seen in a number of OMEGA **BK-7 spatial filter lenses**. A summary of the location and number of damaged optics is provided in Table 70.II. The peak design fluence is the average fluence times a 1.8 intensity modulation factor.

Each lens has one to four damage sites, ranging in size from 50 to $800 \mu\text{m}$, occurring at varying depths within the bulk glass. Photographs of damage sites from an OMEGA lens are shown in Fig. 70.9. Microscopic evaluation revealed damage sites

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Figure 70.8

Damage thresholds for tantala/silica and hafnia(metal)/silica beam-splitter coatings measured at 1054 nm with a 1-ns pulse at *p*-polarization.

with radial fractures forming a circular pattern perpendicular to the beam. This type of fracture pattern is characteristic of platinum-inclusion damage.¹² Previous data reported by LLNL showed that platinum inclusions in phosphate glass cause fractures when irradiated by laser pulses above a threshold value of 2.5 to 3.0 J/cm^2 with a 1-ns pulse.¹³ These fractured sites do grow upon repeated irradiation, but this growth rate is slow with respect to a bimonthly inspection cycle. During inspection cycles, sizes are visually estimated by observation with a white-light source. Quantitative microscopic evaluation is difficult on *in-situ* optics inspection, so an accurate estimate of damage growth rate is not available. Currently, we have not observed any damage propagation from these lenses to other optics in the beam path.

Table 70.II: Summary of the number of optics with bulk damage sites.

Component	Damaged Lenses/ Total Lenses	Design Peak Fluence (J/cm^2) w/1.1-ns pulse
Stage-C relay lens (in/out)	8/60	2.3
Stage-E output lens	7/60	3.0
Stage-F output lens	34/60	3.6

Large-scale fractured damage sites have been observed in OMEGA's stage-C, -D, -E, and -F input **fused-silica spatial filter lenses**. These areas of the system see some of the highest 1ω fluences and exceed the 2.5 J/cm^2 criterion for the use of fused silica as opposed to BK7 glass (see Fig. 70.1 for the peak fluence in the OMEGA laser chain). The highest peak design 1ω fluence is at the F-input lenses with a fluence of 6.6 J/cm^2 . The lowest 1ω fluence seen by these lenses is the C input, which has a fluence of 4.1 J/cm^2 . Figure 70.10 shows a good correlation between the number of damaged lenses and the 1ω cumulative energy per beam plotted over the same time frame. As the 1ω cumulative energy per beam increases, the number of lenses with fractured damage sites also increases, indicating a form of operational damage.

The damage occurs on the exit vacuum side of the lens at random locations, and the sites grow in size on subsequent laser shots. The damage does not correlate to hot spots in the beam and is believed to be initiated at surface particulate sites. Some possible contamination sources are particulate blowoff from the pinhole and debris from fabrication of the spatial filter vacuum tubes. A fraction of the beam strikes the edge of the pinhole, causing material to vaporize and redeposit on the vacuum side of the input spatial filter lens. Investigation into the cause of this damage is ongoing. The fracture pattern from one of these damage sites is seen in Fig. 70.11.

The vacuum surface of an OMEGA spatial filter lens is under tensile stress, and any damage to this vacuum surface can

Figure 70.9

Damage patterns observed on an OMEGA stage-F, BK-7, output spatial filter lens. Fracture patterns indicate platinum-inclusion damage.

Figure 70.10

Observation timeline of OMEGA fused-silica, input spatial filter lenses. The left axis plots the number of damaged spatial filter lenses observed over the lifetime of the laser. The right axis plots the 1ω cumulative energy per beam over that same time frame.

Figure 70.11
Fractured damage site observed on an OMEGA stage-C input, spatial filter lens.

lead to catastrophic crack growth if the flaw size reaches a critical value a_c . The critical flaw depth a_c depends on the shape of the flaw with respect to the applied stresses and can be calculated from¹⁴

$$a_c = \frac{(K_c)^2}{\pi(Y\sigma)^2},$$

where K_c = fracture toughness of the glass, Y = geometrical factor of the flaw, and σ = bending stress induced by atmospheric pressure.

Actual defects on OMEGA spatial filter lenses are shallow and elliptical, and these defects can be simulated with a half-penny-shaped defect ($Y = 1$) that has a surface diameter twice the defect depth. For an OMEGA stage-F input lens, 25 mm thick and 283 mm in diameter, a critical flaw depth of 8 mm is calculated for a half-penny defect on the vacuum side of a lens subjected to a tensile stress of 615 psi. A defect of this size will be easily detected before catastrophic failure occurs.

It is also important to understand how these lenses will fracture to avoid implosion of the lens when it is under a full vacuum load. Lens fracture on Nova and Beamlet was modeled at LLNL, where it was reported that the design of a "fail-safe" lens probably requires a stress $\sigma < 700$ psi.¹⁵ A fail-safe lens would generate one large fracture and would not implode. The

maximum bending stress for an OMEGA stage-F input lens is 615 psi and occurs at the center of the lens. If the model is correct, these lenses under full vacuum should not implode into several fragments when defects reach their critical flaw size but should crack into two pieces and lock together as long as the mount restrains the radial motion of the fragments. Although damage on the vacuum side of OMEGA spatial filter lenses has been observed, there has been no incident where an OMEGA spatial filter lens has cracked into several pieces. For safety concerns, OMEGA optics will be removed when defects reach one-half their critical flaw size (i.e., when the depth is 4 mm). For convenience we measure the surface trace *in situ* and replace any lenses with a surface trace of 8 mm (which would imply a depth of 4 mm if the flaw is penny shaped). We have replaced two lenses that had surface traces of 8 mm. Later depth measurements showed only a depth of 2 mm, well within the criteria for failure.

Optical Property Changes

OMEGA's sol-gel coatings show a significant change in reflectivity when exposed to a vacuum. A guided wave spectrometer is used to measure the reflectivity of optics mounted on OMEGA. While the instrument is not photometrically accurate, the shape of the spectral curve does provide information on the performance of the coating. The spectral data in Fig. 70.12 was taken on a 1ω (1053-nm) sol-gel-coated spatial filter lens. While the S1, non-vacuum-side, reflectivity curve is as expected for a 1ω antireflection coating, the S2, vacuum-side data indicates that the porous sol-gel coating has adsorbed some organic material, which has affected its coating performance. Evaluation of the contaminated sol-gel coating by gas chromatography mass spectrometry detected the presence of

Figure 70.12
Reflectance data measured on a sol-gel-coated spatial filter lens. The S2 (in) surface is the vacuum interface. The S1 (out) surface is not exposed to a vacuum.

vacuum pump oil. This coating problem is seen on all OMEGA sol-gel-coated spatial filter lenses that are pumped by a mechanically pumped system. Hard-oxide dielectric coatings pumped under similar vacuum conditions and sol-gel coatings pumped by an ion sublimation pump show no change in reflectivity after exposure to a vacuum. Since OMEGA's optical layout was designed with the assumption of a 4% light loss per surface, damage is not expected from ghost reflections and no system damage has occurred as a result of this coating problem.¹⁶ Two solutions to this sol-gel-coating problem are being investigated: (1) replace the sol-gel coatings with hard-oxide dielectric coatings and (2) improve the spatial filter pumping system. Sol-gel coatings were used on all fused silica spatial filter lenses since these lenses have the highest 1 ω fluence on the system and sol-gel coatings have, in the past, shown considerably higher damage threshold than hard-oxide coatings. Recent improvements in the hard-oxide coating process have resulted in increased damage thresholds, which now make them viable replacements to the sol-gel type.

Target Chamber Optics

After 1000 target shots, no target chamber optics have been replaced because of damage. The inside target surface of the blastshield optics has degraded from the impact of target debris, and the damage is in the form of shallow pits ranging in size from 50 μ m to 1 mm. Recently, self-focusing damage has been observed in several blastshields' optics, which serve as the vacuum barrier for the target chamber. Damage to the outside blastshield surface (non-vacuum side) is minimal and usually occurs as a result of mishandling during blastshield replacement. Degradation in the blastshields' sol-gel coating occurs with prolonged exposure (over 100 days) to the target chamber. This degradation leads to reduced transmission of laser energy to target center. The on-target imaging system (OTIS) diagnostic is used to measure the irradiation uniformity and can therefore be used as an *in-situ* measurement of the performance of the sol-gel coating. We use the OTIS measurement to recommend replacement of the optics with the lowest UV transmission. The baseline replacement rate of blastshield optics is 20 per month; optics are replaced during a regular maintenance day.

Conclusion

LLE has implemented a plan to maintain the quality of OMEGA optics. This preventative maintenance plan involves frequent inspections and *in-situ* cleaning of optics by a skilled support group as well as proper training of laser operations personnel to prevent optics handling damage. Damaged optics are placed on a critical optics list and are tracked frequently for damage growth. This critical optics list is used to report the current condition of OMEGA optics to the Laser Facility Manager.

This article has focused on specific optics that have started to damage after 18 months of OMEGA laser operation. Components that have *not* damaged include frequency-conversion crystals, polarizers, calorimeters, and liquid-crystal optics. In addition, there has been no propagating bulk or surface damage to components downstream of damaged optics.

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A Single-Beam, Ponderomotive-Optical Trap for Energetic Free Electrons

Traditionally, there have been many advantages to using laser beams with Gaussian spatial profiles in the study of high-field atomic physics. High peak intensities are achieved due to their focusability, they are easily generated in most laser systems, and their focal region can be described analytically. However, free electrons interacting with such a field are rapidly accelerated out of the high-intensity region via the ponderomotive force (which is proportional to the gradient of the laser intensity).^{1,2} As a result, the electrons' interaction time with the intense portions of the laser focus can be much shorter than the pulse duration. In order to observe harmonic generation from oscillating free electrons, one must first control the expulsion of electrons from the focal region.³ This confinement can be most easily achieved by creating an intensity minimum at the focus, thereby using the ponderomotive force to push the electrons toward the central minimum. If this intensity minimum is non-zero, then the electrons can interact with intense fields while remaining trapped. Such ponderomotive trapping has been proposed in the past,^{4,5} and specific laser-based schemes to trap electrons in the radial direction have been described.^{6,7} To our knowledge, such focal spots have not been generated with a high-power laser. In addition, these proposed traps would not confine electrons in the axial (laser propagation) direction. In this article we will present a novel scheme to trap electrons in three dimensions with a single laser beam.

Numerical Results

A "trapping" focal region consists of an intensity minimum surrounded on all sides by higher intensities. To understand how to create such a region, it is important to adopt a formalism to describe the propagation of light. Huygens' principle, which states that we can consider every point on a wave front to be a point source for a spherical wave, makes it possible to find the field distribution in some "observation" plane if the field distribution in an earlier "source" plane is known.⁸ A computer program has been written to calculate the three-dimensional focal region from an arbitrary, monochromatic laser beam. The "source" plane is the field distribution at the lens, and the "observation" plane is scanned through the focal region. The incident intensity distribution can be created mathematically

(Gaussian, super-Gaussian, flat-top, etc.) or can be entered as a digitized image. The incident phase front is assumed to be uniform, and the lens is assumed to be perfect (although the addition of aberrations is possible). The beam can be passed through phase and amplitude masks on-line in an effort to alter the shape of the focal region.

By passing the beam through amplitude masks, losses are introduced into the system, yet the shape of the focal region is changed only slightly. Figure 70.13(a) shows an incident Gaussian beam [where w is the $1/e^2$ (intensity) radius]. Figure 70.13(b) shows the intensity distribution of the focal region as a function of z (axial or laser propagation direction) and r (radial direction). The Rayleigh range (z_0) and the $1/e^2$ waist (w_0) are shown in the figure. Figure 70.13(c) shows the focal spot at $z = 0$. Blocking the center of the beam [Figs. 70.13(d)–70.13(f)] or blocking an annular section [Figs. 70.13(g)–70.13(i)] has little effect on the intensity distribution near the focus. A centrally peaked distribution is formed in all cases.⁹

As Casperson has shown,¹⁰ modulating the phase of the incident beam can result in deep far-field amplitude modulation with little loss in total energy. This can be accomplished by placing phase masks in the path of the beam. Of particular interest is the effect of binary phase plates, which add either zero or π phase to portions of the beam. If half of the Huygens' spherical waves encounter a π -phase shift while the other half do not, we can expect complete destructive interference where they meet. For a focused laser beam, this occurs at the center of the focal region. A simple, two-zoned binary phase plate is shown in Fig. 70.14(a). By passing a Gaussian beam through this plate and focusing it, an altered intensity distribution [Fig. 70.14(b)] is formed. It is important to note that this distribution is not symmetric about the z axis and, therefore, would confine particles in only one dimension. The focal spot is shown in Fig. 70.14(c). A beam similar to the TEM_{01}^* mode (commonly referred to as a "donut" mode) can be generated by passing a Gaussian TEM_{00} through a smoothly varying helical phase plate [Fig. 70.14(d)]. Such a phase plate has been generated for millimeter¹¹ as well as optical wavelengths.^{12,13}

The resulting focal region [Figs. 70.14(e) and 70.14(f)] confines electrons in the radial direction only. It is also difficult to "tune" the trap minimum away from zero. A scheme has been proposed that combines a donut beam focus with a centrally peaked focus.⁶ The added complexity of mixing two beams results in a focal region with a tunable center intensity, but the trap would not confine particles in the axial direction.

The binary phase plate shown in Fig. 70.14(g) contains only two zones, yet the focal spot [Fig. 70.14(i)] is surprisingly similar to the one generated by the complicated helical plate. The donut shape near $z = 0$ confines electrons in the radial direction, and the centrally peaked regions away from $z = 0$ provide axial trapping [see Fig. 70.14(h)]. The focal region is shown in detail in Fig. 70.15. The π region of the phase plate

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Figure 70.13

The effects of amplitude masking on the intensity distributions of a focused Gaussian beam. The first column shows three different near-field intensity distributions. The second column shows the intensity distributions as a function of z and r (the center of each image corresponds to $z = 0, r = 0$). The third column shows the focal spots. In order to show the maximum contrast, each image is normalized to its peak value. Note that all of the focal regions are centrally peaked.

has a diameter of $1.65 w$ [where w is the $1/e^2$ (intensity) radius of the incident Gaussian beam]. This results in a π -phase shift for half of the incident field. The calculated intensity in the focus is normalized to the peak intensity in the absence of the phase plate. The r and z positions are normalized to the unaltered beam waist w_0 ($1/e^2$ radius) and the unaltered beam Rayleigh range z_0 , respectively. Figure 70.15(a) shows a surface plot and Fig. 70.15(b) shows a contour plot of the trapping region. There is an exact zero in intensity at the center of the

focus, with intensity walls ranging from $\sim 8\%$ to $\sim 30\%$ (of the unaltered peak intensity) in all directions. The trapping region has a volume of complete trapping of $\sim 2w_0^2 z_0$ (bound by the solid contour line of 8.2%). By changing the size of the π region of the phase plate (or, equivalently, by changing the size of the incident beam), the destructive interference at the center of the focus will not be complete, creating a non-zero minimum. Figure 70.16(a) shows the calculated focal distribution for a Gaussian beam incident on a phase plate with a

ES410

Figure 70.14

The effects of phase masking on the intensity distributions of a focused Gaussian beam. The first column shows three different phase masks. The second column shows the intensity distributions as a function of z and r (the center of each image corresponds to $z = 0, r = 0$) when the phase mask is placed before the incident beam. The third column shows the focal spots. In order to show the maximum contrast, each image is normalized to its peak value.

π -region diameter of $2.20 w$. Figure 70.16(b) shows a contour plot of the trapping region. In this case, the bottom of the trap is $\sim 17\%$ of the unaltered beam peak intensity, and the trap walls range from $\sim 24\%$ to $\sim 50\%$. The trapping volume is $\sim w_0^2 z_0$. This focal region is ideal for trapping electrons in a high-field region.

Experimental Setup

The phase masks described above can be made by using photolithographic techniques. However, for the particularly simple geometry that produces a suitable trapping region, a novel segmented wave-plate approach has been developed. A half-wave plate is typically used to rotate the polarization angle of an incident, linearly polarized beam of light. This is a result of the retardation of π phase between two orthogonal incident polarizations. However, if part of the incident beam can be

made to travel through the plate as a fast wave while the other travels as a slow wave, a π -phase shift will be created between the two portions. To accomplish this, a 4-cm disk was cut from the center of an 8-cm mica half-wave plate. The disk and annulus were then mounted individually on cross hairs and aligned to the beam path (as shown in Fig. 70.17). In Fig. 70.17, the laser is polarized in the vertical direction, and the annulus is arranged such that its e axis is vertical, while the disk has its o axis vertical. This results in a π -phase shift of a portion of the incident field. The size of the disk was chosen such that approximately half of the incident field passed through it. This resulted in near-zero intensity at the center of the focus. By rotating the two pieces as a unit, one can rotate the polarization of the incident beam without introducing additional optics or changing the shape of the focal region. Also, as long as one does not require a single polarization direction in the focal

E8339

Figure 70.15

(a) A computer-generated trapping focal region with zero intensity at its center (π -region diameter = $1.65 w$). (b) A contour plot of the trapping region. The volume of complete trapping is bound by the solid contour line at 8.2%.

E8340

Figure 70.16

(a) A computer-generated trapping focal region with 17% intensity at its center (π -region diameter = $2.20 w$). (b) A contour plot of the trapping region. The volume of complete trapping is bound by the solid contour line at 23.6%.

Figure 70.17

The experimental setup for the segmented wave-plate technique for generating a trapping focal region.

region, the depth of the trap can be tuned by rotating the central disk away from 90° . This results in a polarization rotation of the inner portion of the beam, which will put some of the field into the noninterfering orthogonal polarization, thereby filling in the minimum. The extreme case is to rotate the disk a full 90° , back to its original position, which produced an ordinary, centrally peaked focus.

Experimental Results

To image the altered focal spot, the wave-plate arrangement was aligned to the beam path of the tabletop terawatt laser (T^3) directly before the focusing lens ($f = 212$ cm).¹⁴ The focal region was imaged with a CCD camera coupled to a $10\times$ microscope objective (Fig. 70.18). The camera-objective combination was moved together to map out the focal region. A single laser shot was taken at each position, and each image was minimally smoothed and background subtracted. The scan was generated with ~ 40 -mJ, ~ 2 -ps infrared laser pulses. By firing the final, single-pass amplifier, energies of ~ 1 J are achieved with no noticeable change in the near- or far-field characteristics of the beam. The unaltered beam's peak intensity typically reaches 10^{18} W/cm².

Figure 70.19(a) shows the focal spot created without the wave-plate arrangement in place. The beam diameter is approximately $1.5\times$ the diffraction limit. Figure 70.19(b) shows the focal spot with the wave-plate arrangement set to 90° . Despite the deviation of the unaltered beam from the diffraction limit, there is a well-defined minimum in the center of the altered spot. Figure 70.19(c) shows the focal spot with the wave-plate arrangement in place, but set to 0° . Here, the wave plate is set to its original, unaltered configuration, and the beam has returned to its original shape. Figures 70.20(a) and 70.20(b) show the lineouts of the altered beam in the x and y directions at $z = 0$. The x and y positions are normalized to the experimentally determined beam waist of the unaltered beam. Figure 70.20(c) shows the intensity of the beam at $r = 0$ as a function of z . The intensities in Fig. 70.20 are normalized to the peak intensity of the unaltered beam.

A more detailed scan was made with a continuous-wave beam (the oscillator of the T^3 laser). Figure 70.21 shows two slices through the focal volume along the z axis. Figure 70.21(a) shows the intensity distribution in the plane $y = 0$, and Fig. 70.21(b) shows the distribution for $x = 0$. In each

Figure 70.18

The experimental setup for imaging the altered focal region. Here $z = 0$; i.e., the object plane of the microscope objective coincides with the focal plane of the lens.

Figure 70.19

Focal spots (a) without the wave-plate arrangement in place, (b) wave-plate setup in place and set to 90° , and (c) wave-plate setup in place and set to 0° .

Figure 70.20

Lineouts of the trapping focal region. (a) Lineout in the x direction, (b) lineout in the y direction, and (c) the intensity at $r = 0$ as a function of z . All axes are normalized to the experimentally determined values for the unaltered beam (spot size for the x and y positions, Rayleigh range for the z positions, and unaltered beam peak intensity for the intensity values).

Figure 70.21

Contour plots of the trapping focal region. (a) In the plane $y = 0$ and (b) in the plane $x = 0$. The experimentally observed volume of complete trapping is bound by the solid contour line at 10%.

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plot, there is a region of complete trapping bound by intensity walls of 10% (solid line) of the unaltered beam's peak intensity. Although the incident beam size was not matched with the size of the wave-plate pieces, a sufficiently large portion of the incident field was shifted to create a trapping region. These experimentally obtained distributions (using a pulsed or cw beam) exhibit a local minimum in intensity at the *origin along all three dimensions*, making them suitable focal regions for ponderomotive trapping.

Summary

In summary, we have produced, for the first time to our knowledge, a tunable, single-beam, three-dimensional, ponderomotive-optical trap for free electrons with a high-peak-power laser system. Neutral atoms and macroscopic, low-index particles could also be trapped in the low-field region. We have presented a novel, segmented-wave-plate approach as an alternative to a phase mask, along with experimental confirmation of the intensity distribution of the altered focal region. Experiments to image the linearly and nonlinearly scattered light (both 1ω and 2ω) from the trapped electrons are being carried out.

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Appendix A: Huygens' Principle

Huygens' principle states that we can consider every point on a wavefront to be a point source for a spherical wave.⁸ As a result, if one knows the field in some "source plane (S')," it is then possible to add the contributions from the individual spherical waves to find the field distribution in an "observation plane (S)" (see Fig. 70.22). The spherical wave emanating from a tiny element of area $\Delta a'$, surrounding the point $(x', y', 0)$ in the source plane, results in a contribution to the field in the observation plane,

$$\Delta \varepsilon(x, y, z) = -(i/\lambda) \left(e^{ikr} / r \right) \varepsilon(x', y', 0) \Delta a', \quad (A1)$$

where

$$r = \left[(x - x')^2 + (y - y')^2 + z^2 \right]^{1/2}, \quad (A2)$$

$k = 2\pi/\lambda$, $\Delta a' = dx' dy'$, and $\varepsilon(x', y', 0)$ is the electric field (magnitude and phase) in the source plane. For most applications, we may make the Fresnel approximation

$$z^3 \gg (\pi/4\lambda) \left[(x - x')^2 + (y - y')^2 \right]_{\max}^2 \quad (A3)$$

and sum up all of the field contributions with the integral

$$\varepsilon(x, y, z) \approx - \left(i e^{ikz} / \lambda z \right) \iint \varepsilon(x', y', 0) \exp \left\{ ik \left[(x - x')^2 + (y - y')^2 \right] / 2z \right\} dx' dy'. \quad (A4)$$

To calculate the field at the point (x, y, z) , $\varepsilon(x', y', 0)$ must be known for all (x', y') . Equation (A4) can then be solved numerically.

For the case of a laser beam incident on a positive lens, the intensity distribution can be assumed to be Gaussian (as is typical for many laser systems) or it can be determined experimentally by imaging the near-field distribution. Determining the phase of the incident field is more complicated, as it involves interferometric methods.¹⁵ If the laser is well collimated and is found to focus to a spot size close to the diffraction limit, then the phase front can be assumed to be uniform. The effect of the lens is to simply add a phase proportional to the thickness of the lens material at any point. A plano-convex lens results in a "spherical" phase distribution,

$$\Phi(r') = k \left(f^2 - r'^2 \right)^{1/2} \quad (A5)$$

where f is the focal length of the lens and $r'^2 = x'^2 + y'^2$ is the square of the radial location in the source plane. If we assume that $f^2 \gg r'^2$ (the paraxial approximation), then we can perform a binomial expansion,

$$\Phi(r') = kf \left(1 - r'^2 / 2f^2 - r'^4 / 8f^4 - \dots \right). \quad (A6)$$

The first term adds a constant phase to the entire beam and can be discarded, as can the higher-order terms since they are vanishingly small. By keeping only the quadratic term, we are left with a "perfect" lens in the context of this approach (in practice, aspherical lenses are not purely parabolic). A Gaussian laser beam that has passed through this lens will have the field distribution

E8510

Figure 70.22

The geometry of the Huygens' wave approach to the propagation of light. The source and observation planes are separated by a distance z . The locations in the source plane are denoted with a prime ('), while the locations in the observation plane are left unprimed.

$$\epsilon(x', y', 0) = A \exp\left(-r'^2/w^2 - ikr'^2/2f\right), \quad (A7)$$

where A is the field amplitude and w is the beam radius at the $1/e$ point in the electric field amplitude (which is equivalent to the $1/e^2$ point in intensity). It should be pointed out that all calculations in both this article and the computer codes use the scalar-wave approach.

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A Diode-Pumped Nd:YLF Master Oscillator for the OMEGA Laser

The OMEGA laser is a 60-beam laser-fusion system capable of producing a total of 30 kJ of ultraviolet (351-nm) energy on target,¹ where the temporal profile of the optical pulse applied to a laser-fusion target can be specified in advance. To accomplish this, the laser system consists of a master oscillator, a pulse-shaping system,² regenerative amplifiers (regens), large-aperture regenerative amplifiers (LARA), rod and disk amplifier chains, and frequency converters. A schematic of the laser chain is shown in Fig. 70.23. The pulse-shaping system (Fig. 70.24) is based on LLNL's conceptual design³ and includes integrated-optics (IO) modulators driven by electrical waveform generators. The pulse-shaping system produces an optical pulse with a specific temporal pulse shape at the nanojoule energy level, which seeds the regens. The temporal profile of this low-energy seed pulse, when amplified and frequency tripled by the laser system, will compensate (1) for the temporal distortions caused by gain saturation in the regens and amplifiers and (2) for distortions caused by the tripling process and will produce the desired pulse shape on target. The overall gain of the OMEGA laser system from the pulse-shaping modulator to the target is approximately 10^{14} . This large gain puts extremely severe requirements on stability, reproducibility, and reliability of the pulse-shaping system and master oscillator performance. This article details the OMEGA requirements, describes the new diode-pumped mas-

ter oscillator, and presents results from this newly fielded master oscillator.

OMEGA's Master Oscillator Design Requirements

The master oscillator for the OMEGA laser system must satisfy a number of requirements:

1. The wavelength of the oscillator must match that of OMEGA's Nd:phosphate glass amplifiers (1053 nm). This requirement automatically defines the choice of the laser medium for the master oscillator to Nd:YLF.
2. The oscillator must operate in a single-frequency regime. This requirement enables the stable and reproducible generation of complex optical pulse shapes and assures optimal frequency conversion.
3. The master oscillator must be *Q*-switched to produce enough energy to split among four or more pulse-shaping channels and to properly seed the various regens.
4. The duration of the master oscillator's *Q*-switched pulse must be >100 ns, and the timing jitter must be within ~10 ns to provide a constant pulse amplitude within the required 10-ns pulse-shaping window.

Figure 70.23
OMEGA laser system for ICF experiments.

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Figure 70.24

Block diagram of the pulse-shaping system for the OMEGA laser.

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5. The master oscillator's Q -switched pulse must be externally synchronizable to the OMEGA timing system.

At present a flashlamp-pumped Nd:YLF laser is used on OMEGA. Although this laser satisfies all of the above conditions,⁴ maintenance and reliability considerations have led us to develop a new diode-pumped version with significantly improved performance. The long-term maximum amplitude has been increased, and frequency stability has been dramatically improved, while the complexity, maintenance requirements, and footprint of this new laser have shrunk significantly.

Currently there are two major areas of interest in diode-pumped Nd:YLF laser development: high-power cw lasers and short-pulse Q -switched lasers. Even the most powerful cw lasers⁵ cannot provide enough energy within a 10-ns pulse-shaping window to meet our requirement of >100 nJ in a 10-ns pulse in each of four channels. Q -switched, diode-pumped Nd:YLF lasers⁶⁻⁹ are usually designed to generate short (3- to 10-ns) pulses with high peak powers. The short pulse width makes this type of laser unacceptable for satisfying the requirement for a constant pulse amplitude within a 10-ns pulse-shaping window. Passively Q -switched, diode-pumped lasers (for example see Refs. 10 and 11) cannot be precisely synchronized.

To ensure that the laser generates a single frequency, the laser should operate unidirectionally. The most commonly used technique to achieve this involves using an optical diode (Faraday rotator). This technique does require additional intracavity elements and consequently increases resonator losses, thus rendering it unsuitable for our use. Kane and Byer¹² have used nonplanar monolithic ring lasers; however, their use is restricted to laser materials with a large Verdet constant and no birefringence (Nd:YAG is such a material, but Nd:YLF is not). The most suitable way to provide unidirectional operation and Q -switching in a Nd:YLF laser is to employ an acousto-optic (A-O) Q -switch, which can be used for both birefringent¹³ and nonbirefringent¹⁴ materials.

The A-O Q -switched Nd:YLF laser meets most of the important requirements for the OMEGA master oscillator, namely the 1053-nm wavelength, single-frequency operation, high energy, and externally synchronizable requirements. However, close attention must be paid to the long (>100-ns) pulse width generation and high stability requirements.

There are no commercially available lasers with the specifications required for use on the OMEGA laser. In particular the wavelength (1053 nm) and long pulse duration (>100-ns) requirements are not commonly required in commercial systems.

Laser Design

The diode-pumped, single-frequency Nd:YLF laser (Fig. 70.25) is based on a triangular ring cavity with two spherical mirrors and a prism that doubles as an active element. The resonator is very similar to the Nd:YAG laser described in Ref. 15. In that design one of the resonator mirrors was deposited on the laser element; however, in our design we have moved the active element away from one of the mirrors to avoid spatial hole burning.¹⁶ This residual hole-burning effect, when an active element is located at one of the ring cavity mirrors, cannot be fixed in our case by displacing the pumped volume from the cavity TEM₀₀ mode as suggested in Ref. 17 because we use a fiber-coupled pumping diode and the pumped volume is larger than the TEM₀₀-mode volume. We use a fiber-coupled pumping diode so that the water-cooled laser diode can be removed from the laser head for enhanced laser stability. The diameter of the pumped volume for fiber-coupled diodes is relatively large (~0.4 mm) compared to other designs. This large volume reduces the efficiency of the laser but increases the overall stability of the diode-pumped laser. For our application the enhanced stability is more important than the efficiency of the master oscillator. The ring cavity has been chosen to provide traveling-wave unidirectional lasing and hence easier single-frequency operation. Both cavity mirrors are spherical with a 10-cm radius of curvature. The distance between the mirrors is approximately 30 mm. The angle of incidence for both mirrors is ~11°. The end mirror has maxi-

imum transmission for the pump radiation (797 nm) and maximum reflection for the 1053-nm laser wavelength. The output coupler has a 95% reflection coefficient at 1053 nm.

A 4 × 4-mm-cross-section, *a*-cut Nd:YLF element has Brewster-angled entrance and exit surfaces cut for 1053-nm operation (*c* axis is perpendicular to the resonator plane). Its on-axis thickness is 8 mm and has been calculated to compensate for the astigmatism from the two off-axis spherical cavity mirrors.¹⁸ To provide maintenance-free operation no adjustable mounts were used for cavity alignment. The laser was assembled on a 50 × 30 × 30-mm Zerodur base plate to minimize thermal drift and improve long-term stability. A precisely machined metal template allows accurate placement of the carefully machined mirrors (flattened barrels) to eliminate the necessity for fine alignment of these mirrors. The active element is glued to a miniature piezo-translator (Physik Instrumente) for cavity-length adjustment, which is in turn mounted on a Zerodur standoff glued to the Zerodur base plate. The A-O modulator (QS080-2G-RU1, Gooch & Housego) assures a unidirectional prelude phase prior to *Q*-switching. The A-O *Q*-switch has plane-parallel surfaces (AR coated for 1053 nm) and can be aligned without disturbing the cavity alignment. Unidirectional operation is achieved by using the self-feedback mechanism in the A-O *Q*-switch, which is described in detail in Ref. 19.

Figure 70.25
Schematic of the diode-pumped, single-frequency, pulsed OMEGA master oscillator with amplitude and frequency feedback systems.

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The laser is pumped by 700- μ s square pulses at a 5-Hz repetition rate generated by an SDL-3490-P5 fiber-coupled laser diode (nominal 5-W output). The OMEGA front-end laser system operates at 5-Hz repetition rate that reduces thermal problems in the master oscillator and extends the lifetime of the diode. The image of the output 0.4-mm-diam fiber surface is relayed into the active element of the laser by using two 25-mm-focal-length spherical lenses AR coated for 797 nm. Both lenses are aligned and assembled together with the fiber port. The position of the pumped volume can be aligned to the TEM₀₀ cavity mode by adjusting the whole lens-fiber unit.

The low duty cycle of the oscillator obviates the need for a cooling system on the active element. The laser resonator is thermally uncoupled from the environment. The laser requires only two adjustments: alignment of the pumped volume in the

active element and the A-O *Q*-switch position for unidirectional operation. The laser layout is shown in Fig. 70.26.

Experimental Results

We have found that a low rf power of ~ 5 mW_{rf} is sufficient for unidirectional operation. We have achieved a very high contrast for the counter-propagating beams ($>1500:1$). As mentioned above, unidirectional operation is essential to maintain the single-frequency regime. In this operating mode the prelude stage exhibits regular, slowly decaying relaxation oscillations [Fig. 70.27(a)], which cause significant amplitude and temporal jitter of the subsequent *Q*-switched pulse. To stabilize the prelude phase we have developed a negative amplitude feedback with a constant offset. The offset provides the required rf power level to achieve unidirectional operation during the prelude phase. One of the beams diffracted by the A-O modulator is coupled into the 0.4-mm multimode fiber

Figure 70.26
The OMEGA master oscillator layout. It consists of three units: laser cavity (no adjustments required); A-O modulator (adjusted for unidirectional operation); and pump unit (adjusted for maximum laser output).

Figure 70.27
Oscillograms of the prelude phase: (a) single-frequency operation, no amplitude feedback; (b) same with the feedback.

and sent to a diode that generates the feedback signal. Increased feedback increases the rf power to the A-O modulator, thus increasing the cavity losses and stabilizing the laser output. With amplitude feedback stabilization a very smooth prelude phase is observed [Fig. 70.27(b)], and the externally triggerable *Q*-switch leads to high amplitude stability and low temporal jitter of the output pulse.

At the lowest rf power required for unidirectional operation ($\sim 5 \text{ mW}_{\text{rf}}$) the laser generates a $\sim 300\text{-ns}$ pulse of $\sim 5\text{-}\mu\text{J}$ energy. The pulse width is defined by the removable *Q*-switch loss. To maintain unidirectional operation and obtain high stability and low jitter, the laser must work close to threshold. For a given resonator configuration the energy content of the *Q*-switched pulse is closely related to the pulse width. The self-feedback mechanism of the A-O enforced unidirectional operation¹⁹ is crucial for generating long single-frequency pulses because of the very low rf power required for the unidirectional prelude phase and hence low removable loss level. By increasing the offset in the amplitude feedback system and the pumping energy, we have been able to obtain stable single-frequency pulses with widths as short as 30 ns and energies of $\sim 18 \mu\text{J}$. For use in the OMEGA laser chain the laser has been adjusted to generate an externally triggerable single-mode pulse of $\sim 160\text{-ns}$ duration and $\sim 10\text{-}\mu\text{J}$ pulse energy. As a test of the laser stability and reproducibility, 100 laser pulses taken within 0.5 h with the laser operating at 5 Hz were recorded using a

1-GHz, 4-GS/s Hewlett-Packard 54720A digitizing oscilloscope. Figure 70.28(a) shows the average pulse shape (solid line) and its amplitude and temporal rms deviations (dashed lines). Figure 70.28(b) shows the same experimental data as Fig. 70.28(a), but the effect of the build-up jitter has been removed. By comparing Figs. 70.28(a) and 70.28(b) we conclude that the fluctuations seen in Fig. 70.28(a) are mostly due to the build-up time jitter (the long-term temporal jitter of the peak of the pulse is $< 7 \text{ ns rms}$). These data clearly show that the laser demonstrates excellent amplitude stability and pulse-shape reproducibility. The short-term (0.5-h) amplitude stability of the pulse has been measured to be $\sim 0.3\%$ rms. The long-term (8-h) stability is $\sim 0.6\%$ rms. The variation of the pulse width (which has an average FWHM of 159 ns) is $\sim 0.7 \text{ ns rms}$. The achieved pulse duration and the remarkable stability of the pulse provide a constant pulse amplitude, with a variation of $< 1\%$, within the 10-ns pulse-shaping window. To maintain these excellent properties, the output beam must be handled very carefully. Even immeasurably small amounts of back reflections (from such sources as diagnostics, fibers, or fiber couplers) can reduce the stability significantly.

In all the experiments reported here, the laser was operated in the TEM_{00} mode despite the large pump volume. This mode of operation, which involves operating close to threshold with an extended prelude phase, helps the laser maintain a single longitudinal and spatial cavity mode.

Figure 70.28

(a) The average pulse shape (solid line) and amplitude and temporal rms deviations (dashed lines) from the output *Q*-switched pulse for 100 shots taken within 0.5 h . (b) The data as in (a), but with the build-up jitter removed (all three lines, average and average $\pm$ rms fluctuations, virtually overlap). This indicates that all fluctuations in the left graph are mostly due to build-up time jitter of 7 ns rms , which is very small compared to the 159-ns pulse width (FWHM).

The optical efficiency in this Q-switched laser is ~5%. This relatively low optical efficiency is due to several factors: (1) the pump volume is large, exceeding the TEM₀₀ cavity-mode volume; (2) the output coupler has not been optimized; and (3) the pump beam is slightly vignetted by the cavity mirror. However, for this application, optimal efficiency is not of major importance.

Long-term, single-frequency operation and stability are achieved by incorporating a frequency stabilization scheme. In this design 10% of the laser output is expanded and sent through a low-finesse, solid, temperature-controlled etalon to create fringes. Two fiber sensors straddle one of the fringes such that the difference between optical signals in the fibers is zero. Any fringe shift leads to a signed error signal that feeds into the (biased) PZT driver; this corrects the cavity length and drives the error signal to zero.

Conclusions

We have developed a diode-pumped, single-frequency, pulsed Nd:YLF master oscillator for the OMEGA laser system. Output pulses have been generated that have 30- to 300-ns duration and 5- to 18- μ J energy with high amplitude stability (<0.6% rms), low timing jitter (<7 ns rms), and long-term frequency stability of <100 MHz over periods of many hours without observable dual-mode operation. The laser adjusted for 160-ns pulse width and 10- μ J pulse energy satisfies all OMEGA requirements and is currently being incorporated into the OMEGA laser system.

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Heat Transfer from Localized Absorbing Defects to the Host Coating Material in HfO₂/SiO₂ 351-nm High Reflectors

Recent atomic-force microscopy (AFM) studies of laser-induced damage morphology in HfO₂/SiO₂ UV multilayers,¹ caused by 351-nm, 1-ns laser pulses, revealed submicron-sized craters as the dominant damage morphology in these films [see Fig. 70.29(a)]. The origin of such craters was attributed to nanoscale, localized absorbers in the film matrix, the smallest size of which was estimated to be <10 nm. However, at this time, both the exact physical nature of these absorbers and the mechanism of energy transfer to the host coating material in the promotion of damage remain speculative. Empirical evidence¹ suggests that, at some time in the damage evolution, melting occurs; furthermore, near damage threshold, even the smallest absorbing defects are able to induce such melting to extend to about 50 nm inside the coating. For incident fluences 5%–20% above this threshold, the crater-depth distribution shifts dramatically, as measured by AFM (see Fig. 70.29), while for the lower fluences the melt processes and crater formation are largely confined to the top HfO₂ layer [Fig. 70.29(b)]—these processes penetrate into the next-from-the-top, SiO₂ layer at fluences ~20% above threshold [Fig. 70.29(d)].

In this work, by carrying out heat-transfer calculations for the same HfO₂/SiO₂ multilayer system, our goal was to answer the following questions:

- How high a temperature must an absorbing defect reach to support film-crater formation of the type observed experimentally by AFM?
- How does heat *conduction* affect the damage kinetics at these temperatures and for different defect and host characteristics?
- May heat conduction be considered the dominant energy-transfer mechanism, or must other mechanisms be included for the process physics to remain plausible?

In the following, the model assumptions and simulation results will be presented. From these results, one concludes that, for realistically sized defects and absorption conditions

for 351-nm laser light, a damage model commensurate with empirical, AFM evidence *must* account for energy-transfer mechanisms other than heat conduction alone.

Methodology

Heat-transfer modeling was accomplished by using the finite-element analysis code ANSYS² in rigid-grid form. The equation of motion was treated in matrix form

$$[c] \{\dot{T}\} + [K] \{T\} = \{Q\}, \quad (1)$$

where $[c]$ represents the specific-heat matrix, $[K]$ the conductivity matrix, $\{T\}$ the nodal temperature vector, and $\{Q\}$ the nodal heat-flow vector. The model assumed *convective* heat transfer at the film/air interface and *conduction* inside both the defect and the host material. The model geometry, depicted in Fig. 70.30, contains alternating quarter-wave layers ($\lambda = 351$ nm) of HfO₂ and SiO₂ of cylindrical geometry with the r axis along the film/air interface, and the y axis oriented perpendicular to the layer stack. A single, cylindrical defect is placed at the thickness midpoint of the top HfO₂ layer, its size taking on values in accordance with data in Fig. 70.30.

Heat *generation* inside the defect is based on uniform volume absorption of the laser light at a *constant rate* throughout the 1-ns laser pulse. We assume that *non-Beer's law* of absorption holds for these defects that are much smaller than a "skin depth." For an absorbed fraction β of the light normally incident on the defect of length d , the heat-generation rate inside the defect becomes

$$H_{\text{gen}} = \beta \Phi_{\text{th}} / d\tau, \quad (2)$$

where Φ_{th} is the damage-threshold fluence (5.6 J/cm²) and τ is the 1-ns laser-pulse duration. As reliable absorption-coefficient data are absent for most nanoscale absorbers, β was fixed at a value of 0.1. For a value of $d = 27$ nm, this yields a heat-generation rate $H_{\text{gen}} = 2 \times 10^{20}$ J/m³ s. To evaluate the relative influence of various other parameters, such as the

Figure 70.29

(a) Conventional laser-damage craters formed by 351-nm radiation in the HfO₂/SiO₂ film. Crater-depth distributions in the HfO₂/SiO₂ system and the electric field E of the incident light for different laser-fluence conditions: (b) ~5% above threshold fluence, (c) 10%–15% above threshold fluence, and (d) ~20% above threshold fluence.

thermal-conductivity *anisotropy* and defect size and type, this H_{gen} value was initially kept constant. Later, when agreement between empirically obtained AFM results and simulations was sought, H_{gen} values were parametrically varied.

Material Properties

Two types of absorbing defects were considered:

1. *Metallic Hf nanoclusters* with bulk-metal thermal parameter values for K (thermal conductivity), H (enthalpy), and T_m (melting point);

2. *Off-stoichiometric HfO₂ film* with the following parameters:

K – HfO₂ *thin-film* thermal conductivity,
 H – HfO₂ *bulk* enthalpy, and
 T_m – HfO₂ *bulk* melting point.

The absorbing-defect thermal parameter values are listed in Table 70.III.

Thin-film material property values are summarized in Table 70.IV, where K_x , K_y , and K_z are the Cartesian-

Figure 70.30
Model geometry employed for the finite-element analysis of heat transfer in the $\text{HfO}_2/\text{SiO}_2$ system.

Table 70.III: Absorbing-defect thermal parameters.

Defect Type	K ($\text{W/m}^\circ\text{K}$)	H (J/kg)	T_m ($^\circ\text{K}$)
Hf – metal	17.2	Hf – bulk ^a	2506
HfO_2	0.1	HfO_2 – bulk ^a	3085

^aReference 3.

Table 70.IV: Thin-film thermal parameters.

Material	K_y ($\text{W/m}^\circ\text{K}$)	K_x, K_z ($\text{W/m}^\circ\text{K}$)	H (J/kg)	T_m ($^\circ\text{K}$)
HfO_2	0.10	0.10γ	Bulk ^a	3085
SiO_2	0.25	0.25γ	Bulk ^a	1986

$\gamma = K_x/K_y$; $K_x = K_z$; $\gamma_1 = 0.1$; $\gamma_2 = 0.2$; $\gamma_3 = 0.5$; $\gamma_4 = 1$.

^aReference 3.

coordinate components of the thermal conductivity and γ is the anisotropy parameter for the thermal conductivity, i.e., $\gamma = 1$ for perfect isotropy. In these simulations, $\gamma = 0.1$ was chosen for the most anisotropic case. Also, in distinguishing metallic defects from HfO_2 -type defects, *thin-film anisotropy* values were used for HfO_2 -defect properties.

Latent heat of fusion was accounted for in these calculations by enthalpy H .

Modeling Results and Discussion

Instantaneous temperature contours in 2-D space (Fig. 70.31) show the temperature-field evolution within the

Figure 70.31

Temperature contours $[T(^{\circ}\text{K})]$ obtained for the case of large HfO_2 defect and $\gamma = 0.1$ at various times after the beginning of the laser pulse: (a) 1 ns, (b) 10 ns, (c) 20 ns, and (d) 40 ns.

film at various times after the beginning of the laser pulse. These contours represent the case of a “large” HfO_2 defect ($R = 20 \text{ nm}$; $d = 26.7 \text{ nm}$) located in the top HfO_2 layer. Note that in these contours the air/film interface is located at the *bottom* of each grid and the successive film layers stack upward. One finds [Fig. 70.31(a)] that at the end of the excitation pulse the heat generated is essentially confined to the defect itself, within which the temperature reaches a value of $2.7 \times 10^4 \text{ }^{\circ}\text{K}$. As time progresses, the locus of the peak temperature moves from the defect center toward the air/film interface [Fig. 70.31(b)]. As a consequence, 17 ns after the end of the pulse, the interface reaches a temperature of $1.5 \times 10^4 \text{ }^{\circ}\text{K}$. On the nanosecond time scale considered in Figs. 70.31(a)–70.31(d), *convective* heat transport at this interface remains physically irrelevant and, in effect, renders this interface an adiabatic boundary, fostering the modeled temperature buildup.

Critical to the comparison with empirical crater information are the lateral temperature contours, defining the lateral scale of the melting process. One can see that in the highly anisotropic ($\gamma = 0.1$) heat-conduction case represented in these figures, *lateral melt-front propagation* is fairly restricted—the melt radius remains $< 30 \text{ nm}$ inside HfO_2 , which is nearly a factor of 2 smaller than observed during AFM mapping of craters.

To determine the vertical-melt propagation through the film stack, Fig. 70.32 shows the *vertical* temperature around a “large” HfO_2 defect as a function of depth (here the air/film interface is located at the origin) at various times after the end of the laser pulse. From Fig. 70.32 it becomes apparent that within only 3 ns after the pulse, melting starts inside the SiO_2 layer (note the broken, vertical-layer-demarcation line between HfO_2 and SiO_2) and, by 10 ns after the pulse, reaches nearly halfway into the SiO_2 layer. Eventually the system starts to cool down, which sets the limit on the vertical-melt propagation. Similarly, Figs. 70.33 and 70.34 show the maximum temperature reached at any instant during the process at various depths below the defect as a function of conductivity anisotropy and defect type, respectively. For all anisotropy ratios chosen (Fig. 70.33), the SiO_2 layer experiences melting to some extent, up to $\sim 75\%$ of the entire layer. The same strong dependence of melting on anisotropy also holds for the “large” Hf-metal defect (Fig. 70.34), which at the fixed, identical heating rate is completely molten and itself isotropic ($T_m = 2506 \text{ }^{\circ}\text{K}$). In Figs. 70.33 and 70.34 there is a small, systematic error. The code recognizes and accounts for the change in conductivity anisotropy upon melting, but during the cool-down phase reinstates anisotropy as if the material had recovered initial crystallinity. Preliminary simulations show that this error in no way alters the seminal features or conclusions.

G4044

Figure 70.32

Vertical (along the y axis of the model) temperature profile kinetics (1 to 10 ns after the beginning of the laser pulse) for a "large" HfO₂ defect.

G4045

Figure 70.33

Maximum-temperature $T_{\max}$ vertical profiles as a function of conductivity anisotropy γ for a "large" HfO₂ defect case.

G4122

Figure 70.34

Maximum-temperature $T_{\max}$ vertical profiles as a function of conductivity anisotropy γ for a "large" metallic Hf defect case.

Ignoring the differences in absorption between a metallic and an off-stoichiometric defect and simply assuming that, over the laser pulse length, *equal amounts of energy* have been deposited in either one, the relative importance of *defect* heat capacity and internal heat conductivity for energy transfer to the surrounding film is outlined in Fig. 70.35. Due to its higher conductivity, the metallic defect attains a *uniform* temperature and delivers heat rapidly to its surface. This causes the surrounding HfO₂ film to reach higher temperatures than if the defect consisted of off-stoichiometric HfO₂. At the HfO₂/

SiO₂ interface this difference has been mediated away by the poor conductivity of the dielectric film. The melting behavior of the SiO₂ layer is insensitive to the nature of the defect.

The issue now remains how, for equal energy volume density at the end of an excitation pulse, defects of various sizes (i.e., of various surface areas) affect the melting behavior of the host layer as well as that of the next-nearest neighboring layer. For the three sizes (small/medium/large) considered, the volumes scale as 1:6:37. Figures 70.36 and 70.37 plot the maximum temperatures reached inside the HfO₂ and SiO₂ layers for the three cases—Fig. 70.36 representing a dielectric defect and Fig. 70.37 a metallic defect. For the same reason as given in the discussion of Fig. 70.35 (i.e., the higher thermal conductivity of the metal), the metallic defects shed heat effectively *during* the laser pulse, providing for larger maximum-temperature differences between the defects with the various surface-to-volume ratios than for defects of dielectric composition. At threshold, neither dielectric nor metallic, *small* defects contain enough heat to melt the SiO₂ layer, indicating a critical defect size for melting to occur (at threshold) inside the SiO₂ layer. Tables 70.V and 70.VI list the melt-penetration depth ℓ_m normalized, for comparison convenience with AFM measurements, to the HfO₂/SiO₂ layer interface for both types of medium and large defects, as well as ℓ_m dependence upon a range of conductivity anisotropies. The emphasis here is on weak anisotropies: for medium defects, no melt events are recorded, and even for large defects the melt penetration into SiO₂ is very shallow.

G4046

Figure 70.35

Maximum-temperature $T_{\max}$ vertical profiles as a function of defect type. Metallic Hf versus off-stoichiometric HfO_2 .

G4123

Figure 70.36

Maximum-temperature $T_{\max}$ vertical profiles as a function of the defect size for an HfO_2 defect case.

G4124

Figure 70.37

Maximum-temperature $T_{\max}$ vertical profiles as a function of the defect size for a metallic Hf defect case.

By raising the value for H_{gen} and simulating fluence conditions *above threshold*, one may address the question of how high a temperature a given defect must reach to generate damaging melt-penetration depths comparable to experimentally observed ones. For small- and medium-sized metallic defects the corresponding melt-penetration depths ℓ_m and defect peak temperatures are listed in Tables 70.VII and 70.VIII as a function of increasing H_{gen} . From these tables it becomes immediately apparent that experimentally measurable melt penetration occurs only for very high defect temperatures. For either defect size a temperature of 3 to 4×10^4 °K is required. Such temperatures are alarming inasmuch as the physical mechanisms so far ignored in these simulations start to play a role in the overall laser-damage model. They also point toward the inability of the current model to simulate laser damage by *thermal-conductivity-dominated* energy transfer alone.

Table 70.V: Melt-penetration depth ℓ_m (nm) values, large-defect case, as a function of parameter γ .

Defect Type	$\gamma = 0.1$	$\gamma = 0.2$	$\gamma = 0.5$	$\gamma = 1.0$
HfO_2	47	35	18	9
Metallic	49	38	24	14

$$T_{\max}(\text{HfO}_2) = 27,000 \text{ K}; T_{\max}(\text{met}) = 30,000 \text{ K}.$$

Two mechanisms that spring immediately to mind have the potential for significantly altering the kinetics of laser damage by providing sources for energy absorption from the laser field *at a distance from the original defect*: thermoionically emitted electrons and blackbody radiation in both the IR and, more importantly, the UV portion of its spectrum. Still, during the excitation pulse, when the thermal energy is mostly bottled up inside the defect and its temperature is highest, free carriers directly emitted into or ionized by UV photons within the surrounding dielectric layer stack provide starting electrons for "avalanche-type" or other energy-transfer processes at locations that, in effect, constitute an electronic *ballooning* of the defect. Table 70.IX emphasizes this aspect by listing the fraction of integrated blackbody radiation with photon energy above the HfO_2 bandgap as a function of defect temperature, emitted by a heated defect into its surrounding. From Table 70.IX we see that 10% of the emitted blackbody spec-

trum lies above 5.6 eV for temperatures as low as 1×10^4 °K. At the temperatures predicted by the current simulations this fraction increases to nearly 90%. The potential significance of such a UV ionization mechanism was pointed out in 1978 by Manenkov and his group.⁴ Based on this earlier work, the Grenoble Group⁵ evaluated the role of such preionization in single-layer films on melting and mechanical breaking of the film matrix.

Such a "larger-effective-defect" mechanism is also a plausible answer to the earlier-mentioned discrepancy between the calculated lateral-melt penetration range and the empirical AFM data. While part of this discrepancy can be assuaged by choosing a less-severe conductivity anisotropy, the lateral seeding of heating sources by ionizing radiation is an equally provocative concept whose merit must be tested in expanded simulations.

Table 70.VI: Melt-penetration depth ℓ_m (nm) values, medium-defect case, as a function of parameter γ .

Defect Type	$\gamma = 0.1$	$\gamma = 0.2$	$\gamma = 0.5$	$\gamma = 1.0$
HfO_2	13	5	0	0
Metallic	14	6	0	0

$T_{\max}(\text{HfO}_2) = 25,000$ K; $T_{\max}(\text{met}) = 17,000$ K.

Table 70.VII: Melt-penetration depth ℓ_m , as a function of heat-generation rate H_{gen} , small-defect case.

H_{gen} (W/m ³)	$T_{\max}$ (°K)	ℓ_m (nm)
2×10^{20}	8551	0
4	18760	1
5	22110	4
6	24750	6
7	31890	9
8	37750	12
10	42990	15

Table 70.VIII: Melt-penetration depth ℓ_m , as a function of heat-generation rate H_{gen} , medium-defect case.

H_{gen} (W/m ³)	T_{max} (°K)	ℓ_m (nm)
1×10^{20}	8652	2
2	16600	14
3	27070	23
4	36970	30
5	44880	35

Table 70.IX: Fraction of thermal radiation from the absorbing defect with photon energy exceeding HfO₂ bandgap (5.6 eV).

T (°K)	F_{ion} (%)
10,000	10
20,000	57
30,000	80
40,000	89

Summary

In considering thermal conduction as the only energy transport mechanism, the current simulations show that relevant comparison with empirical damage morphology necessitates such high defect temperatures ($T > 30,000^\circ\text{K}$) that other physical mechanisms start to attain dominance. These calculations point toward a mechanism in which the thermally trapped defect (thermally trapped on a nanosecond time scale) emits ionizing radiation into the surrounding film matrix. In response to such ionization, a volume larger than the defect itself becomes absorptive to the laser field and mediates direct energy transfer to the film at locations distant from the defect.

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Signatures of Target Performance and Mixing in Titanium-Doped Target Implosions on OMEGA

In recently published work we discussed the signatures of mixing¹ and of target performance² provided by incorporating a titanium-doped layer within the shell of imploding targets. We report here on the results of an experiment on OMEGA using such targets. We observe the predicted absorption features formed by the titanium layer: absorption lines due to $1s-2p$ transitions in titanium ions of incomplete L shell, as well as K -edge absorption in cold titanium. We also observe oscillations due to extended x-ray absorption fine structure (EXAFS) above the Ti K edge. The observed spectra are simulated by a radiative-transport model that uses *LILAC* profiles and *OPLIB* opacity values. The observed core temperature and shell areal density are found to fall short of the one-dimensional predicted values. We attribute this shortfall to a lack of irradiation symmetry, which leads to hydrodynamic instability and mixing [distributed phase plates (DPP's) were not used in this experiment, although the SSD modulator was activated]. The introduction of a sharp density jump at the interfaces of the titanium layer introduces a Rayleigh-Taylor (RT) unstable layer, which makes the target more unstable. Signatures of mixing include the emission of the He- α line of Ti^{+20} due to titanium migrating to the core, as well as the EXAFS spectrum indicating cold titanium close to peak compression. The spectral signatures due to the titanium layer are shown to provide very useful signatures of target performance.

A titanium layer was placed within a polymer-shell target, between a mandrel and an overcoat layer. The overcoat is much thicker than the laser-ablated layer to prevent emission of titanium lines from laser heating. Likewise, the titanium layer was sufficiently far from the inner-shell interface to ensure that the titanium lines are not emitted in the absence of mixing.¹ The titanium layer is thus cold for the most part and should give rise to absorption lines.³ The measured area enclosed within the absorption lines yields the areal density of the titanium and, by implication, of the total compressed shell. The intensity distribution of the absorption lines yields the temperature of the cold titanium layer. Additionally, the intensity slope of the continuum at higher photon energies yields the core temperature. These parameters are deduced from the measured spectra

and compared with simulation predictions. Highly performing targets require only a low percentage ($\sim 1\%$) of titanium doping to give significant absorption;² the need to use a pure-titanium layer in this experiment is an indication of lower compression.

The sharp density jump at the interfaces of the titanium layer becomes unstable during acceleration and deceleration; however, a thin layer cannot support unstable modes of long wavelength. A reduction factor multiplying the RT growth rate, due to the finite thickness of an unstable layer, has been derived by Landau and Lifshitz.⁴ One-dimensional code simulations of shot 8207 (see below) show that during the acceleration phase (when the shell and Ti layer are still thin) the reduction factor ranges from ~ 0.2 to ~ 0.5 for modes $\ell = 10$ to $\ell = 50$. Thus, the contribution of the titanium interface to the overall RT growth should be small.

In this article we concentrate on a single shot on the OMEGA laser⁵—shot No. 8207. The target was a polymer shell of $882\text{-}\mu\text{m}$ inner diameter and $4.6\text{-}\mu\text{m}$ thickness, coated with a $0.3\text{-}\mu\text{m}$ layer of titanium, a $15\text{-}\mu\text{m}$ layer of polymer, and finally a $0.1\text{-}\mu\text{m}$ layer of aluminum as a shinethrough barrier.⁶ The fill gas was atmospheric air. The 19.9-kJ laser pulse had a 1-ns flat top (to within $\pm 5\%$) with 150-ps rise and fall times. SSD⁷ was used in this experiment, but not DPP's.⁸

Experimental Results

Spectrally dispersed images were obtained with a Kirkpatrick-Baez x-ray microscope of resolution $\leq 5\text{ }\mu\text{m}$, fitted with a dispersing grating.⁹ The images show a highly nonuniform core of $\sim 60\text{-}\mu\text{m}$ FWHM, indicating an unstable implosion. This is further indicated by the spectral results discussed below.

The x-ray spectrum in the range of ~ 4 to 9 keV was recorded by a Ge(111) crystal spectrometer fitted with a $100\text{-}\mu\text{m}$ -wide spatially resolving slit. The slit acts to discriminate core radiation from coronal emission. The calibrated x-ray spectrum obtained with a Ge(111) diffracting crystal is shown in Fig. 70.38. Ti absorption lines are seen around 4.6 keV as well

as absorption above the Ti *K* edge at 4.966 keV. These absorption features take place when the radiation emitted by the hot, compressed core is absorbed by the colder Ti layer within the shell. The calibration procedure is explained in more detail in Ref. 10. For the calibration of the Kodak DEF film we relied on the measurements and modeling of Henke *et al.*,¹¹ for the calibration of the Ge(111) crystal we relied on the measurements of Burek and Yaakobi,¹² which agree well with the Darwin-Prins model calculations of Henke *et al.*¹³ Figure 70.39 shows the measured and calculated crystal reflectivity. It should be noted that the crystal sensitivity is a weak function

Figure 70.38

Calibrated, time-integrated x-ray spectrum obtained on shot 8207. Ti absorption lines are seen around 4.6 keV and the Ti *K* edge at 4.966 keV. These absorptions take place when the radiation emitted by the hot, compressed core is absorbed by the colder Ti layer within the shell.

Figure 70.39

Measured¹² and calculated¹³ integrated reflectivity of Ge(111) crystal.

of energy in the range of interest (4 to 8 keV). In Fig. 70.40 we magnify the part of the spectrum showing various absorption features that are of particular interest in this experiment. As explained below, the lines around 4.6 keV are absorbed within a cool titanium layer ($T_e \sim 200$ to 500 eV), whereas radiation above the Ti *K* edge at 4.966 keV is absorbed by much colder titanium. The EXAFS oscillations also indicate cold titanium. On the other hand, the Ti He- α line is emitted in a high-temperature region (~ 1 keV).

Figure 70.40

Part of the spectrum of Fig. 70.38, where absorption features appear. The lines around 4.6 keV are absorbed within a cool titanium layer ($T_e \sim 200$ to 500 eV), whereas radiation above the Ti *K* edge at 4.966 keV is absorbed by much colder titanium. The EXAFS oscillations also indicate cold titanium. On the other hand, the Ti He- α line is emitted in a high-temperature region.

K-Edge Absorption

The spectrum shows a drop in intensity at the *K* edge of cold titanium at 4.966 keV. This indicates the existence of cold titanium around the time of peak compression, when the observed core emission of continuum radiation is significant. To estimate the temperature of the absorbing titanium we note that for each electron ionized from a titanium atom, the *K* edge shifts to higher energies. A simple method for estimating the energy of these edges is discussed in Ref. 14. When the 12 *M*-shell electrons in Ti are removed, the edge shifts progressively from 4.966 keV up to about 5.2 keV; when the eight *L*-shell electrons are removed, the edge shifts progressively from ~ 5.2 to ~ 6.2 keV. The sharpness of the *K*-edge absorption indicates that no more than a few *M*-shell electrons are ionized, indicating a temperature lower than ~ 10 to 20 eV. The EXAFS spectrum discussed below indicates a still lower temperature.

The drop above the *K* edge corresponds to an absorption by a cold titanium layer of an areal density of 1.3 mg/cm²

(equivalent to a 2.9- μm -thick solid titanium foil). This is 9.3 times higher than the original areal density of the titanium layer, 0.135 mg/cm². Using this value of areal density we can derive the spectrum emitted by the core before being attenuated by the cold titanium layer. This is simply given by $I(E)/\exp[-\rho\Delta r \tau(E)]$, where $I(E)$ is the observed spectrum and $\tau(E)$ is the absorption (or opacity) of cold titanium; the result is shown in Fig. 70.41. The continuum is indeed seen to fall along a smooth exponential curve whose slope indicates a core electron temperature of 0.94 keV. This could have also been deduced from the high-energy part of the continuum in Fig. 70.38, where the attenuation by the cold titanium is negligible.

Figure 70.41

The calculated core emission before being absorbed by the cold titanium layer. The core electron temperature obtained by the continuum slope is 0.94 keV.

EXAFS Spectrum

Above the K edge the spectrum shows oscillations due to EXAFS. EXAFS¹⁵ is observed in solids (even amorphous ones) and is due to the interference of the photoelectron waves emitted from the absorbing ion and those reflected from the neighboring ions. It was previously observed in laser-heated flat targets¹⁶ but not in imploded targets. The analysis of the EXAFS oscillations yields the interparticle distance and, hence, the density. EXAFS data for solid titanium have been published¹⁷⁻¹⁹ and can be used for comparison.

The theory of EXAFS¹⁵ yields an expression for $\chi(k) = \mu(k)/\mu_0(k) - 1$, where $\mu(k)$ is the absorption (or opacity) and $\mu_0(k)$ is the absorption of the isolated atom (i.e., without the EXAFS oscillations). The wave number of the ejected photoelectron k is given by $\hbar^2 k^2 / 2m = E - E_K$, where

E is the absorbed photon energy and E_K is the energy of the K edge. The expression for $\chi(k)$ is

$$\chi(k) = \sum_j N_j F_j(k) \exp[-2\sigma^2 k^2 - 2R_j/\lambda(k)] \times \sin[2kR_j + \phi_j(k)] / kR_j^2, \quad (1)$$

where N_j is the number of atoms surrounding the absorbing atom at a distance R_j . The backscattering amplitude $F_j(k)$ and phase shift factor $\phi_j(k)$ for titanium were taken from the detailed calculations by Teo and Lee,²⁰ and the mean free path of the ejected electron in titanium, $\lambda(k)$, was taken from calculations of Blanche *et al.*¹⁷ The vibration amplitude σ^2 in the Debye-Waller factor and the interparticle distance R_j will be adjustable parameters in fitting the experimental EXAFS spectrum to Eq. (1).

Figure 70.42 shows the calculated absorption $\mu(k)$ of the cold titanium layer in shot 8207 (shown above to have an areal density of 1.3 mg/cm²), compared with an EXAFS spectrum of a 6- μm titanium foil absorber, obtained with synchrotron irradiation.¹⁹ The latter curve was multiplied by the ratio of thicknesses (2.9/6) in order to normalize it to the effective thickness (2.9 μm) of the absorbing layer in shot 8207. For both curves, the absorption was reduced by the much smaller absorption due to L -shell electrons, given by the absorption below the K edge (thus the absorption curves equal zero below the edge). The oscillations in the present result have a slightly larger period (thus a smaller R) and larger amplitude than the solid-titanium oscillations; both are consistent²¹ with slightly

Figure 70.42

Comparison of the absorption due to the cold titanium layer in shot 8207 and in a 6- μm -thick solid titanium foil, irradiated by synchrotron radiation. The latter curve was normalized to the effective thickness (2.9 μm) of the titanium layer in shot 8207.

higher than solid density in our case. The more gradual rise above the K edge in the curve obtained here is partly the result of finite spectral resolution and partly due to the ionization of a few M -shell electrons.

The isolated-atom absorption curve $\mu_0(k)$ is given by a smooth average through the EXAFS oscillations; thus, the ratio $\chi(k) = \mu(k)/\mu_0(k) - 1$ retains the changes in the absorption due only to the EXAFS oscillations. Before fitting the experimental EXAFS spectrum to Eq. (1), a Fourier transformation is applied (yielding the radial charge distribution around the absorbing atom) to single out the contribution of only the first term in Eq. (1), that of the closest shell of atoms around the absorbing atom. Also, rather than transforming $\chi(k)$, the quantity $k^3\chi(k)$ is normally chosen,¹⁵ putting more weight on the part of the spectrum away from the K edge (i.e., on higher k values). This choice is advantageous because Eq. (1) is imprecise at small k values; also, the effect of uncertainty in the value of E_K (due to valence-band energy structure) diminishes with increasing k values. The range of 2 to 10 $\text{\AA}^{-1}$ was used for the analysis of $k^3\chi(k)$. Below 2 $\text{\AA}^{-1}$ the curve is negligibly small; above $k \sim 10 \text{\AA}^{-1}$ the EXAFS spectrum is obscured by the rise of the Ti He- β line (albeit weak) at 5.582 keV and its dielectronic satellites.

The result of Fourier transforming the quantity $k^3\chi(k)$ (using a Hanning window²² to avoid spurious oscillations) is seen in Fig. 70.43. Only one peak, at 2.2 $\text{\AA}$, has a significant contribution. This distance is not exactly equal to the interparticle distance because of the effect of the phase factor $\phi_j(k)$ in

Figure 70.43

The result of Fourier transforming the $k^3\chi(k)$ curve obtained from Fig. 70.42. The peak at 2.2 $\text{\AA}$ corresponds to the shell of ions closest to the ion absorbing the photon.

Eq. (1). We next filter out this peak (extending from ~ 1 to $\sim 3 \text{\AA}$) by multiplying the curve of Fig. 70.43 by an appropriate super-Gaussian curve, and finally apply an inverse Fourier transformation back to the k space. The resulting curve (Fig. 70.44) was fitted by the first term of Eq. (1), multiplied by k^3 . As mentioned above, the functions $F(k)$, $\lambda(k)$, and $\phi_j(k)$ for titanium have been calculated in great detail, and thus the two remaining parameters, σ^2 and R , serve as adjustable parameters in the fit; no additional normalization was applied. Figure 70.44 shows the best fit, with $\sigma^2 = 0.013 \text{\AA}^2$ and $R = 2.6 \text{\AA}$. The parameter R depends primarily on the oscillation frequency of $k^3\chi(k)$, whereas σ^2 depends primarily on the amplitude decay rate at large k values (the amplitude rises at small k values because of the k^3 factor). Additionally, σ^2 determines the location and magnitude of the maximum in the $k^3\chi(k)$ curve. It should be emphasized that no amplitude normalization was applied to obtain this fit.

In solid titanium each atom is surrounded by two sub-shells of six atoms each, at distances of 2.889 $\text{\AA}$ and 2.940 $\text{\AA}$. The unresolved peak in Fig. 70.43 thus corresponds to $N = 12$ backscattering atoms at an average distance of 2.6 $\text{\AA}$ [hence the choice $N = 12$ in Eq. (1)]. This indicates a titanium layer of density ~ 1.4 higher than that of solid titanium, i.e., a density of 6.4 g/cm³. The Debye-Waller parameter σ (the amplitude of ion thermal oscillations) is proportional to $T_i^{1/2}$. By comparison with room-temperature EXAFS spectra of titanium, the spectrum of Fig. 70.44 corresponds to an ion temperature below ~ 1 eV. The values of density and temperature deduced here from the EXAFS spectrum are similar to those found with

Figure 70.44

Fitting Eq. (1) to the experimental $k^3\chi(k)$ curve, obtained by the inverse Fourier transform of the curve in Fig. 70.43. Two adjustable parameters were used (interparticle distance R and Debye-Waller amplitude σ^2) with no additional normalization.

an EXAFS measurement of laser-compressed foils.¹⁶ However, the ion temperature deduced here is lower than the electron temperature estimated above ($T_e < 10$ to 20 eV) from the sharpness of the K -edge absorption.

Absorption Lines

The absorption lines seen in Fig. 70.40 indicate that part of the titanium shell is ionized, in addition to the cold part giving rise to the K -edge and EXAFS absorptions. The application of absorption lines to determine the properties of the compressed shell was discussed in detail in Ref. 2. We show in Fig. 70.45 an expanded view of the absorption-line manifold obtained on shot 8207, compared with opacity calculations for $T = 300$ eV and $\rho = 6.4$ g/cm³. The calculations were taken from the LTE opacity library OPLIB,²³ which was also used in the simulations described in the next section. The validity of assuming LTE populations in the calculation of absorption lines has been discussed in Ref. 2. The absorption lines around 4.6 keV correspond to transitions of the type $1s-2p$ in titanium ions of an incomplete L shell: Ti^{+13} to Ti^{+20} , formed when core-emitted continuum traverses the cold titanium layer. The designation Li in Fig. 70.45 stands for lithium-like titanium ion, and likewise for the other designations. Each peak contains several lines that are unresolved mostly because of broadening due to the finite source size. For example, the peak marked C (carbon-like) consists of 35 transitions of the type $1s^2 2s^2 2p^2 - 1s 2s^2 2p^3$. The intensity distribution of the absorption peaks depends mostly on the temperature but also on the density. For example, the cases where the dominant peaks are the C-like or the N-like peaks (as in shot 8207) correspond to a temperature that changes from ~ 200 to ~ 500 eV for an assumed density in the range of 1 to 50 g/cm³. Thus, with no

knowledge of the density, a comparison as in Fig. 70.45 can be used to obtain only a rough estimate of the temperature in the cold shell. The energies calculated by OPLIB are approximate, and Fig. 70.45 shows them to be slightly different than the measured values.

In addition to absorption lines, Fig. 70.45 shows the emission of the He- α line of Ti^{+20} . The emission of He-like lines from a doped layer that is sufficiently removed from the interface was shown¹ to be a signature of mixing. Indeed, code simulations described in the next section show that these lines *should not have been seen* in the spectrum, in the absence of mixing.

Next we deduce the shell areal density from the absorption lines, using the relation²

$$\int \ln[I_0(E)/I(E)] dE = \rho \Delta r \int k(E) dE. \quad (2)$$

The left-hand term is an integral over the measured absorption-line manifold [if film density is proportional to $\ln(I)$, this term equals the area enclosed within the absorption lines in a film-density plot]. Although the opacity integral depends on the temperature and density of the absorbing region, it has been shown² that using the observed intensity distribution of the absorption peaks enables us to deduce the areal density without having to know these parameters. We calculate the opacity integral $\int k(E) dE$ for titanium, using the OPLIB opacity library,²³ for a wide range of temperatures and densities. When the results are grouped by the peak of highest opacity, the integral values fall within relatively narrow ranges, which can be understood as follows: Ionization increases with increasing

Figure 70.45

Comparison between the measured Ti absorption lines from shot 8207 and the opacity calculated by OPLIB. The OPLIB calculation is for $T = 300$ eV and $\rho = 6.4$ g/cm³. The designation Li stands for $1s-2p$ transitions in Li-like Ti, and likewise for the other designations.

temperature but decreases with increasing density (because the three-body recombination rises faster with density than other, two-body processes). Thus, a given ionization level is obtained for pairs of (T, ρ) values where both increase simultaneously. Since opacity increases with density but decreases with temperature, the integrated opacity for such pairs tends to be reasonably constant. We show in Fig. 70.46 the calculated opacity integral, arranged according to the strongest opacity peak (marked Li for cases where the lithium-like peak is the strongest, etc.). Thus, simply knowing the strongest absorption peak narrows down the uncertainty in the integrated opacity and, hence, in the deduced areal density. For example, for cases where the C-like peak is the strongest, the uncertainty in the integrated opacity is $\pm 25\%$. Applying Fig. 70.46 and Eq. (2) to the spectrum of shot 8207, we find $\rho\Delta r \sim 0.4 \text{ mg/cm}^2$ for the titanium layer with a temperature in the range of ~ 200 to 500 eV . Adding the areal density found above for the cold titanium (1.3 mg/cm^2), we find a total $\rho\Delta r$ for the titanium layer of 1.7 mg/cm^2 . In adding the areal densities we have assumed that the line- and K -edge absorptions occur at the same time, signifying a temperature gradient within the titanium layer. This is indeed indicated by Fig. 70.40. Suppose the two absorptions take place at two different times: the K edge is absorbed at time t_1 and the lines are absorbed at a time t_2 . The observed fluence above the K edge ($\sim 1.2 \times 10^{16} \text{ keV/keV}$) consists then of three parts: (a) the fluence that was emitted at t_1 but survived the K -edge absorption, (b) the fluence emitted

at t_2 , and (c) the fluence emitted at other times. Thus, the fluence at t_2 must be smaller than the observed fluence above the K edge. However, this level of fluence is much smaller than the depth of the absorption lines ($\sim 4 \times 10^{16} \text{ keV/keV}$), indicating insufficient fluence to absorb the lines in those times.

Using the measured $\rho\Delta r$ of the titanium layer we can estimate the $\rho\Delta r$ of the entire compressed shell. Figure 70.47 shows the ratio $Q = (\rho\Delta r)_{\text{total}}/(\rho\Delta r)_{\text{Ti}}$ calculated by the one-dimensional hydrodynamic code *LILAC*²⁴ for shot 8207 as a function of the compression ratio (defined as r_0/r , where r is the radius of the gas-shell interface and r_0 is its initial value). Also shown is the result of a simple geometrical model where the unablated part of the shell is assumed to be incompressible as it implodes inward. Since charge-collector measurements (and simulations) indicate that an overcoat thickness of $\sim 9 \mu\text{m}$ is ablated, the model uses the inner $10.9 \mu\text{m}$ as the unablated part of the shell. Thus, a lower bound $Q \sim 10$ can be used, giving $\rho\Delta r > 17 \text{ mg/cm}^2$ for the compressed shell.

Figure 70.46

Area under the opacity-spectrum peaks of titanium (corresponding to the absorption-line manifold around 4.6 keV). For temperatures in the range of 300 to 1000 eV and densities in the range of 1 to 50 g/cm^3 , the points are grouped according to the strongest absorption peak in the opacity spectrum (Li stands for cases where the lithium-like peak is the strongest, and likewise for the other designations).

Figure 70.47

The ratio of $\rho\Delta r$ for the total shell and for the titanium layer. The model assumes that the unablated shell implodes incompressibly. The compression ratio is defined as r_0/r , where r is the radius of the gas-shell interface and r_0 is its initial value.

It should be pointed out that the use of time-integrated spectra yields an underestimate of the areal density because emission during times when absorption is insignificant diminishes the absorption contrast (typically² by a factor of ~ 2).

Comparison with Simulations

In this section we compare the observed spectrum with that predicted by *LILAC*.²⁴ The Ti spectral signatures are shown to reflect the reduced target performance that was a consequence of the poor irradiation uniformity in this experiment.

The simulated spectra were computed by applying a radiation-transport model¹ to *LILAC*-generated profiles, using the OPLIB opacity library. As discussed elsewhere,^{1,2} using the LTE opacity library is adequate for simulating the continuum radiation and the absorption lines. However, this would not be true for simulating the emission of helium-like and hydrogen-like titanium lines.

Figure 70.48 shows the comparison for the higher-energy part of the spectrum, and Fig. 70.49 the lower-energy part, where the absorption features appear. In both cases the simulated spectra are integrated in time up to the time marked in

Figs. 70.48 and 70.49. The integration up to 2 ns is essentially an integration over the whole implosion. The simulated curves show some flattening below ~ 7 keV due to shell attenuation,^{25,9} and only higher-energy radiation is unattenuated and can be used to determine the core temperature. However, the experimental curve does not show this flattening, indicating a less-compressed shell. It also shows a cooler core than predicted: $T = 0.94$ as compared with $T = 1.8$ keV. The maximum temperature in the simulation ($T \sim 3$ keV) was higher than that indicated by the time-integrated continuum slope ($T = 1.8$ keV), and thus it can be expected that the peak temperature in this experiment is ≈ 1 keV. The lower tempera-

E8361

Figure 70.48

Comparison of the measured and simulated x-ray continuum spectra for shot 8207. The simulated curves are integrated up to the marked times. The slopes yield the core electron temperature.

E8362

Figure 70.49

Comparison of measured and simulated x-ray spectra for shot 8207, in the region showing absorption features. The simulated curves are integrated up to the marked times.

ture and compression result in a much lower core-radiation fluence than predicted. Figure 70.48 shows that the implosion in this experiment falls short of the 1-D code predictions; in particular, the target fails to undergo the final compression stage that is predicted to occur between ~ 1.86 ns and peak compression at ~ 1.9 ns.

Similar conclusions can be drawn from the comparison in Fig. 70.49. The simulated time-integrated spectrum (< 2 ns) shows much deeper absorption lines than observed, indicating a higher areal density, and little K -edge absorption, indicating a higher shell temperature. In fact, the saturated absorption of the simulated absorption lines shows that almost all radiation emitted after ~ 1.84 ns at the wavelengths of these lines is absorbed. As time progresses, the simulated spectra show a gradual shift from K -edge absorption to line absorption, indicating heating of the titanium layer by the core. Also, the absorption-line manifold is seen to shift toward higher ionizations, again indicating an increase in shell temperature. By contrast, the experimental curve shows (a) a lower total absorption, indicating a lower areal density [see Eq. (2)], and (b) prominence of K -edge absorption over line absorption, indicating a cooler shell. As in Fig. 70.48, the results show that the implosion is effectively aborted after ~ 1.86 ns. At intermediate times (~ 1.84 to 1.86 ns) K -edge and line absorptions occur simultaneously, which was shown above to be the case in the experiment.

We also note absorption features above ~ 5 keV in the simulated spectra but not in the measured spectrum. These are K -edge absorption in ionized Ti, as well as transitions of the type $1s-3l$ (K -shell to M -shell transitions) in various titanium ions.

The measured $\rho\Delta r$ of the titanium layer was found above to be ~ 1.7 mg/cm². The same $\rho\Delta r$ value in the simulation occurs at 1.7 ns, consistent with the conclusion of incomplete compression suggested above. At 1.7 ns the predicted titanium density is 11 g/cm³, compared with the value 6.4 g/cm³ deduced from the EXAFS spectrum. The predicted electron temperature is ~ 50 eV, compared with the estimate $T < 10$ to 20 eV from the sharpness of the K -edge absorption. The simulated spectra also predict no emission of helium-like Ti lines (even though LTE opacity modeling tends to overestimate their intensity).

The appearance of the He- α line in the experiment (see Fig. 70.45) is clear evidence of shell-core mixing,¹ which transports titanium into high-temperature regions. Titanium

can be driven out to the ablation region (during acceleration) or into the compressed core (during deceleration). To determine which is the case in this experiment, we examine the spatial profile of the He- α line, as measured by the crystal spectrometer. Figure 70.50 shows the profiles (in the target plane) at the wavelength of the He- α line, as well as at a nearby wavelength; the net He- α emission is given by their difference. The sharpness of the core emission profile indicates a smaller core than the slit width (100 μ m), consistent with the 60- μ m width measured by the x-ray microscope. Figure 70.50 clearly shows that most of the He- α emission comes from the core, indicating mixing during the deceleration phase. The peak core temperature (≥ 1 keV) is sufficiently high to excite the He- α line. The observed intensity of the He- α line can be used to obtain a rough estimate of the amount of titanium reaching the core due to mixing. To that end, we repeat the simulations of Fig. 70.49 while assuming that a small fraction of the titanium mass is uniformly spread throughout the shell layers lying inward of the original titanium layer. Figure 70.51 shows the result for the case of mixing 3% of the total titanium mass into the core. The He- α line (as well as its adjacent satellite line) now appears in the calculated spectrum. A comparison is made between the experimental spectrum and the simulated spectrum integrated up to 1.86 ns, as suggested by Figs. 70.48 and 70.49. In calculating the spectrum of Fig. 70.51, a correction was made to account for the LTE calculations' overestimate of the intensity of the emission lines. A steady-state, non-LTE model was used to calculate the relative populations of the lithium-like, helium-like, and hydrogen-like Ti species, from which the intensity of the major

Figure 70.50
Spatial profile of emission at the wavelength of the Ti He- α line (4.75 keV), and at a nearby wavelength. The difference (i.e., the net He- α emission) indicates mixing of titanium into the compressed core.

Figure 70.51

Simulation of mixing: the simulated curve (integrated up to 1.86 ns) assumes that 3% of the total titanium mass is uniformly spread throughout the shell layers inward from the original titanium layer. Note the appearance of the He- α line due to mixing, in agreement with the time-integrated, measured spectrum.

lines due to collisional excitation follows. The required ionization, recombination, and excitation rates were taken from Ref. 26.

The reduced target performance evidenced in Figs. 70.48 and 70.49 as well as the presence of mixing is the result of inadequate irradiation uniformity when not using phase plates to smooth the laser irradiation, as well as the result of the density jump at the titanium layer interfaces. The titanium spectral signatures studied in this work should prove very useful in tracking the improved target performance with improved laser uniformity.

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Dephasing Time of an Electron Accelerated by a Laser Pulse

In a recent paper¹ we described the motion of electrons in the electromagnetic field of a circularly polarized laser pulse propagating through a plasma. Electrons that are in front of the pulse initially can be accelerated to high energies and extracted easily. Although this direct acceleration scheme is less than ideal because the pulse can generate a parasitic wake, its simplicity is noteworthy. The wake fields produced by short pulses have been observed recently,^{2,3} and future experiments will study the interaction of electron bunches with these wake fields. One would only need to change the timing of an electron bunch in these experiments to test the scientific feasibility of direct acceleration. In this article we study the dephasing time of an electron accelerated by a pulse of infinite width to determine the propagation time and plasma length required to observe direct acceleration.

In the following sections, the trajectory of a charged particle is determined analytically for a representative pulse profile; the dephasing time of an accelerated particle is determined; its dependence on the speed, length, and intensity of the pulse along with the injection energy of the particle is studied in detail; and finally, the main results are summarized.

Particle Motion in a Planar Field

The motion of a particle, of charge q and mass m , in an electromagnetic field is governed by the equation⁴

$$d_\tau(u_\mu + a_\mu) = u^\nu \partial_\mu a_\nu, \quad (1)$$

where u^μ is the four-velocity of the particle divided by c , τ is the proper time of the particle multiplied by c , and a^μ is the four-potential of the field multiplied by q/mc^2 . The metric four-tensor $g_{\mu\nu} = \text{diag}(1, -1, -1, -1)$.

For a planar field a^μ has y and z components that are functions of t and x . It is convenient to denote the transverse (two-vector) component of a^μ by $\mathbf{a}$, the transverse component of u^μ by $\mathbf{v}$, and the longitudinal components of u^μ by γ and u . In this notation, the transverse component of Eq. (1) is

$$d_\tau(\mathbf{v} + \mathbf{a}) = 0. \quad (2)$$

For a particle that is in front of the pulse initially, and is not moving transversely,

$$\mathbf{v} = -\mathbf{a}. \quad (3)$$

By using Eq. (3), one can rewrite the longitudinal components of Eq. (1) as

$$d_\tau \gamma = \partial_t(v^2/2), \quad d_\tau u = -\partial_x(v^2/2). \quad (4)$$

For a circularly polarized field

$$a^\mu = (0, 0, a \cos \phi, a \sin \phi) / \sqrt{2}. \quad (5)$$

We assume that the phase $\phi = t - sx$, where $s < 1$ is the inverse phase speed of the pulse, and the amplitude a is a function of $\psi = t - rx$, where $r > 1$ is the inverse group speed of the pulse. Equations (3)–(5) were solved in Ref. 1 for a particle that is at rest initially. The solution of these equations for a particle that is moving initially is similar. Since the ponderomotive potential $v^2/2$ is independent of ϕ , it follows from Eqs. (4) that

$$d_\tau(r\gamma - u) = 0. \quad (6)$$

By combining Eq. (6) with the definition of γ , one can show that

$$\gamma = \frac{r(r\gamma_0 - u_0) \mp \omega}{r^2 - 1}, \quad (7)$$

$$u = \frac{(r\gamma_0 - u_0) \mp r\omega}{r^2 - 1},$$

where

$$\omega = \left[(r\gamma_0 - u_0)^2 - (r^2 - 1)(1 + v^2) \right]^{1/2}. \quad (8)$$

In Eqs. (7) the $-$ sign applies to the case in which $\gamma > ru$, which corresponds to a particle that is moving more slowly than the pulse, and the $+$ sign applies to the case in which $\gamma < ru$, which corresponds to a particle that is moving more quickly than the pulse. By using the fact that $1 = \gamma_0^2 - u_0^2$, one can rewrite Eq. (8) in the convenient form

$$\omega = \left[(\gamma_0 - ru_0)^2 - (r^2 - 1)v^2 \right]^{1/2}. \quad (9)$$

For the case in which $\gamma_0 = 1$ and $u_0 = 0$, Eqs. (7) and (9) reduce to the corresponding equations of Ref. 1. A particle that is moving more slowly than the pulse initially will be repelled by the pulse if $\omega = 0$. For this to happen the pulse intensity must equal the repelling intensity¹

$$a^2 = 2(\gamma_0 - ru_0)^2 / (r^2 - 1), \quad (10)$$

in which case the gain in particle energy¹

$$\delta\gamma = 2(\gamma_0 - ru_0) / (r^2 - 1). \quad (11)$$

For completeness, a covariant analysis of particle motion in a circularly polarized field is given in **Appendix A**, and a brief description of particle motion in an elliptically polarized field is given in **Appendix B**.

Equations (3), (7), and (9) define u^μ as a function of ψ . By combining the equation $d_\tau\psi = \gamma - ru$ with Eqs. (7), one can show that

$$d\tau/d\psi = \pm 1/\omega(\psi), \quad (12)$$

where the $+$ sign applies to the case in which $\gamma < ru$ and the $-$ sign applies to the case in which $\gamma > ru$. If the solution of Eq. (12) can be inverted, u^μ can be expressed as an explicit function of τ .

To illustrate the particle motion we consider the simple profile

$$a(\psi) = e \sin(\pi\psi/2lr), \quad (13)$$

where e^2 is the peak intensity of the pulse and l is its full-width at half-maximum. For this profile

$$\omega(\psi) = (\gamma_0 - ru_0) [1 - m^2 \sin^2(\pi\psi/2lr)]^{1/2}, \quad (14)$$

where

$$m^2 = (r^2 - 1)e^2 / 2(\gamma_0 - ru_0)^2 \quad (15)$$

is the ratio of the pulse intensity to the repelling intensity.

When $m < 1$, the pulse overtakes the particle completely. In this case ψ varies between 0 and $2lr$, and the solution of Eq. (12) is

$$\tau(\psi) = [2lr/\pi(\gamma_0 - ru_0)] F(\pi\psi/2lr, m), \quad (16)$$

where F denotes the incomplete elliptic integral of the first kind, of modulus m .⁵ It follows from Eqs. (7) and (9) that

$$t(\psi) = \frac{r(r\gamma_0 - u_0)\tau(\psi) - \psi}{r^2 - 1}, \quad (17)$$

$$x(\psi) = \frac{(r\gamma_0 - u_0)\tau(\psi) - r\psi}{r^2 - 1}.$$

The particle motion is illustrated in Fig. 70.52 for the case in which $\gamma_p = 30$, $\gamma_0 = 7$, and $e^2 = 7$ [The Lorentz factor γ_p is defined in the first of Eqs. (22)]. In Fig. 70.52(a) the phase, normalized to lr , is plotted as a function of time, normalized to $\gamma_p^2 l$. As the particle is accelerated by the front of the pulse, the rate of phase slippage decreases. However, since the peak intensity of the pulse is lower than the repelling intensity, the particle speed never equals the pulse speed and the pulse overtakes the particle. As the particle is decelerated by the back of the pulse, the rate of phase slippage increases. It is evident from Fig. 70.52(a) that the deceleration time equals the acceleration time. In Fig. 70.52(b) the longitudinal momentum is plotted as a function of the normalized time. Although the particle speed never exceeds the pulse speed, the energy

associated with the transverse particle motion allows the particle momentum to exceed the pulse momentum. Because the longitudinal momentum is a symmetric function of time, the deceleration distance equals the acceleration distance.

When $m > 1$, the particle is repelled by the pulse. In this case ψ increases from 0 to $(2lr/\pi)\sin^{-1}(1/m)$ as the particle ascends the ponderomotive potential and decreases from $(2lr/\pi)\sin^{-1}(1/m)$ to 0 as the particle descends the ponderomotive potential. The solution of Eq. (12) is

$$\tau(\psi) = \begin{cases} [2lr/\pi m(\gamma_0 - ru_0)]F(\theta, 1/m), \\ [2lr/\pi m(\gamma_0 - ru_0)][2K(1/m) - F(\theta, 1/m)], \end{cases} \quad (18)$$

where

$$\theta(\psi) = \sin^{-1}[m \sin(\pi\psi/2lr)] \quad (19)$$

Figure 70.52

Particle trajectory for the case in which $\gamma_p = 30$ and $\gamma_0 = 7$. The pulse intensity $e^2 = 7$ is slightly lower than the repelling intensity [Eq. (10)]. (a) Normalized phase ψ/lr plotted as a function of the normalized time $t/\gamma_p^2 l$; (b) longitudinal momentum u plotted as a function of the normalized time.

and K denotes the complete elliptic integral of the first kind, of modulus m .⁵ The first form of Eq. (18) applies to the ascent and the second form applies to the descent. Equations (17) apply to both the ascent and descent, provided that τ is defined by Eqs. (18) and (19). The particle motion is illustrated in Fig. 70.53 for the case in which $\gamma_p = 30$, $\gamma_0 = 7$, and $e^2 = 10$. In Fig. 70.53(a) the normalized phase is plotted as a function of the normalized time. Initially, the pulse overtakes the particle and the rate of phase slippage is positive. Since the peak intensity of the pulse is higher than the critical intensity, the particle is accelerated until its speed equals the pulse speed and the rate of phase slippage is zero. Subsequently, the particle overtakes the pulse and the rate of phase slippage is negative. The descent time is longer than the ascent time because the time dilation associated with a particle moving faster than the pulse is larger than that associated with a particle moving slower than the pulse. In Fig. 70.53(b) the longitudinal momentum is plotted as a function of the normal-

Figure 70.53

Particle trajectory for the case in which $\gamma_p = 30$ and $\gamma_0 = 10$. The pulse intensity $e^2 = 10$ is slightly higher than the repelling intensity [Eq. (10)]. (a) Normalized phase ψ/lr plotted as a function of the normalized time $t/\gamma_p^2 l$; (b) longitudinal momentum u plotted as a function of the normalized time.

ized time. Since the particle does not reach the peak of the pulse, the x component of the ponderomotive force is always positive and the longitudinal momentum of the particle increases monotonically. Because the longitudinal momentum is an asymmetric function of time, the descent distance is longer than the ascent distance.

Dephasing Time of an Accelerated Particle

The previous analysis shows how an intense pulse repels a charged particle that is in front of the pulse. The relation between the pulse intensity, the particle injection energy, and the gain in particle energy was studied in Ref. 1. In this section the time required for the pulse to catch and repel the particle, and, subsequently, for the particle to outrun the pulse, is studied. This time is referred to as the dephasing time and is denoted by T . The distance traveled by the particle during the dephasing time is referred to as the dephasing distance and is denoted by X . It follows from Eqs. (17) and (18) that

$$T = \frac{4lr^2(r\gamma_0 - u_0)K(1/m)}{\pi m(r^2 - 1)(\gamma_0 - ru_0)} \quad (20)$$

and $X = T/r$. For future reference, notice that $K(1/m) \rightarrow \log [4/(m^2 - 1)^{1/2}]$ as $m \rightarrow 1$ and $K(1/m) \rightarrow \pi/2$ as $m \rightarrow \infty$.⁵

Formula (20) for the dephasing time exhibits a complicated dependence on the pulse intensity and speed and the initial particle momentum. One can gain insight into the underlying physics by performing a pulse-frame analysis of the acceleration process. In the notation of Ref. 1, $'$ denotes a pulse-frame quantity, the subscript A denotes the initial position of the particle, and B denotes the position at which the particle is repelled.

The pulse-frame energy and momentum of the particle are related to the laboratory-frame energy and momentum by the equations

$$\gamma' = \gamma_P \gamma - u_P u, \quad u' = \gamma_P u - u_P \gamma, \quad (21)$$

where

$$\gamma_P = r/(r^2 - 1)^{1/2}, \quad u_P = 1/(r^2 - 1)^{1/2}. \quad (22)$$

In these equations γ_P is the Lorentz factor associated with the pulse speed $1/r$ and $u_P = (\gamma_P^2 - 1)^{1/2}$. If one uses the linear

group speed of the pulse to estimate the Lorentz factor, $\gamma_P = \omega_0/\omega_e$, where ω_0 is the carrier frequency of the pulse and ω_e is the electron-plasma frequency.⁶

In the pulse frame v^2 is time-independent. It follows from the first of Eqs. (4) that γ' is constant and, hence, that $(u')^2 + v^2 = (u'_A)^2$. Since $dx'/dt' = u'/\gamma'_A$, it follows that

$$T' = 2\gamma'_A \int_{x'_B}^{x'_A} \frac{dx'}{[(u'_A)^2 - v^2(x')]^{1/2}}. \quad (23)$$

In Eq. (23) the factor of 2 arises because the pulse-frame descent time equals the pulse-frame ascent time. The factor of γ'_A arises because of the difference between proper time and pulse-frame time. Provided one ignores the distinction between momentum and velocity, the integral in Eq. (23) represents the ascent time of a nonrelativistic particle in the potential well $v^2(x')/2$. In the pulse frame $a = e \sin(-\pi x'/2l')$, where $l' = \gamma_P l$. For this profile

$$T' = 2\gamma'_A (2l'/\pi) (\sqrt{2}/e) K(\sqrt{2}u'_A/e). \quad (24)$$

The factor of $2l'/\pi$ arises because the ponderomotive force associated with the pulse is inversely proportional to the pulse length. Although Eq. (24) is complicated, the origin of each factor is well understood.

In the pulse frame the particle begins and ends its interaction with the pulse at point A . Since $X' = 0$, it follows that

$$T = \gamma_P T', \quad X = u_P T'. \quad (25)$$

Notice that $X = T/r$, as stated after Eq. (20). It follows from Eqs. (21) and (22) that $r(r\gamma_0 - u_0)/(r^2 - 1) = \gamma_P \gamma'_A$ and $4lr/\pi m(\gamma_0 - ru_0) = 2(2l'/\pi)(\sqrt{2}/e)$. Thus, Eq. (24) and the first of Eqs. (25) agree with Eq. (20).

It is convenient to define the normalized dephasing time

$$\bar{T} = (4\sqrt{2}\gamma'_A \pi e) K(\sqrt{2}u'_A/e), \quad (26)$$

which is the dephasing time divided by $\gamma_P^2 l$. The factor of l was due to the inverse dependence of the ponderomotive force on the pulse length. One factor of γ_P was due to the Lorentz

transformation of the pulse length from the laboratory frame to the pulse frame; the other factor was due to the Lorentz transformation of the dephasing time from the pulse frame to the laboratory frame. These factors do not depend on the physical origin or shape of the potential well in which the particle moves. Thus, it was inevitable that they should be the same as the factors that control the dephasing time of an electron in the laser beat-wave accelerator⁷ or the laser wake-field accelerator.^{7,8} For completeness, a brief analysis of the particle motion and dephasing time associated with these indirect acceleration schemes is given in **Appendix C**.

Just as a pulse-frame analysis of the acceleration process fosters insight into the dephasing time, so also does it foster insight into the energy gain. In the pulse frame the particle energy is constant, and the final particle momentum has the same magnitude as the initial particle momentum and the opposite sign: $\delta\gamma' = 0$ and $\delta u' = 2|u'_A|$. It follows from these results and Eqs. (21) that

$$\delta\gamma = 2u_P|u'_A|, \quad (27)$$

in agreement with Eq. (11).

The normalized dephasing time is plotted as a function of pulse intensity in Fig. 70.54 for the case in which $\gamma_P = 30$. In Fig. 70.54(a) the injection energy $\gamma_A = 7$. The solid line denotes the exact dephasing time [Eq. (26)], and the broken line denotes the approximate dephasing time $2\sqrt{2}\gamma'_A/e$. For the chosen values of γ_P and γ_A the approximate dephasing time is $6.4/e$. When the pulse intensity is close to the repelling intensity, the particle lingers near the peak of the pulse and the dephasing time is long. As the pulse intensity increases, point *B* moves toward the front of the pulse and the dephasing time decreases. In the high-intensity regime this decrease is gradual. Since the pulse energy located behind point *B* is wasted, there is little to be gained by using pulse intensities that exceed the critical intensity by more than a factor of 2. Since the injection energy is constant, so also is the energy gain [Eq. (27)]. In Fig. 70.54(b) the injection energy

$$\gamma_A = \gamma_P \mu - \left[(\gamma_P^2 - 1)(\mu^2 - 1) \right]^{1/2}, \quad (28)$$

where $\mu = (1 + e^2/4)^{1/2}$ is a measure of the pulse intensity. This choice of injection energy ensures that the repelling intensity is one-half of the pulse intensity. For this injection

energy $u'_A = -e/2$, $\gamma'_A = (1 + e^2/4)^{1/2}$, and the saturation time is $1.7(1 + 4/e^2)^{1/2}$, independent of γ_P . In the low-intensity regime the dephasing time is long because the ponderomotive force is weak. In the high-intensity regime the dephasing time is almost independent of pulse intensity because the increase in ponderomotive force that accompanies an increase in pulse intensity is offset by the corresponding decrease in injection energy. It follows from Eq. (27) and the preceding discussion that the energy gain equals $u_P e$. As the pulse intensity increases, the energy gain increases and the required injection energy decreases.

The normalized dephasing time is plotted as a function of injection energy in Fig. 70.55 for the case in which $\gamma_P = 30$. In Fig. 70.55(a) the pulse intensity $e^2 = 10$. The solid line denotes the exact dephasing time [Eq. (26)], and the broken line denotes the approximate dephasing time $2\sqrt{2}\gamma'_A/e$. For the chosen values of γ_P and e the approximate dephasing time is

Figure 70.54

Normalized dephasing time [Eq. (26)] plotted as a function of pulse intensity for the case in which $\gamma_P = 30$. (a) The particle injection energy $\gamma_A = 7$. (b) The particle injection energy [Eq. (28)] ensures that the repelling intensity [Eq. (10)] is one-half of the pulse intensity.

P1717

Figure 70.55
Normalized dephasing time [Eq. (26)] plotted as a function of particle injection energy for the case in which $\gamma_P = 30$. (a) The pulse intensity $e^2 = 10$. (b) The pulse intensity [Eq. (29)] is twice the repelling intensity [Eq. (10)].

$0.89\gamma'_A$. When the injection energy is close to the repelling energy, the particle lingers near the peak of the pulse and the dephasing time is long. As the injection energy increases, point B moves toward the front of the pulse and the dephasing time decreases. In the high-energy regime the dephasing time is almost independent of the injection energy because $\gamma'_A \approx 1$. The energy gain decreases as the injection energy increases. In Fig. 70.55(b) the pulse intensity

$$e^2 = 4(u'_A)^2 \quad (29)$$

is twice the repelling intensity and the dephasing time is $1.7\gamma'_A/|u'_A|$. In the low-energy regime the dephasing time is almost independent of the injection energy because $\gamma'_A \approx \gamma_P$ and $u'_A \approx -\gamma_P$. The ratio $\gamma'_A/|u'_A|$ is almost independent of γ_P . In the high-energy regime the dephasing time is long and the energy gain is small because $\gamma'_A \approx 1$ and $|u'_A| \ll 1$.

Summary

The motion of an electron in the electromagnetic field associated with a circularly polarized laser pulse of infinite width was studied analytically. When the pulse intensity is lower than the repelling intensity [Eq. (10)], the pulse overtakes the electron completely. When the pulse intensity is higher than the repelling intensity, the electron is repelled by the pulse and eventually outruns it. The time taken for the electron to outrun the pulse is called the dephasing time and is the product of two terms. The first term is $\gamma_P^2 l$, where γ_P is the Lorentz factor associated with the pulse speed and l is the pulse length. The second term [Eq. (26)] depends on the pulse intensity, the pulse shape, and the electron injection energy. As a rough guideline, the second term is of order unity unless the pulse intensity is close to the repelling intensity. For a pulse of finite width, an electron that is not close to the pulse axis initially will be expelled from the pulse by the radial component of the ponderomotive force.⁹ Further work is needed to quantify this snowplow effect.

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Appendix A: Covariant Analysis of the Particle Motion in a Planar Field

The resolution of Eq. (1) into longitudinal and transverse components is facilitated by the introduction of the four-vector k^μ , which is defined by the equation $\psi = k^\nu x_\nu$, and the four-vector l^μ , which is defined by the equations $l^\nu l_\nu = -k^\nu k_\nu$, $l^\nu k_\nu = 0$, and $l^\nu a_\nu = 0$, where a^μ is the transverse four-potential of a planar field of arbitrary polarization. In the laboratory frame $k^\mu = (1, r, 0, 0)$ and $l^\mu = (r, 1, 0, 0)$. By using these four-vectors, one can write

$$x_\mu = y_\mu + \psi k_\mu / k^\nu k_\nu + \theta l_\mu / l^\nu l_\nu, \quad (A1)$$

where y_μ is transverse and $\theta = l^\nu x_\nu$. In a similar way, one can write

$$u_\mu = v_\mu + (k^\nu u_\nu) k_\mu / k^\nu k_\nu + (l^\nu u_\nu) l_\mu / l^\nu l_\nu, \quad (A2)$$

where $v_\mu = dy_\mu/d\tau$ is transverse, $k^\nu u_\nu = d_\tau \psi$, and $l^\nu u_\nu = d_\tau \theta$.

The transverse component of Eq. (1) is

$$d_\tau(v_\mu + a_\mu) = 0, \quad (\text{A3})$$

from which it follows that

$$v_\mu(\tau) = v_\mu(0) + a_\mu(0) - a_\mu(\tau). \quad (\text{A4})$$

Equation (A4) is the analog of Eq. (3).

By using Eq. (A4), one can rewrite the right side of Eq. (1) as $-\partial_\mu(v^\nu v_\nu/2)$. Since $v^\nu v_\nu$ was assumed to be a function of ψ , $\partial_\mu = k_\mu d_\psi$. It follows from these results that the longitudinal components of Eq. (1) are

$$d_\tau(k^\mu u_\mu) = -k^\mu k_\mu d_\psi(v^\nu v_\nu/2), \quad d_\tau(l^\mu u_\mu) = 0. \quad (\text{A5})$$

It follows from the second of Eqs. (A5) that

$$l^\mu u_\mu(\tau) = l^\mu u_\mu(0). \quad (\text{A6})$$

One way to obtain an expression for $k^\nu u_\nu$ is to use the identity $u^\nu u_\nu = 1$, which can be rewritten as

$$v^\nu v_\nu + (k^\nu u_\nu)^2 / k^\nu k_\nu + (l^\nu u_\nu)^2 / l^\nu l_\nu = 1. \quad (\text{A7})$$

It follows from Eq. (A7) that

$$[k^\nu u_\nu(\tau)]^2 = [l^\nu u_\nu(\tau)]^2 + k^\nu k_\nu [1 - v^\nu v_\nu(\tau)]. \quad (\text{A8})$$

Equations (A6) and (A8) are the analogs of Eqs. (7) and (8). By using the expression for $k^\nu u_\nu(0)$ that follows from Eq. (A8), and Eq. (A6), one can show that

$$[k^\nu u_\nu(\tau)]^2 = [k^\nu u_\nu(0)]^2 + k^\nu k_\nu [v^\nu v_\nu(0) - v^\nu v_\nu(\tau)]. \quad (\text{A9})$$

Equation (A9) is the analog of Eq. (9). Another way to obtain an expression for $k^\nu u_\nu$ is to change the independent variable in the first of Eqs. (A5) from τ to ψ . Since $d_\tau \psi = k^\nu u_\nu$, the first of Eqs. (A5) becomes

$$d_\psi[(k^\nu u_\nu)^2/2] = -k^\mu k_\mu d_\psi(v^\nu v_\nu/2), \quad (\text{A10})$$

from which Eq. (A9) follows.

Finally, since $v^\nu v_\nu$ is a function of ψ , Eqs. (A4), (A6), and (A9) express u_μ as a function of ψ . To express u_μ as a function of τ one must invert the solution of the phase equation

$$d\tau/d\psi = \pm 1/\omega(\psi), \quad (\text{A11})$$

where ω is the square root of the terms on the right side of Eq. (A9).

Appendix B: Guiding-Center Motion in a Planar Field

Equation (6) is valid when the v^2 terms in Eqs. (4) are independent of ϕ . To satisfy this condition we assumed that the field is circularly polarized and that the particle is in front of the pulse initially and not moving transversely. Equations (7), (9), and (10) follow from Eq. (6) and the definition of γ , which requires that

$$d_\tau(\gamma^2 - u^2 - v^2) = 0. \quad (\text{B1})$$

For the elliptically polarized field

$$a^\mu = (0, 0, a_y \cos \phi, a_z \cos \phi), \quad (\text{B2})$$

where $a_y = a\delta$ and $a_z = a(1 - \delta^2)^{1/2}$, the v^2 terms in Eqs. (4) are not independent of ϕ , and Eq. (6) is not valid. However, the particle motion is known to consist of a fast oscillation about a guiding center and a guiding-center drift that varies slowly. In a vacuum, the guiding-center motion is governed by the equation⁹

$$d_\tau \langle u_\mu \rangle = -\partial_\mu \langle a^\nu a_\nu \rangle / 2, \quad (\text{B3})$$

where $\langle \rangle$ denotes a ϕ -average and $\langle a^\nu a_\nu \rangle = -a^2/2$. We expect Eq. (B3) to provide a reasonable description of the guiding-center motion in a rarefied plasma, in which the phase speed of the field is slightly higher than the speed of light. Equation (B3) has associated with it the conservation equation

$$d_\tau \langle \langle u^\nu \rangle \langle u_\nu \rangle + \langle a^\nu a_\nu \rangle \rangle = 0. \quad (\text{B4})$$

Since the ponderomotive potential $a^2/4$ is independent of ϕ , it follows from Eq. (B3) that

$$d_\tau(r\langle\gamma\rangle - \langle u\rangle) = 0. \quad (\text{B5})$$

Equation (B5) is the analog of Eq. (6). Since $\langle v\rangle$ is constant, Eq. (B4) reduces to

$$d_\tau(\langle\gamma\rangle^2 - \langle u\rangle^2 - a^2/2) = 0. \quad (\text{B6})$$

Equation (B6) is the analog of Eq. (B1). Thus, for a particle that is in front of the pulse initially, $\langle\gamma\rangle$ and $\langle u\rangle$ are given by Eqs. (7) and (9), in which v^2 is replaced by $a^2/2$, and the repelling conditions are described by Eq. (10).

Appendix C: Particle Motion in a Planar Electrostatic Field

The four-potential of an electrostatic field can be written as

$$a_\mu = pk_\mu/k^\nu k_\nu + ql_\mu/l^\nu l_\nu, \quad (\text{C1})$$

where k^μ and l^μ were defined in **Appendix A**. We assume that a_μ is a function of ψ , from which it follows that $\partial_\mu = k_\mu d_\psi$. Since the electrostatic field is unaffected by the gauge transformation $a_\mu \rightarrow a_\mu + \partial_\mu b$, where b is an arbitrary function of ψ , p is redundant. In the Lorentz gauge $p = 0$.

By substituting decomposition (C1) in Eq. (1) and contracting the resulting equation with k^μ , one can show that

$$d_\tau(k^\mu u_\mu + p) = (k^\nu u_\nu) d_\psi p - (l^\nu u_\nu) d_\psi q. \quad (\text{C2})$$

Since $k^\nu u_\nu = d_\tau \psi$, the p terms in Eq. (C2) cancel, as they must do. By substituting decomposition (C1) in Eq. (1) and contracting the resulting equation with l^μ , one can show that

$$d_\tau(l^\mu u_\mu + q) = 0. \quad (\text{C3})$$

It follows from Eq. (C3) that

$$l^\mu u_\mu(\tau) = l^\mu u_\mu(0) + q(0) - q(\tau). \quad (\text{C4})$$

One way to obtain an expression for $k^\nu u_\nu$ is to use the identity

$u^\nu u_\nu = 1$, which can be rewritten as

$$[k^\nu u_\nu(\tau)]^2 = [l^\nu u_\nu(\tau)]^2 - l^\nu l_\nu. \quad (\text{C5})$$

Another way to obtain an expression for $k^\nu u_\nu$ is to solve Eq. (C2) directly. By changing the independent variable from τ to ψ , one can rewrite Eq. (C2) as

$$d_\psi \left[(k^\nu u_\nu)^2 / 2 \right] = d_\psi \left[(l^\nu u_\nu)^2 / 2 \right], \quad (\text{C6})$$

from which Eq. (C5) follows.

Since q is a function of ψ , Eqs. (C4) and (C5) express u_μ as a function of ψ . To express u_μ as a function of τ , one must invert the solution of the phase equation

$$d\tau/d\psi = \pm 1/\omega(\psi), \quad (\text{C7})$$

where ω is the square root of the terms on the right side of Eq. (C5).

In the wave frame $k'_\mu = (0, l, 0, 0)$ and $l'_\mu = (l, 0, 0, 0)$, where $l = (r^2 - 1)^{1/2}$. It follows from these results that $\psi = -lx'$, $k^\nu u_\nu = -lu'$, $l^\nu u_\nu = l\gamma'$, and $q = l\phi'$, where ϕ is the electrostatic potential. Thus, Eq. (C4) can be rewritten as

$$\gamma'(\tau) = \gamma'(0) + \phi'(0) - \phi'(\tau) = 0, \quad (\text{C8})$$

Eq. (C5) can be rewritten as

$$[u'(\tau)]^2 = [\gamma'(\tau)]^2 - 1, \quad (\text{C9})$$

and Eq. (C7) can be rewritten as

$$d\tau/dx' = \mp 1/\omega(x'), \quad (\text{C10})$$

where ω is the square root of the terms on the right side of Eq. (C9). The dephasing time of an accelerated particle can be determined from Eq. (C10) in a manner similar to that described previously. In particular, by considering the relations between laboratory-frame and wave-frame quantities, one can show that the dephasing time is proportional to $\gamma_W^2 \lambda$, where γ_W is the Lorentz factor associated with the phase speed of the wave and λ is the wavelength.

The potential associated with a large-amplitude plasma wave is described by elliptic functions. Simple formulas for the injection energy and energy gain associated with this potential were determined by Esarey and Piloff.¹⁰ The dephasing times associated with this and other potentials were studied by Teychenné, Bonnaud, and Bobin.^{11,12}

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The Effect of Increased Irradiation Uniformity on Imprinting by 351-nm Laser Light

The experimental program¹ at LLE supports the national inertial confinement fusion (ICF) effort by performing experiments on OMEGA² to investigate the requirements for attaining ignition using direct-drive targets on the National Ignition Facility. One of the primary challenges in direct-drive ICF is to minimize perturbations in the target that are created by nonuniformities in the drive laser. These imprinted perturbations can seed the Richtmeyer–Meshkov (RM) and Rayleigh–Taylor (RT) instabilities that can amplify target nonuniformities sufficiently large enough to destroy an implosion. Typical high-gain implosions are expected to experience unstable growth factors in excess of 500.

One of the primary experimental efforts is to reduce imprinting by controlling irradiation nonuniformities. Other methods being studied include modifications to the target (foam buffers) and general reduction of growth rates; these methods will be reported on in the future. Eventually, successful ICF implosions may require some level of each of these methods. Here, we report on planar-target experiments that quantify the effect of changes to irradiation nonuniformity on imprinting. In addition, we report on “perturbing-beam” experiments—a novel method to measure the times over which imprinting occurs.

Drive nonuniformities conventionally are expressed in terms of spherical-harmonic ℓ -modes that describe the number of perturbation wavelengths contained in the circumference of a spherical target. The amplitudes of the low-order modes ($\ell \leq 8$) are affected by the configuration, pointing, and focusing of the beams, and by beam-to-beam power balance. In contrast, imprinting occurs predominantly in the high-order modes ($\ell \geq 8$), which largely result from nonuniformities in the intensity distributions of the individual beams. To reduce these nonuniformities, beam-smoothing techniques for glass lasers have traditionally relied on the physical optics associated with optical arrays placed in each of the laser beams. Distributed phase plates³ (DPP’s) act to randomize the phase across the beam wavefront, producing a speckle pattern with high-frequency modulations. Smoothing by spectral dispersion⁴ (SSD)

produces a variable frequency across the wavefront, producing multiple modes that are displaced in the target plane and, when averaged over time, smooth the DPP speckle nonuniformities. The DPP’s transform the long-wavelength beam nonuniformities (typical of high-power lasers) to wavelengths sufficiently short enough that displacements of the speckle pattern (caused by SSD) smooth out most of the nonuniformities.

Distributed polarization rotators⁵ (DPR’s) change the beam polarization across the wavefront (and therefore across the DPP), thereby diminishing the interference responsible for the high-frequency modulations in the DPP pattern. Since beams of orthogonal polarization do not interfere, the uniformity can be increased in proportion to the square root of the number of polarization states one produces in the wavefront, which in the case of the current DPR’s is 2.

To assess the irradiation uniformity, equivalent-target-plane (ETP) images of single beams with various beam-smoothing techniques have been analyzed. Beams with (1) DPP’s only, (2) DPP and 1-D SSD (1.5 Å), and (3) DPP and 2-D SSD (0.6 Å × 1.5 Å) were compared and, when measured over the duration of a 1-ns pulse, had time-averaged nonuniformities of 98%, 25%, and 12%, respectively.⁶ A recent implementation of DPR’s for OMEGA is shown in Fig. 71.1. A wedge of KDP (a birefringent material) is inserted such that the beam polarization bisects the ordinary and extraordinary axes. Since the element is wedged, the two orthogonally polarized beams emerge from the KDP at different angles, thereby producing two speckle patterns that are shifted by $\approx 80 \mu\text{m}$ [see Fig. 71.1(a)]. Since the two beams are orthogonally polarized, the two speckle patterns do not interfere and, therefore, produce a $1/\sqrt{2}$ ($N = 2$ beams) reduction in nonuniformity. This increase occurs instantaneously and is therefore particularly important to direct drive because irradiation imprinting occurs early in the drive and in some cases before appreciable SSD smoothing takes place. Figure 71.1(b) compares the uniformity of single beams with DPP’s only (i.e., no SSD), with and without DPR’s. To produce these plots, the square of the Fourier transform of the ETP images is integrated azimuthally

Figure 71.1

(a) Schematic of a DPR (a KDP wedge in the initial implementation on OMEGA) and (b) nonuniformity spectrum on target with and without DPR's for an OMEGA beam with phase plates (DPP's) but no SSD.

at constant radii, i.e., constant spatial frequency. The graphs are the square of the contribution (σ_k^2) to total nonuniformity as a function of spatial frequency. (The upper limit in frequency is the diffraction limit of the OMEGA lenses.) These graphs indicate that the $1/\sqrt{2}$ factor applied to σ_{rms} for the 2-D SSD result (above) should produce a time-integrated uniformity of 8.5% rms. Contributions to nonuniformity from film noise prevented this comparison from being performed on the 2-D SSD beams. We expect to use electronic imaging to further study the highly uniform beams.

Imprinting Experiments

Several multifoil, planar-target experiments were performed on OMEGA to study hydrodynamic instabilities. These experiments, which began on Nova as a collaboration between Lawrence Livermore National Laboratory (LLNL) and LLE,⁷ study two aspects of direct-drive target stability: growth rates and irradiation imprinting. Using x-ray radiography, the target trajectories and the growth of pre-imposed mass perturbation were measured in accelerated targets. These experiments were well simulated by hydrocodes, providing confidence that both the energy coupling⁸ and unstable growth⁹ are well modeled for these experiments. This confirmed our experience in similar experiments conducted on Nova.

The effect of irradiation nonuniformity on laser imprinting was studied using unperturbed CH₂ targets driven by five overlapped UV beams at 2×10^{14} W/cm² in 3-ns square

pulses. The targets were backlit with x rays produced by a uranium backlighter irradiated at $\sim 2 \times 10^{14}$ W/cm² (using 12 beams). X rays transmitted through the target were imaged using a framing camera with 8- μ m pinholes filtered with 20 μ m of Be and 12 μ m of Al. This yielded highest sensitivity at an average photon energy of ~ 1.3 keV. Face-on radiographic images at different times were analyzed to measure the growth of perturbations seeded by imprinting. Since the initial perturbations are too small to be detected by this method, RT growth (driven by the target acceleration) is employed to amplify the imprint perturbations to detectable levels. Unfortunately, the addition of this highly nonlinear "amplifier" somewhat complicates interpretation of the results.

Figure 71.2 shows a series of radiographic images taken at 2.4 ns for each of four different laser configurations (all with DPP's): (a) no SSD, (b) 1-D SSD (1.5 Å), (c) 2-D SSD (0.6 Å $\times$ 1.5 Å), and (d) DPR's and 2-D SSD (0.6 Å $\times$ 1.5 Å). These configurations are listed in order of increasing on-target uniformity, which can be estimated by multiplying the time-integrated single-beam nonuniformities by $1/\sqrt{5}$ ($N=5$ beams), yielding 44%, 11%, 5.4%, and 3.8%, respectively, for the four configurations. Note, however, that at time equals zero, the first three cases present identical nonuniformities (i.e., the DPP speckle patterns) to the target. As time progresses, SSD provides time-averaged smoothing, the level of which depends upon the degree of SSD (bandwidth and number of dimensions). In contrast, the DPR's provide a $1/\sqrt{2}$ increase in

uniformity instantaneously (at time zero and all times) because two speckle patterns (from each beam) are continuously incident on the target.

Figure 71.3 depicts the nonuniformity "power spectra" obtained at different times from radiographic images of targets irradiated with various laser configurations. The spectra for the images in Fig. 71.2 (at 2.4 ns) are the top line in each graph. An estimate of the noise level for the diagnostic is shown as the gray regions. The nonuniformities depicted by these graphs result from acceleration-driven growth that is seeded by imprinting. Each point in these graphs is the azimuthal integral at constant radius (constant frequency) of the two-dimensional Fourier transform of the radiographic images. The summation

of each graph over all frequencies provides the square of the total rms deviation from a smooth envelope, i.e., the nonuniformities in that image. Note that, at a given time, the nonuniformities in these images decrease as the irradiation uniformity is improved. Since the pulse shape and drive intensities were the same, the growth rates are expected to be similar for all cases. The higher uniformity of the accelerated targets is, therefore, indicative of lower initial perturbation amplitudes, i.e., less imprinting. The smallest amount of imprinting occurs for the case with 2-D SSD and DPR's. Inferences of the actual imprint amplitudes will be made using hydrodynamic simulations that have been normalized using the growth-rate data and include instrument response functions and noise levels.

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Figure 71.2

Radiographic images of 25- μm CH_2 foils taken at 2.4 ns and irradiated by 3-ns UV pulses having a variety of laser conditions (all with DPP's): (a) no SSD, (b) 1-D SSD, (c) 2-D SSD, and (d) DPR's and 2-D SSD. The latter images show higher uniformity of the accelerated target.

E8399b

(Shot Nos. 8544, 8557, 8545, 8567)

Figure 71.3

Measured nonuniformity power spectra of the x-ray backlighter transmission through accelerated CH foils subjected to a variety of laser-irradiation conditions. Fourier analysis of various x-ray framing camera images is shown for the times indicated. The maximum wave number shown (400 mm^{-1}) corresponds to a spatial wavelength of 16 μm .

This analysis is performed on the optical depth of the target (the natural log of the intensity-converted images). Using calibration targets, it is estimated that ratios of local intensities in Fig. 71.2(a) indicate that many RT "bubbles" have nearly penetrated the foil. The power spectra belie this point because the Fourier amplitudes correspond to modulations of less than $1\text{ }\mu\text{m}$. This is because the target nonuniformities comprise a wide range of frequencies, so the individual amplitudes are relatively small. (This is in contrast to the single-mode experiments where the images are dominated by the imposed modulations.) Furthermore, one cannot estimate growth rates by tracking the amplitudes of individual modes at various times because the radiographic contrast of these measurements is sufficiently low that a significant amount of RT growth (linear and possible nonlinear) is needed to observe the target nonuniformities. This implies that saturation and possibly mode competition and coalescence dramatically affect the shape of the power spectrum. (In fact, many of our results exhibit a shift to longer wavelengths for late times.) Also, the MTF of the device has not been included in the measurements; if it were, the noise-dominated signal at shorter wavelength would appear as a dominant factor. To properly analyze these experiments we are investigating other image-analysis techniques and will use hydrodynamic simulations that include the instrument response. This is particularly important for spatial wavelengths $20\text{ }\mu\text{m}$ and lower.

The time dependence of imprinting is important for determining the requirements for beam-smoothing rates. Imprint-

ing occurs until sufficient coronal plasma is produced and laser nonuniformities (those with wavelength shorter than the distance from the critical to the ablation surfaces) are smoothed by that plasma. Simulations¹⁰ indicate that for pertinent ICF conditions, irradiation nonuniformities decouple from the ablation surface after 300 to 500 ps. As shown above, SSD reduces imprinting, yet the addition of DPR's shows continued reduction of imprinting. This indicates that imprinting occurs before this level of SSD smoothing is fully effective.

To understand the time over which imprinting occurs, a series of preliminary experiments were conducted on OMEGA before DPP's and SSD were available. In these experiments, five beams without DPP's accelerated initially smooth targets, and a sixth beam of smaller diameter was added at different times and with different energies. The relative effect of the sixth beam is gauged by observing the nonuniformities in the accelerated target. One can estimate the duration of imprinting by determining the perturbing-beam delay that no longer produces observable effects in the accelerated target.

For these experiments, the nominal drive beams were $1000\text{ }\mu\text{m}$ in diameter and the perturbing beam was $400\text{ }\mu\text{m}$. The drive intensity was $\sim 1 \times 10^{14}\text{ W/cm}^2$ in 1-ns Gaussian pulses. Qualitatively, dramatic differences can be seen in the radiographic images for different conditions of the perturbing beam. Figure 71.4 depicts images from two shots with different timing of the perturbing beam. The upper images [Fig. 71.4(a)] show that, when timed 150 ps before the drive pulses, the

Figure 71.4

Framing camera images at various times relative to the peak of the Gaussian 1-ns drive beams for two shots: (a) perturbing beam at -150 ps and (b) perturbing beam at $+300\text{ ps}$. The effect of the $400\text{-}\mu\text{m}$ -diam perturbing beam on the accelerated target uniformity is diminished when it arrives sufficiently late in time.

perturbing beam imprinted its features (smaller circle of nonuniformities) onto the target. The perturbing beam also produced similar effects when it was timed coincident (0 ps) with drive beams. Figure 71.4(b) shows the diminished effect of the perturbing beam when it is delayed by 300 ps with respect to the drive beams. This is likely the result of sufficient plasma formation (by the primary beams) to allow smoothing of the irradiation nonuniformities produced by the perturbing beam. In both cases, the perturbing beam had an intensity equal to the total drive-beam intensity, i.e., $\sim 1 \times 10^{14}$ W/cm². When the intensity of the perturbing beam was reduced to 7% of the drive beams' intensity, its effect was no longer observed in the radiographic images. The mottled patterns in Fig. 71.4 are the growth of features imprinted by the 1-mm drive beams. Similar features can be seen in the outer annulus of the images in Fig. 71.4(a). We expect that refinement of these experiments will provide insight into imprinting mechanisms and their time dependence, as well as provide information on the effectiveness of uniformity enhancements to the laser.

The next step in this effort will be to study the growth and mitigation of short-wavelength imprinting and to diagnose smaller amplitude perturbations. By relying less on the RT instability to amplify the perturbations, measurements performed at earlier times may provide better estimates of the spectrum and amplitude of imprinting. These measurements will require higher-resolution x-ray imaging (<10 μ m) and higher-contrast measurements and will be attempted using a Kirkpatrick-Baez microscope or curved-crystal x-ray imaging, and high-resolution imaging of a 250-eV backlighter. In addition, imprint mitigation and growth-rate reduction will be studied using target modifications such as foam buffers and high-Z dopants.

Summary

The mitigation of the deleterious effects of laser imprinting in direct-drive ICF targets is crucial to LLE's experimental program. The primary approach—increased irradiation uniformity—has been shown to reduce imprinting levels, particularly for a new beam-smoothing device called distributed polarization rotator (DPR). The perturbing-beam experiments presented here demonstrate a novel way to study the time dependence of imprinting—an important issue for the design of certain uniformity schemes. To study imprinting earlier in time and perturbations with shorter wavelengths, we are investigating other diagnostics for radiography, including a number of advanced x-ray and XUV diagnostics.

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Single-Mode Rayleigh–Taylor Growth-Rate Measurements with the OMEGA Laser System

Understanding the growth of perturbations due to hydrodynamic instabilities is important to the design of direct-drive targets for the National Ignition Facility (NIF). Direct-drive inertial confinement fusion (ICF) has the potential to be more efficient than indirect-drive ICF; however, there is the potential for greater growth of hydrodynamic instabilities due to the lower ablation velocities, which will prevent a NIF target from igniting. The primary instability of concern is the Rayleigh–Taylor (RT) instability, where a less dense fluid “supports” a more dense fluid and as such merits a thorough investigation. Direct-drive targets are subject to the RT instability twice during an ICF target implosion: first during the ablative acceleration phase of the implosion and second when the deceleration phase begins core assembly. Target imperfections and laser-irradiation nonuniformities act as amplitude seeds for the acceleration RT instability, while interior target imperfections and feedthrough of the ablation surface growth to the fuel–pusher interface act as amplitude seeds for the deceleration RT instability.

The importance of the RT instability problem is reflected in the number of experiments conducted^{1–5} to determine the RT growth rates in directly driven systems. Experiments of interest have used laser irradiation of 527 nm and 351 nm and have employed a number of beam-smoothing schemes. It is the development of these advanced laser-irradiation smoothing techniques like induced spatial incoherence (ISI)⁶ and smoothing by spectral dispersion (SSD)⁷ that has enabled the RT growth experiments to yield quantitative data. Without the beam-smoothing techniques, irradiation nonuniformity can dominate the experimental observables, even in initial-large-amplitude, single-mode, mass-perturbation experiments.

The RT growth rate at the ablation surface has typically been described by the Takabe⁸ or modified Takabe dispersion formula given by

$$\gamma = \alpha \cdot \sqrt{k \cdot g} - \beta \cdot k \cdot V_a$$

and

$$\gamma = \alpha \cdot \sqrt{\frac{k \cdot g}{1 + k \cdot L}} - \beta \cdot k \cdot V_a,$$

respectively, where γ is the growth rate, k is the wave number of the perturbation, g is the acceleration of the interface, V_a is the ablation velocity, α is a constant approximately equal to 0.9, β is a constant with an approximate value of 3, and L is the density scale length at the ablation surface. The importance of thermal conduction⁹ has been investigated and leads to a modification of the Takabe formula for the growth rate through changes in the values α and β that depend on the details of the thermal transport at the ablation surface.

Details of the dispersion formula for RT growth rates remain an active subject of research at this time. Previous measurements¹⁰ of the growth rate, conducted using x-ray radiography, have yielded values that are substantially lower than that calculated by either of the above formulae. The need for nonthermal electron transport has been proposed as the cause for these lower growth rates.

The results from a series of single-mode RT-instability growth experiments performed on the OMEGA laser system using planar targets are presented. Planar targets with imposed mass perturbations were accelerated using five to six 351-nm laser beams overlapped with a total overlapped intensity up to 2.5×10^{14} W/cm². Experiments were performed with both 3-ns ramp and 3-ns flat-topped temporal pulse shapes. The use of distributed phase plates (DPP's)¹¹ and SSD resulted in a laser-irradiation nonuniformity of 4%–7% over a 600- μ m-diam region defined by the 90% intensity contour.

The temporal growth of the modulation in optical depth was measured using through-foil radiography and was detected with an x-ray framing camera. The temporal-growth evolution of the measured variation in the optimal depth for both 31- μ m and 60- μ m wavelength perturbations was found to be in good agreement with *ORCHID* simulations when the experimental details, including noise, were included. In addition, it has been found that for current experiments care must be taken in

relating the actual ablation-surface RT growth rate from the experimental observables used in these experiments.

Experimental Configuration

A schematic of the experimental layout (Fig. 71.5) shows the orientation of the three foils and diagnostics used in these experiments. The foils are supported by a 140- μm boron-coated carbon fiber. Five to six beams are incident onto the 20- μm -thick CH-foil drive target, which is mounted on a 50- μm -thick Mylar washer with a 1-mm hole in the center. A 2-mm $\times$ 2-mm piece of 25- μm -thick uranium is mounted at the other end of the boron-coated carbon fiber 9 mm from the drive foil. Six to twelve beams were used to irradiate the uranium foil in order to generate the x rays needed to radiograph the accelerated foil. A third foil composed of 6- μm -thick aluminum was centered between the drive and backlighter foils to act as a bandpass filter for x rays between 1.0 and 1.5 keV. This prevents the very-low-energy x rays from the uranium target from preheating the accelerated foil.

The typical mass-modulated accelerated foil is composed of a 20- μm -thick CH foil that had perturbations imposed on the side irradiated by the laser. The 20- μm thickness was chosen because it represents about two attenuation depths for the 1.0- to 1.5-keV x rays used for radiography. Perturbation wavelengths of either 31 μm or 60 μm were imposed onto the foils with initial amplitudes of 0.45 μm and 0.5 μm , respectively. The face of the foil irradiated by the laser was overcoated with a 500-Å to 1000-Å layer of Al to prevent shinethrough.¹² (Foils used for acceleration measurements had no initial imposed mass perturbations.)

The primary diagnostic is an x-ray framing camera.¹³ The pinhole array is composed of eight 8- μm pinholes arranged in a checkerboard pattern to minimize interference from adjacent images. The framing camera used for the perturbation amplitude measurement has a magnification of 14.1 ± 0.1 with

the pinhole array located 24.8 mm from the rear surface of the accelerated foil. A second x-ray framing camera measured the spatial and temporal emissions of the backlighter at a magnification of 6 with a standard 16-pinhole array. Each pinhole is 10 μm in diameter. Data from the backlighter framing camera are shown in Fig. 71.6. The images are timed so that the backlighter reaches full emission in the end of the second strip (i.e., between images 3 and 4). Each image is separated by 50 ps, and the timing of the backlighter emission is known to about 50 ps. The gating pulses for uniform strips of each of the two framing cameras are recorded by an oscilloscope, and the time of each image relative to the drive beams is known to about 100 ps. The temporal resolution of both x-ray framing cameras was 80 ps. The current noise levels on the framing camera allow the instrument to measure perturbations with amplitudes $\geq 0.4 \mu\text{m}$. As a result, the early-time amplitude of the accelerated foils was at the detection limit and, therefore, difficult to resolve.

Laser Irradiation

Typical planar-foil experiments use two independent drivers so that separate pulse shapes could be used for the drive and backlighter beams. The beams were overlapped onto their respective targets with a radial displacement accuracy of 35 μm from the center of each foil.

The backlighter pulse shape used was a 3-ns-wide flat-topped pulse. Each backlighter beam was focused to a 1.5-mm-diam spot at the 5% intensity contour and had an energy of ~ 400 J for a total of 3.2 to 4.8 kJ on the uranium backlighter.

The two pulse shapes used for the drive beams are shown in Fig. 71.7. The first was a 3-ns ramp pulse, and the second was a 3-ns flat-topped pulse. Each of the drive beams was focused to a ~ 900 - μm -diam spot size (at the 5% intensity contour) and incorporated DPP's with an on-target intensity distribution measured to be $I \propto \exp[-(r/r_0)^4]$. The resultant five- to six-

Figure 71.5

Experimental schematic of the target and primary diagnostic. The target is constructed from three foils: a uranium backlighter, an aluminum debris shield, and a CH drive foil with an imposed perturbation. The primary diagnostic is an x-ray framing camera with a 14.1 magnification provided by 8- μm pinholes located 24.8 mm from the accelerated target.

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beam overlapped spot had a region of constant intensity of $\sim 600\text{-}\mu\text{m}$ diameter. The drive beams were temporally smoothed by 2-D SSD with a total bandwidth of 0.135 THz at 351 nm with an angular dispersion of $33\text{ }\mu\text{rad}/\text{\AA} \times 33\text{ }\mu\text{rad}/\text{\AA}$ (IR bandwidths of $1.5\text{ }\text{\AA} \times 0.7\text{ }\text{\AA}$). The total drive intensity for six irradiation beams peaked at $2.5 \times 10^{14}\text{ W/cm}^2$ for the ramp pulse and $2.0 \times 10^{14}\text{ W/cm}^2$ for the flat-topped pulse.

Figure 71.6

Framing camera image of backlighter x-ray emission. The four strips are timed from bottom to top and right to left. The values for t_0 reflect the trigger time of the strip relative to the start of the backlighter pulse.

Figure 71.7

Laser illumination pulse shapes used to accelerate the CH target. The 3-ns ramp pulse (dashed curve) has a peak intensity of $2.5 \times 10^{14}\text{ W/cm}^2$, and the 3-ns flat-topped pulse (solid curve) has an intensity of $2.0 \times 10^{14}\text{ W/cm}^2$.

The uniformity of a single beam with a DPP and SSD at the maximum bandwidth for these experiments was measured by imaging the focal spot onto UV-sensitive film and analyzing the intensity distribution.¹⁴ The resultant nonuniformity for a single beam was measured as 12% integrated over $\sim 1\text{ ns}$. The overlay of six of these beams onto the drive target is, therefore, expected to give an overall nonuniformity of 5%.

Acceleration Measurements

The measured acceleration for a $20\text{-}\mu\text{m}$ -thick CH foil driven with the 3-ns ramp pulse with a peak intensity of $2.5 \times 10^{14}\text{ W/cm}^2$ is shown in Fig. 71.8. The data are shown as squares with error bars that reflect the accuracy of the position measurement. The round points are the position of the foil as simulated by the 1-D hydrodynamics code *LILAC*. The agreement between the measured foil position and the *LILAC* simulation is excellent and shows that the target drive is being accurately modeled.

While the data shown in Fig. 71.8 are consistent with a constant acceleration, the *LILAC* simulation shows that the acceleration (Fig. 71.9) for a 3-ns ramp pulse is not constant and changes by a factor of 2 during the experimental measurement. This points to the importance of complete simulations rather than the application of simplified models that generally assume constant acceleration.

Figure 71.8

Comparison of the measured acceleration of a $20\text{-}\mu\text{m}$ -thick CH foil with a *LILAC* 1-D hydrodynamic simulation. The data are shown as squares with error bars and the calculation as circles.

E8453

Figure 71.9

The calculated acceleration of a 20- μm -thick CH foil illuminated with a 3-ns ramp pulse with a peak intensity of $2.5 \times 10^{14} \text{ W/cm}^2$. The acceleration predicted by the simulation is not constant in time but continues to increase during the target illumination.

RT Growth Measurements

Face-on radiography was used to measure the RT growth of imposed perturbations. The modulation in the optical depth is calculated by analyzing a $250\text{-}\mu\text{m} \times 250\text{-}\mu\text{m}$ region of the target for each temporal frame. The backlighter spatial distribution is taken out by fitting a two-dimensional, fourth-order polynomial to the natural log of the intensity. The fourth-order polynomial fit was used to model the x-ray intensity distribution from the backlighter for shots where no drive foil was present. It has also been verified by using higher-order polynomial fits (up to eighth order) and showing no significant improvement in the quality of the polynomial fit as determined from the value of χ^2 . This fit is then subtracted from the data, leaving the difference. This analysis tacitly assumes that the low-order polynomial adequately models the backlighter emission profile and does not affect the perturbation measurement. The analyzed region for the 31- μm -wavelength perturbation is chosen to be exactly eight wavelengths across the frame. A 2-D Fourier transform (FFT) then yields the optical depth at the perturbation wavelength. One dimension of the FFT is used to measure the imposed perturbation temporal evolution, and the orthogonal direction is used to determine the residual noise in the measurement.

Data from three target shots (plotted in Fig. 71.10) were taken with the 3-ns ramp pulse with a peak intensity of $2.5 \times 10^{14} \text{ W/cm}^2$. The measured data are shown as either diamonds, squares, or triangles. The error bars are calculated by analyzing separate regions in the x-ray images. The results are averaged and the standard deviation is plotted as error bars. The framing camera was triggered at three different times to study the evolution of the perturbation over 2.5 ns. The thick solid line in Fig. 71.10 is the optical depth calculated by the 2-D hydro-

namics code *ORCHID*. It agrees with the measured optical depth late in time but not at early times. The thin solid line is an estimate of the noise in the experimental measurement as determined using the FFT technique outlined above. We thus estimate a minimum measurable optical depth of 0.033 as a noise source. The dashed line in Fig. 71.10 is the sum of the optical depth as calculated by *ORCHID*, and the noise and results are a good fit over the entire measurement time.

E8402

Figure 71.10

A comparison of the experimentally measured optical depth and the predictions of the *ORCHID* 2-D hydrodynamic simulations. Data are shown for three shots as diamonds, squares, and triangles; the calculated optical depth is shown as a solid line. Good agreement with the experimental data is achieved only when the experimental noise is added to the calculated optical depth.

The 3-ns square pulse data are shown in Figs. 71.11(a) and 71.11(b). The growth of an initial perturbation with a wavelength of 60 μm and an initial amplitude of 0.5 μm is shown in Fig. 71.11(a), and the growth of a 31- μm -wavelength perturbation with an initial amplitude of 0.45 μm is shown in Fig. 71.11(b). The data are plotted as points, and the *ORCHID* simulation is plotted as a solid line. In these data the experimental noise has been subtracted from the data before being plotted with the simulation output. Several shots are plotted together so that the temporal development of the modulation in the optical depth can be shown. The agreement between the variation in the optical depth measured with x-ray radiography and the *ORCHID* simulations for both cases is quite good over the temporal range of interest to determining the linear RT growth rate.

The *ORCHID* calculations are a complete end-to-end simulation of the experiment and include the backlighter spectrum,

Figure 71.11

A comparison of the experimentally measured optical depth and the predictions of the *ORCHID* 2-D hydrodynamic simulations for both the 60- μm and 31- μm perturbations. Data are shown as diamonds; the calculated optical depth is shown as a solid line. Good agreement with the experimental data is achieved only when the experimental noise is subtracted from the measured optical depth. These are data for a 3-ns square drive pulse.

all filtering of the x-ray framing camera, and a correction for the camera resolution. The comparison of the optical depth calculated from the simulation with the experimental data shows the quality of the simulation. A simple comparison of the growth rate as calculated by a simulation and the growth rate taken from a set of experimental data is not sufficient. Detailed comparisons of the data from a series of planar-foil experiments and their associated numerical simulations show that associating the measured growth rate of the modulation in optical depth with that due to the ablative RT instability must be done with great care, if done at all. During this study, we have identified two important sources of departure in the growth rates inferred from variations in the optical depth as compared to determinations based on variations in the actual displacement at the ablation surface. The first source of discrepancy that must be accounted for is the "noise" in the measurement. Figure 71.10 indicates that using the measured variation in the optical depth yields a growth rate value of 1.6 ns^{-1} , while *ORCHID* (without noise) predicts approximately 2.1 ns^{-1} . Therefore, the accurate determination and treatment of experimental noise are important in determining growth rates. The second, and possibly more important, cause of problems in determining growth rates from the variations in the optical depth has been found to be the initial shock propagation through the foil. The *ORCHID* simulations of these experiments have shown that for the pulse shape-foil combinations used, the shock propagation through the foil has

a large influence on the temporal evolution of the modulation in optical depth early in time. The comparison within a given simulation of the growth rate from the evolution of the optical depth with the growth rate determined from the evolution of the ablation surface gives substantially different results. In all cases of large-initial-amplitude mass perturbations considered to date, we have found that the growth rate determined from the modulation in the optical depth is lower than that determined from the modulation in spatial position of the ablation surface (illustrated in Fig. 71.12). Analysis of the *ORCHID* results shows that the propagation of the initial shock through the foil results in a change in the $\rho\Delta x$ of the foil (predominantly through a change in the density) even before the ablation surface has accelerated (or moved) appreciably. A simple approximation to the temporal history of the variation in the $\rho\Delta x$ data has a functional form of $a = a_0 \cdot e^{\gamma t} + \text{constant}$, where the constant is due to the initial shock propagation.

To check this finding one can take the time derivative of $\rho\Delta x$. A closer analysis of the *ORCHID* simulation of a 31- μm -wavelength perturbation is plotted for the time of linear growth (0.8 ns to 1.1 ns) in Fig. 71.13. The data plotted as diamonds are the amplitude of the optical depth of the planar target as determined from the uranium backlighter. The amplitude of the $\rho\Delta x$ at the ablation interface is plotted as squares. The Δx amplitude from the variation in position at the ablation surface is plotted as x's. The growth rate as determined by the $\rho\Delta x$ and

Figure 71.12

ORCHID simulation results for the growth of the perturbation at the ablation surface and the modulation in the optical depth are shown for one of the 31- μm , 0.45- μm initial amplitude experiments.

the optical depth are in agreement, but there is a significant difference in the growth rates as determined by radiography (optical depth) and the growth rate as determined from the ablation surface. The data plotted as triangles are the time derivative of the amplitude from the $\rho\Delta x$. When the time derivative of $\rho\Delta x$ is taken, the value of γ as determined by the variation in the displacement of the ablation surface is recovered. This behavior was also found for the 60- μm , 0.5- μm -initial-amplitude case and for both pulse shapes considered. (This second reason for growth-rate departures is reminiscent of the experimental noise effect encountered during the analysis of the framing camera data.)

Summary

The OMEGA laser facility has been used to study the growth of perturbations due to the RT instability at the ablation interface. The targets were accelerated with 351-nm illumination at $2.5 \times 10^{14} \text{ W/cm}^2$ for 3-ns ramp pulses and at $2.0 \times 10^{14} \text{ W/cm}^2$ for 3-ns flat-topped pulses. The comparison of target acceleration with 1-D hydrodynamics simulation is good and suggests that the laser drive is correctly modeled. The comparison for the growth of initial perturbations as measured by the change in optical depth agrees with 2-D ORCHID simulation when the measurement noise is included. It is important to compare the values of the measured optical depths and not simply the growth rates in order to understand the underlying physics and infer RT growth rates.

Figure 71.13

The growth of the perturbation as a function of time is plotted for three different simulation quantities: optical depth, $\rho\Delta x$, and ablation-surface amplitude. Both the optical depth and $\rho\Delta x$ show growth rates that are substantially different from the ablation-surface amplitude. The time derivative of $\rho\Delta x$ shows a growth rate that agrees with the ablation-surface amplitude.

The pulse shapes used in our experiments produce both accelerations and ablation velocities that evolve in time. This implies that care must be taken when determining growth rates from experimental data to ensure the validity of the underlying assumptions. The growth rates can be used to compare experimental data for targets with similar perturbations when the temporal region studied coincides with regions where the acceleration and ablation velocity are relatively constant.

It is important that the numerical simulations include as many details associated with the actual experiment and diagnostic configuration as possible. The comparisons between the hydrodynamic simulations and the experimental data should be made with the temporal evolution in optical depth modulation. If good agreement is found between the experimental data and the simulations, then the ablation-surface growth rate as determined from the simulation can be used with some level of confidence, as the RT growth rate for the initial perturbation used in the experiment.

It is evident from the effect of the noise on the comparison with calculations that future experiments will need to address this issue. Instruments and experimental techniques are needed that will reduce the level of noise so smaller initial perturbation amplitudes can be studied. Instruments are also needed that will allow shorter-wavelength perturbations to be studied.

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Fabrication of Polyimide Shells for Use as ICF Targets

The current target-shell material for direct-drive experiments on OMEGA is a low-density hydrocarbon polymer fabricated using a low-pressure plasma polymerization process. This material meets the immediate requirements for an inertial confinement fusion (ICF) target shell: low density; smooth surfaces (outer and inner); spherical and concentric; and the capability of being fabricated with dopants at discrete radial positions in the shell wall. Future ICF experiments, especially those conducted at cryogenic temperatures, would benefit if the targets possessed additional properties: higher tensile strength; larger elastic moduli; greater room-temperature permeability; greater radiation resistance; higher thermal conductivity and lower electrical conductivity; and greater opacity at ~ 351 nm. These desirable properties will allow targets to be filled more rapidly with DT; be cooled faster to the DT triple point (19.8 K); survive higher temperature gradients in a cryostat; better resist the damaging effects of β -decay from tritium; and more readily accommodate augmented layering techniques such as RF-coupled joule heating. These properties will also allow decreasingly thin-walled shells to be used to contain the DT fuel; the goal is a 1-mm-diam target with a 1- μ m-thick wall.

Polyimide is the only polymeric material with the potential for meeting these additional material requirements. Scientists at Lawrence Livermore National Laboratory were the first to suggest polyimide¹ targets for the National Ignition Facility (NIF). If targets could be fabricated with the same strength as commercial polyimide films, then cryogenic targets could be filled and transported at room temperature and frozen only immediately prior to the experiment. (The requirements for a NIF target are 2.2-mm outer diameter, 160- μ m walls, and a DT room-temperature fill pressure of ~ 360 atm.) These simplifications would substantially reduce the complexity and cost of implementing indirect-drive cryogenic targets on the NIF.

Postulating and theoretically validating polyimide shells as a viable ICF target design were necessary first steps. Fabricating polyimide shells to experimentally quantify the material properties is the next phase. No polyimide shells have been

made until now because the available polyimide fabrication techniques could not be adapted to current target-fabrication technology. Two recent developments, one in the area of ICF target-fabrication methods and the other in the area of polyimide fabrication for microelectronics applications, have provided a method for fabricating polyimide targets: (1) the development of a decomposable polymer mandrel for making the present ICF targets and (2) the development of a technique for vapor-depositing polyimide films. This study employs the combination of these two techniques to produce polyimide shells and presents property data of the resultant material.

Background

1. Preparation of Solution-Cast Polyimide Films

Polyimides are usually produced in a two-step synthesis:² First, an aromatic diamine (e.g., ODA) is dissolved in a polar solvent such as dimethylformamide. An aromatic dianhydride (e.g., PMDA) is then added to the solution to form the soluble polyamic acid. The polyamic acid solution is spun-cast onto flat substrates and then thermally cured at 300°C to form the polyimide. The chemical structures are shown in Fig. 71.14. There is no straightforward way this technique can be adapted to make ICF targets using established microencapsulation methods. Furthermore, there are technical issues with this technique that may render the material unsuitable as a target shell material: polyamic acids are hydrolytically unstable; solvents may remain in the film; adverse interaction may occur between the solvent and the substrate (i.e., mandrel); and the viscosity of the polymer solution may make it difficult to control the thickness and uniformity of the shell wall.

2. Preparation of Vapor-Deposited Polyimide Films

The vapor-deposition-polymerization^{3,4} (VDP) technique of making polyimide film was developed to provide a solventless method of depositing conformal dielectric coatings. The process involved three steps: (1) dianhydride and diamine monomers were sublimed under high vacuum, (2) the vapor phase monomers reacted to form polyamic acid on the surface of the substrate, and (3) the polyamic acid was thermally cured at 300°C to form polyimide. The advantages of this pro-

Figure 71.14

Formation of polyimide from pyromellitic dianhydride (PMDA) and oxydianiline (ODA) precursors.

cessing technique are the lack of detrimental solvent effects (i.e., vacuoles) and the uniformity of the film thickness. This technique was used in this study to make OMEGA-size ICF targets by depositing the film on spherical depolymerizable mandrels.

3. Fabrication of Targets Using Depolymerizable Mandrels

The depolymerizable mandrel technique was recently developed to fabricate plastic ICF targets.⁵ The mandrels were either solid beads or hollow shells made from poly(α -methylstyrene) (PAMS) that were overcoated with plasma polymer using the traditional low-pressure glow discharge polymerization process.⁶ Heating the overcoated mandrel to 300°C decomposes the mandrel into its individual monomer units that permeate through the shell wall. The importance of the PAMS formulation is that it undergoes complete depolymerization into gaseous monomers at a relatively low temperature (160°C–300°C), where the plasma polymer outer shell is thermally stable.

Currently, the PAMS mandrels are thin-walled (~20- μ m) shells (0.70- to 0.95-mm diam) fabricated by the micro-encapsulation method.⁷ (This method was initially developed to

make polystyrene shells for ICF experiments but is better suited for making mandrels.) A critical requirement of the mandrel, which has been demonstrated,⁸ is a smooth, spherical surface.

4. Comparison of Target-Material Relevant Properties

Table 71.I compares the relevant properties of the current target-shell material against the properties of polyimides fabricated using both solution-cast and vapor-phase-deposition techniques.

There are many formulations of solution-cast polyimide films, and the material properties are well defined. The values listed here are the optimal values reported for two types of commercially available polyimide: Upilex-S and Kapton. The values listed for vapor-deposited polyimide films are from the literature. However, these values possess greater uncertainties and are likely to be more process dependent. The strength and modulus values listed for ICF shells are the best estimates derived from a limited set of burst and buckle tests of 1-mm-diam shells.

Fabrication and Characterization

1. Materials

Pyromellitic dianhydride (PMDA) (1,2,4,5-benzenetetracarboxylic dianhydride; 97% purity) and 4,4'-oxydianiline (ODA) (99+% purity; zone refined) were received from Aldrich Chemical and used without further purification. Poly(α -methylstyrene) mandrels were received from General Atomics (850- to 950- μ m inner diam; ~20- μ m wall thickness; MW 400,000) or fabricated in-house (700- to 950- μ m inner diam; ~10- μ m wall thickness; MW 97,500).

2. Deposition Equipment

The deposition system consists of a vacuum chamber, mechanical pump, and two identical evaporation sources, one each for PMDA and ODA (Fig. 71.15). The chamber was a stainless steel cube with an adapting tee (6-liter internal volume) and three glass viewports. The system was pumped with a mechanical pump (Alcatel model 2010) through two liquid nitrogen foreline traps and achieved a base pressure of 1 mTorr. Each evaporation source was a copper cylinder containing a 0.3-cm-outer-diam $\times$ 5-cm cartridge heater (Watlow Firerod; 50 W max) and a K-type thermocouple. The temperature was maintained within $\pm 1.0^\circ\text{C}$ of the set point using a proportional-integral-derivative controller (Omega CN9000).

Each evaporation source contained ~50 mg of monomer, which was sufficient material for an 8-h process cycle. The

Table 71.I: A comparison of available relevant material properties of existing ICF shells and two types of polyimide.

	Current ICF Plasma Polymer Shell	Vapor-Deposited Polyimide Film	Solution-Cast Polyimide Film
Ultimate tensile strength (MPa) at 300 K at cryogenic temperatures	75±15 ^a 160±12 at 40 K ^b	120–170 ^{c,d} –	420 ^e 500 at 77 K ^e
Elastic modulus (GPa) at 300 K at cryogenic temperatures	3.1 ^a 9.7 at 15 K [*]	2.3–4.0 ^{c,d} –	9.3 ^e 12 at 77 K ^e
Density (g/cm ³)	1.039	1.39 ^d	1.42 ^f
Coefficient of thermal conductivity (W/m·K)	0.04 at 20 K ^{**}	– –	0.82 at 293 K ^g 0.08 at 20 K ^g
Dielectric constant	2.5 (at 1 MHz) ^{**}	2.91 (at 1 MHz) ^c	3.4 (at 1 GHz) ^f
Radiation resistance (electron) (dose required to reduce the maximum elongation by 50%)	–	–	6 × 10 ⁷ Gy ^f (1 × 10 ¹⁴ Ci·s)
Permeability at 300 K (mol·m/m ² ·Pa·s)	6.5 × 10 ^{–15} (for D ₂) ^a	1 × 10 ^{–19} (for O ₂) [*]	2.9 × 10 ^{–15} (for H ₂) ^h

* data for a polyamide film.

** data for a polystyrene film.

^a R. Gram, UR/LLE (private communication).

^b N. Alexander, General Atomics (private communication).

^c Ref. 3.

^d M. Iijima, Y. Takahashi, High Perform. Polym. **5**, 229 (1993).

^e H. Yamaoka, K. Miyata, O. Yano, Cryogenics **35**, 787 (1995).

^f DuPont Technical Bulletin. Product information for Kapton Type-H film.

^g H. Yokoyama, Cryogenics **35**, 799 (1995).

^h Ref. 15.

optimal processing conditions that yielded a uniform distribution of each monomer in near-stoichiometric proportions are listed in Table 71.II. The partial vapor pressure of PMDA and ODA at these conditions was 0.35 and 0.15 mbar⁹ (0.26 and 0.11 Torr), respectively.

The preferred experimental configuration for making particulate-free coatings is to suspend the substrate upside down in the vacuum chamber with the evaporation sources directly beneath the substrate. This configuration cannot be used with the existing bounce-coating method for coating shells, so the evaporation sources were built to direct the monomer flow downward onto the substrate. The substrates were either flat witness plates (NaCl and silicon wafers) mounted on a thermally isolated support or spherical mandrels placed in a bounce pan. The bounce pan consisted of a 2-cm curved dish with a piezoelectric element attached to the base. A 5-V, 3-kHz signal applied to the piezo-material maintained the shells in a gentle, random motion during the deposition process. Unlike previous techniques used to coat shells (e.g., plasma parylene

Figure 71.15
Schematic of the deposition chamber: (1) vacuum chamber, (2) deposition source, (3) heater, and (4) bounce pan containing shells.

Table 71.II: Processing conditions used to deposit polyamic acid films from PMDA-ODA monomers.

PMDA temperature	160°C
ODA temperature	140°C
System pressure	1 to 30×10^{-3} Torr
Separation between evaporation sources and substrate	5 cm
Diameter of substrate that received a uniform coating	3 cm
Coating rate on shells	0.8–1.0 $\mu\text{m/h}$
Coating rate on flat substrates	$\sim 3 \mu\text{m/h}$

and plasma polymer), the VDP-polyimide technique does not require a plasma to initiate or accelerate the polymerization process. However, the plasma had an essential serendipitous benefit of preventing shells from sticking together by overcoming the static charge that developed as targets rubbed against each other. (Shells bouncing in the sheath of the RF plasma developed a similar potential that swamped the friction-induced static charge and caused the shells to mutually repel.)

The experimental equipment was modified to add a mechanism to overcome the induced static charge. First, a copper electrode was added inside the vacuum chamber to ignite an argon plasma (glow discharge). The pressure was increased from 1 to ~ 30 mTorr (using argon), and the plasma was maintained for the minimum period necessary to dislodge the shells from the pan (~ 15 s every 15 min). An alternative method of dissipating the static charge used a thermionic electron source directed at the pan. A tungsten filament was positioned ~ 1 cm from the pan, and a constant current (6 mA) was maintained through the filament. The resulting flood of electrons negatively biased the shells and prevented them from sticking together. This process was independent of pressure, which permitted the base pressure of 1 mTorr to be maintained during deposition. Both techniques successfully prevented shells from sticking to themselves or the pan during the >6 -h processes, and shells made by both techniques were evaluated in subsequent tests.

3. Heating Cycle

A programmable vacuum furnace (Ney Centurion VPM) was used to thermally convert (i.e., imidize) the polyamic acid shells to polyimide, and to decompose the inner PAMS mandrel. The furnace's cylindrical heated zone was 10-cm diam $\times$ 6.3-cm high, and the progress of the depolymerization process was monitored through a glass viewport. The heating cycle was performed in a static air environment. The heating rate was controlled with an accuracy of $\pm 0.1^\circ\text{C/min}$.

Each shell was placed in a separate glass tube in the furnace to protect it from fratricidal impact should any shell burst during curing. The optimum heating cycle (Fig. 71.16) required the temperature to be raised from 25°C to 140°C at 2.0°C/min , with a 2-h dwell, then raised to 150°C (5-h dwell), 180°C (12-h dwell), 220°C (12-h dwell), and 300°C (24-h dwell), at a rate of 0.1°C/min . After the final heating stage the furnace cooled to room temperature at $\sim 0.5^\circ\text{C/min}$.

Figure 71.16

Curing cycle used to imidize the polyamic acid and depolymerize the PAMS mandrel.

4. Chemical Analysis—FTIR Measurements

The monomer precursors were co-deposited onto NaCl substrates for Fourier transform infrared (FTIR) characterization. FTIR spectra were taken on the as-deposited, partially cured, and fully cured films, as shown in Fig. 71.17. The spectrum of the as-deposited film showed absorption peaks at ~ 1850 and 1780 cm^{-1} due to the carboxylic acid anhydride and peaks at ~ 1650 and 1550 cm^{-1} due to amide couplings $\{[\text{C}=\text{O}(\text{CONH})]$ and $(\text{C}-\text{NH})$, respectively}. These spectra indicated that the as-deposited film consisted of nonreacted

Figure 71.17

FTIR spectra of the as-deposited, partially cured, and fully cured PMDA-ODA/polyamic acid/polyimide film. The characteristic polyimide peaks appear as the material cures.

monomers and polyamic acid. After curing at 300°C for 60 min, the imide absorption peaks were dominant: 1780 cm⁻¹ (C=O asym. stretch), 1720 cm⁻¹ (C=O sym. stretch), 1380 cm⁻¹ (C-N stretch), and 725 cm⁻¹ (C=O bend), and the amide coupling and carboxylic acid anhydride absorption peaks had disappeared. These spectra agree with similar results published in the literature.^{3,10} [The formal name for this polymer is poly(N,N'-bis(phenoxyphenyl) pyromellitimide) and is commonly known as PMDA-ODA polyimide. A solution-cast version is sold under the trade name Kapton®.]

5. Shell Dimensions

The outside diameter of the shells was measured to an accuracy of $\pm 1 \mu\text{m}$ using a Nikon Optiphot optical microscope equipped with a HMOS Micro-Measure digital scaler. The wall thickness was determined from SEM micrographs of cross sections through the shell with an accuracy of $\pm 0.3 \mu\text{m}$. (The targets were too opaque to use interferometry to measure the wall thickness.)

The sphericity was determined by measuring the outer diameter along a single axis and then repeating the process for additional axes. The sphericity s of a shell was defined as $s = \left[1 - (D_{\text{max}} - D_{\text{avg}}) / D_{\text{avg}} \right] \times 100\%$, where D_{max} is the maximum diameter and D_{avg} is the average diameter from 12 measurements. The sphericity of the polyimide shells was $>99\%$, a value consistent with plasma polymer shells fabricated by the decomposable mandrel technique.

Imidizing the polyamic acid and depolymerizing the mandrel caused the resulting polyimide shell to be smaller and thinner than the original polyamic acid shell. The polyimide shell's outer diameter was 1% to 2% smaller than the original diameter. The polyimide wall was 20% to 30% thinner than the wall of the polyamic acid shell. The range of dimensions of all the polyimide shells is listed in Table 71.III.

Table 71.III: Typical dimensions of polyimide shells successfully fabricated in this study.

Inner diameter	700 to 950 μm
Wall thickness	4 to 13 μm
Sphericity	$> 99\%$
Shrinkage:	
wall	20%–30%
outer diameter	1%–2%
Density	1.21 to 1.40 g/cm ³

6. Elemental Composition—RBS Measurements

The elemental composition of the fully imidized polyimide film was measured using Rutherford backscattering (RBS) analysis. The RBS technique has the advantage of providing absolute quantitative data without calibrating against NIST-traceable standards. Also, the technique can be used to study a single shell (probe size $\sim 5\text{-}\mu\text{m}$ -diam circle) and can provide depth-resolved compositional data.

The data were acquired using the Dynamitron at the Nuclear Laboratory at the State University of New York at Albany. The recoil data were fitted using the RUMP program. A 2-MeV He beam was used to acquire the backscattered signal. The total charge from the beam was kept below 6 μC to prevent the beam from heating the film and changing the elemental composition.

The measured elemental composition of one of the polyimide films was 50 at.% carbon; 37 at.% hydrogen; 9 at.% oxygen; and 4 at.% nitrogen. The uncertainty was $\sim 20\%$. The theoretical composition of a PMDA-ODA polyimide is 56 at.% carbon; 26 at.% hydrogen; 13 at.% oxygen; and 5 at.% nitrogen.

7. Burst and Buckle Pressure Measurements

The equipment consisted of a pressure cell with opposing windows, a microscope, a CCD camera, and a pressure transducer (Omega PX621) with an accuracy of ± 1 psi. The shells were viewed through a microscope with transmitted light. The shell's image, the gas pressure in the cell, and the lapsed time were all recorded on videotape.

To burst (or buckle) a shell, the pressure in the chamber was lowered (or raised) at a faster rate than the rate the internal shell pressure changed due to permeation through the shell wall. Typical pressure-rate changes ranged from 0.01 atm/s to 3.5 atm/s for buckle tests and up to 50 atm/s for burst tests. After the shell failed, the video tape was reviewed frame-by-frame to determine the exact rupture pressure. If the permeation time constant of the shell was known, the pressure difference across the shell wall at the point of failure could be accurately calculated, and the strength measurement was corrected for this change.

8. Permeability and Time Constant Measurements

a. For shells. The permeability of a shell was measured¹¹ by permeating gas into the shell and then transferring it into a (pre-baked) chamber of known volume (2.4 cm^3). The pressure rise in the chamber was monitored as a function of time, and the data were fitted with the expected exponential function. A correction term was added to this function to account for outgassing from the chamber walls ($\sim 4 \times 10^{-11}$ Torr $\cdot\ell/\text{cm}^2\cdot\text{s}$). The equation gives the shell's permeation time constant (for the specific gas) and the total quantity of gas in the shell at the time the pressure measurement began. The technique allowed permeation time constants as short as 1 min and as long as 50 h to be measured with an accuracy greater than 95%.

The permeation time constant τ was determined from the plot of pressure rise versus time. The material permeability coefficient K_p was determined using the permeation time constant of a thin spherical container, $\tau^{-1} = (K_p A R_g T) / w V$, where A is the shell's mean surface area, R_g is the ideal gas constant, T is the absolute temperature, w is the shell wall thickness, and V is the shell volume. For a thin-walled spherical shell this simplifies to $K_p = w \langle r \rangle / 3 \tau R_g T$, where $\langle r \rangle$ is the average shell radius.

b. For flat films. The equipment and test method were based upon the ASTM procedure, which estimates the steady-state rate of transmission of a gas through plastics in the form of a film, sheeting, laminates, and plastic-coated papers or fabrics.¹² A flat, freestanding polyimide film was mounted in a gas transmission cell to form a sealed semibarrier between two chambers. One chamber contained the test gas at a specific high pressure, and the other chamber, initially evacuated, received the permeating gas. The rate of gas transmission through the polyimide film was determined from the measured pressure rise in the downstream chamber. The permeability of 1-cm-diam, 14.7- μm -thick polyimide films was then measured using deuterium and nitrogen. Equipment accuracy was tested by comparing the measured permeability of commercial polyimide films (Kapton and Upilex-R) against the reported values; the agreement was within 50%.

9. Roughness Measurements

Surfaces of the shells and flat films along with cross sections of the shell walls were inspected using a scanning electron microscope (Cambridge Instruments, model S200). The micrographs provided a qualitative measurement of the surfaces' morphology and roughness. Quantitative measurements were obtained using atomic force microscopy (AFM): flat polyimide films on silicon witness plates were analyzed using a Nanoscope III instrument (Digital Instruments); and shells were analyzed using an AFM modified to measure 1-mm-diam shells. The latter technique, performed by General Atomics, measured the roughness along great arcs of a circle around a sphere. Generally, the flat films were very smooth and the spheres were rough (data discussed in detail later).

10. Mechanical Behavior—Stress Measurements

The residual stress that develops in the polymer during the imidizing cycle must be known so that the heating cycle can be optimized to minimize the induced stress. Shells that crack during curing behave this way because of the rate and magnitude at which the stress develops in the shell wall. In this

process, the change in the residual stress was expected to be significant; imidization involves an appreciable mass loss and a substantial change in the structure and morphology of the polymer as the polyamic acid is converted into an increasingly high-molecular-weight polyimide.

Thin polyamic acid films were deposited on high-aspect-ratio silicon (100) substrates (1 cm × 0.1 cm × 0.25 mm thick). The curvature of the substrate was measured prior to, and following, deposition using a profilometer (Rank Taylor Hobson Talysurf). The measured change in the radius of curvature (ΔR) due to the polymer film yielded the total stress (σ_{total}) in the film:¹³

$$\sigma_{\text{total}} = [E_s / 6(1 - \nu_s)] [t_{\text{substrate}}^2 / t_{\text{film}}] [\Delta R^{-1}],$$

where E_s is the elastic modulus of the silicon substrate, ν_s is the Poisson ratio, and t is the thickness of the film/substrate. The total stress was the sum of the intrinsic stress and the induced thermal stress (due to the thermal coefficient mismatch between the silicon substrate and the deposited film) that arises because the film was imidized at an elevated temperature and the stress was measured at room temperature. The thermal stress (σ_{thermal}) is given by

$$\sigma_{\text{thermal}} = [E_f / (1 - \nu_f)] [\alpha_{\text{silicon}} - \alpha_{\text{polyimide}}] \Delta T,$$

where E_f and ν_f are the elastic modulus and Poisson ratio of the polyimide film, respectively, α is the coefficient of thermal expansion for the silicon/polyimide, and ΔT is the difference between the cure and measurement temperatures. Subtracting the thermal stress determined the intrinsic component of the residual stress. These data are discussed later.

Discussion

1. Optimizing the Processing Conditions

The first task was to define the operating parameters required to deposit thin films of polyamic acid with a stoichiometric balance of PMDA and ODA monomers. Optical macrographs and FTIR spectra of the films deposited on flat NaCl witness substrates showed the importance of (1) positioning the substrate relative to the PMDA and ODA evaporation sources and (2) the temperature of the evaporation sources. The film deposited at the center of the substrate experienced an equivalent flux of each monomer and was opaque. FTIR spectra of this region identified polyamic acid as the primary

component. The perimeter of the substrate possessed a film that was thinner and more transparent than the center and contained excess PMDA or ODA, depending upon the substrate's position relative to the evaporation sources. Rotating the substrate during deposition improved the uniformity of the film's thickness and composition. These experiments defined the optimal distance between the evaporation sources and the substrate, the required temperature of the evaporation sources, the maximum diameter of the substrate that was uniformly coated with polyamic acid, and the resultant deposition rate (see Table 71.II). Examples of the FTIR data for this analysis are shown in Fig. 71.18 and compared to spectra reported in the literature.³ Here, the relation between the film's composition and the evaporation source temperature was demonstrated. When the source temperatures for PMDA and ODA were 160°C and 120°C, respectively, excess PMDA was present in the film, and the characteristic dianhydride peaks at 1780, 925, 905, and 720 cm⁻¹ were proportionally greater than the peaks attributed to ODA. Raising the ODA temperature to

Figure 71.18

FTIR spectra of the as-deposited films made with varying ODA deposition temperatures: (a) $T_{\text{ODA}} = 122^\circ\text{C}$, $T_{\text{PMDA}} = 160^\circ\text{C}$; (b) $T_{\text{ODA}} = 140^\circ\text{C}$, $T_{\text{PMDA}} = 160^\circ\text{C}$. Spectrum (b) is identical to the polyamic acid spectrum as reported in the literature. Spectrum (a) shows anhydride peaks due to excess PMDA.

130°C provided an observed improvement, but the best results (based upon comparisons with the literature) were obtained with the source temperatures of 160°C and 140°C for PMDA and ODA, respectively. The spectrum shown in Fig. 71.18(b) is identical to that reported in the literature as a 1:1 molar ratio polyamic acid film.³

The next phase was to define the thermal cycle required to convert the polyamic acid into the polyimide. Typically, VDP polyamic acid films were heated to 175°C for 30 min, then to 300°C for 30 min. The cycle used for the current films included a slower ramp and was interrupted at 100, 140, 180, and 220°C to monitor the changing chemical and mechanical composition of the film. FTIR spectra of the polyamic-acid-coated NaCl substrate were acquired, and the change in curvature of the coated Si substrate was measured.

Conversion of polyamic acid into polyimide was evident at the first inspection at 100°C. Two characteristic absorption peaks were present: the C-N stretch frequency at 1380 cm^{-1} and the C=O asymmetric stretch at 1720 cm^{-1} (Fig. 71.17). This observation is consistent with reports that state VDP polyamic acid films begin imidization at lower temperatures than do solution-cast films.¹⁴ As the heating cycle progressed, the imide peaks became more prominent and the polyamic acid absorption peaks (at 1650 to 1550 cm^{-1}) diminished. Anhydride peaks (1850 cm^{-1}) were still visible after heating to 220°C but disappeared after the final curing stage at 300°C. The changes in the stress that accompanied these chemical changes are described later.

The heating cycle used above was inadequate for imidizing polyamic-acid-coated PAMS shells; the simultaneous depolymerization of the PAMS mandrel overpressurized the shells, causing them to burst. A considerably slower heating ramp (see Fig. 71.16), with longer dwells at critical temperatures,

was required to balance the rate at which PAMS decomposed with the rate at which α -methylstyrene permeated through the polyamic acid/polyimide/PAMS shell wall. This heating cycle was determined empirically since the mechanical and permeability properties of the wall could not be predicted. Several events complicated the process because they occurred simultaneously: excess monomer was subliming; PAMS was depolymerizing; the polyamic acid was reacting to form the polyimide; the softening PAMS was slumping, and the weight of the viscous material was adding a nonuniform strain to the shell wall; and the shell wall was resisting an increasing burst-pressure as the α -methylstyrene partial pressure increased with increasing temperature. Identifying a successful heating cycle is the most important component of making thin-walled, high-aspect-ratio polyimide shells. Currently the thinnest shells successfully fabricated have 4- μm walls and 700- to 950- μm inner diameters.

2. Morphology of the Polyimide Material

Polyimide films (3- to 15- μm thickness) deposited on flat substrates were evaluated using optical, electron, and atomic-force microscopy and were observed to be very smooth. The rms roughness was 0.81 nm with most of the power attributed to features with a wavelength between 1 and 0.1 μm (Fig. 71.19). These results indicate that the deposition and imidizing process provided a material with a smooth surface that is acceptable for use in ICF experiments.

In contrast to the flat film, the shells possess substantial roughness. The rms roughness of a polyamic-acid shell deposited using plasma assistance was ~ 800 nm. This quantified roughness was consistent with SEM micrographs that showed a nodular surface with a high density of nodes 1 to 5 μm in diameter [Fig. 71.20(a)]. This roughness may be attributed to either the special processing conditions required to coat shells (the plasma- or thermionic-assisted technique) or a prob-

Figure 71.19
Power spectrum of a flat polyimide film. The rms roughness is 0.81 nm.

lem specific to forming a rigid polyimide structure in a spherical geometry.

One reason that shells possess a greater roughness than flat films is the high background pressure (20 to 30 mTorr) required by the plasma-assisted technique. A high pressure causes reacting monomers to experience thermalizing gas phase collisions before they reach the shell's surface. This reduces the mobility of the adsorbed atoms (adatoms), which prevents the nascent layer from fully relaxing before the subsequent atomic layer arrives at the surface and freezes the current distribution of adatoms in place. This loss of surface diffusivity generates nodular features with sizes inversely proportional to the adatom diffusion distance. These effects on the magnitude and frequency of the nodes are compounded by statistical fluctuations in the vapor flux, and the amplitude of the roughness grows with increasing film thickness, due to self-shadowing effects.

Figure 71.20

Outer surface (a) and cross section through the wall (b) of a polyimide shell made in the presence of a plasma. The nodular structure is seen through the shell wall.

Polyimide shells made by the plasma-assisted method were rougher and more nodular than the shells made by the thermionic emission method. Examination of the cross sections of fractured shells deposited with plasma assistance showed that the nodular morphology and internodular porosity extend through the wall to the inner surface (Fig. 71.20), while shells deposited with thermionic assistance possessed denser walls with no visible porosity and a smooth inner surface (Fig. 71.21). Despite the improved smoothness achieved at lower processing pressures, the smoothness of the shells deposited using thermionic assistance did not qualitatively approach the smooth-

Figure 71.21

Outer surface (a), cross section through the wall (b), and inner surface (c) of a polyimide shell made in the presence of thermionic emission. The wall appears to be fully dense. The indentations on the inner surface are imprints from imperfections and debris on the PAMS mandrel.

ness of the flat films. (A quantitative comparison remains to be performed.) These observations suggest that the increased pressure contributes to, but is not solely responsible for, the shell's roughness. The only remaining differences between the processing conditions used to make flat films and spheres are the bouncing motion, spherical geometry, and substrate material of the shells, from which we conclude that one or all of these effects contribute to the observed roughness.

The inner surface of the shells deposited by thermionic emission appeared similar to the surface of plasma polymer shells currently made using the PAMS mandrel. Shallow domes (~ 1 - to $5\text{-}\mu\text{m}$ diameter and $<10\text{ }\mu\text{m}$ deep) were sparsely distributed over the inner surface [Fig. 71.21(c)]. These features are replicates of the underlying mandrel imprinted into the overcoated polyimide as a consequence of the highly conformal mapping capability of the deposition process.

3. Material Property: Permeability

The rate at which hydrogen isotopes permeate through the polyimide material is a major factor that will determine the viability of vapor-deposited polyimide films as target shell material. A high room-temperature permeation rate will allow targets to be rapidly filled to the designated pressure (maximum is 1500 atm) for cryogenic ICF experiments. A slow room-temperature permeation rate will require that the targets be filled at elevated temperatures (200°C max); otherwise the material cannot be used. It is necessary that the material be impermeable below 26 K and desirable, although unlikely, that the material be permeable above ~ 50 K. This latter capability will prevent a pressure differential from developing across the shell wall during the cooling cycle (due to temperature gradients in the permeation vessel).

The permeation time constant τ of deuterium for a PAMS shell at 298 K was $\tau = 1.2$ to 3.9 min (τ depends on the shell dimensions). The permeation time constant for the PAMS shells with a polyamic acid coating was $\tau = 1.12$ h. (The polyimide shells discussed here were made using the thermionic emission technique.) After the polyamic acid was imidized and the mandrel removed, the permeation time constant could not be measured. There was insufficient gas in the target to obtain a pressure reading on the apparatus; the time constant was either less than ~ 45 s or greater than ~ 50 h. The results were similar when nitrogen—a less permeable gas—was used (a time constant could not be measured with the current equipment). Unfortunately, literature values could not clarify whether the time constant was short or long: permeation rates obtained from the literature for solution-cast polyimide films

with a similar composition indicate a short time constant for hydrogen,¹⁵ ~ 2.4 min, while the limited data available for vapor-deposited films suggest a very long time constant.¹⁶

The permeability of the shell material and underlying mandrel was measured after curing to increasingly higher temperatures of the imidizing cycle to determine at which point in the cycle the shell became more, or less, permeable. (The results are listed in Table 71.IV.) (These data are also a measure of the changing composition and structure of the shell.) The permeability of the PAMS mandrel was high ($K_p = 4.7 \times 10^{-15}$ mol $\cdot\text{m}/\text{m}^2\cdot\text{Pa}\cdot\text{s}$), which yields $\tau = 1.2$ min. The permeability of the polyamic-acid-coated mandrel ($6.4\text{ }\mu\text{m}$ of polyamic acid) was considerably smaller ($K_p = 0.074 \times 10^{-15}$ mol $\cdot\text{m}/\text{m}^2\cdot\text{Pa}\cdot\text{s}$), which yields $\tau = 74$ min. Heating the shell to 100°C for 1.5 h decreased τ to 17 min. Further heating to 120, 140, and 150°C lowered τ to 8 to 10 min. After the 165°C cure stage, the time constant could not be measured. These data indicated the shell was becoming more permeable as imidization progressed, but more positive evidence was necessary to determine whether or not the fully imidized shell was permeable.

Table 71.IV: Permeation data for polyamic acid/polyimide shells (D_2 at 298 K).

Curing Condition	τ (min)	$K_p \times 10^{17}$ (mol $\cdot\text{m}/\text{m}^2\cdot\text{Pa}\cdot\text{s}$)
PAMS mandrel	1.18	470
Polyamic acid + PAMS mandrel: as-deposited	74.4	7.4
After 100°C at 1.5 h	16.9	32
After 120°C at 1.5 h	8.54	64
After 140°C at 3 h	10.8	51
After 150°C at 5 h	8.52	64

Since critical permeation data were unobtainable for polyimide shells, a more sensitive technique using flat, vapor-deposited polyimide films was tested to obtain the values. (The results are listed in Table 71.V.) The deuterium permeability was 2.0×10^{-14} mol $\cdot\text{m}/\text{m}^2\cdot\text{Pa}\cdot\text{s}$ at 298 K, and the nitrogen permeability was 8.0×10^{-15} mol $\cdot\text{m}/\text{m}^2\cdot\text{Pa}\cdot\text{s}$ at 298 K. While we cannot exclude the possibility that the rapid permeation may be due to interconnected pores or pinholes in the film, no pinholes have been observed (using both AFM and SEM scans) and the size and density of any porosity are unknown. These permeability values were used to calculate time constants for typical OMEGA-size shells (Table 71.V). The time constant

Table 71.V: Measured permeability of different shell materials used to calculate time constants for an OMEGA-size shell.

	Permeability at 25°C (mol•m/m ² •Pa•s)		Time constant for 1-mm-outer-diam, 5-μm-wall shell	
	D ₂	N ₂	D ₂	N ₂
PAMS	4.7 × 10 ⁻¹⁵	1.6 × 10 ⁻¹⁶	1.2 min	34 min
Polyamic acid/PAMS	8.4 × 10 ⁻¹⁷	—	66 min	—
VDP polyimide	2.0 × 10 ^{-14*}	8.5 × 10 ^{-15*}	17 s	39 s

*Data obtained using flat films.

for a 1-mm-outer-diam, 5-μm-wall shell would be 17 s for deuterium and 39 s for nitrogen. These very short values may explain why we were unable to measure time constants of actual shells.

4. Mechanical Properties and Behavior of Polyimide Shells

The elastic modulus and tensile strength of the polyimide shells must be known to define the necessary conditions to deliver DT-filled cryogenic targets: (1) the greater the elastic modulus, the greater the buckling pressure the target can withstand, and (2) the greater the tensile strength, the greater the burst pressure the target can endure. When combined with the permeation rate, these values determine the maximum rate at which the target can be filled and cooled for ICF cryogenic experiments.

Shells with a diameter between 800 and 900 μm and a wall thickness between 5 and 9 μm were repeatedly pressurized and depressurized using the protocol and equipment described previously. The shell's buckle- and burst-pressure differential was measured with an accuracy of ±5 psi (0.34 atm). The initial pressurization rate was 0.01 atm/s and the maximum pressure was 45 atm. The depressurizing rate was 50 atm/s. (The minimum pressurization rate was based on initial estimations for permeability and elastic modulus of polyimide.) All of the shells survived every test. The pressurizing rate was increased to 3.5 atm/s, the maximum that could be achieved, to minimize permeation that would decrease the pressure differential across the shell wall. Shells deposited using the plasma-assisted technique buckled at an overpressure of ~36 atm. The failure mechanism was atypical of polystyrene and plasma polymers: an indent developed in the shell and the shell did not catastrophically disintegrate (Fig. 71.22). These shells exhibit a greater fracture toughness than do existing ICF targets. The shells deposited by the thermionic-assisted technique were more durable than those deposited by the plasma-assisted technique and did not buckle or burst.

The maximum pressure rating of the test apparatus was increased to 105 atm; shells deposited using thermionic assistance were repeatedly tested at incrementally higher pressures. The 8-μm-wall polyimide shell survived the maximum pressurization rate of 3.5 atm/s to 105 atm. Using the calculated permeation rate for nitrogen through polyimide (measured above), the maximum pressure differential across the shell wall (the buckle pressure) was ~60 atm, and, using the standard expression for the crush pressure,

$$P_{\text{buckle}} = \frac{2E}{\sqrt{3(1-\nu^2)}} \left(\frac{w}{r} \right)^2,$$

where E is the Young's modulus, ν is the Poisson ratio, w is the wall thickness, and r is the shell radius;¹⁷ the modulus of the material was calculated to be ~15 GPa. For comparison, the

Figure 71.22 Plasma-assisted polyimide shell buckled at an external overpressure of ~36 atm. The shell did not brittle fracture, indicating a high fracture toughness.

largest reported literature value for the modulus of a solution-cast polyimide is 9.3 GPa. (A 1-mm-diam, 10- μ m-thick shell would crush with an external overpressure of 44 atm.)

All attempts to burst the shells deposited by the thermionic-assisted technique were unsuccessful. Shells were gradually filled with nitrogen at a rate of ~ 4 atm/min to a final pressure of 110 atm and rapidly depressurized at the maximum rate possible (~ 50 atm/s). Assuming that negligible permeation occurred during the 2-s depressurizing cycle, the burst pressure was greater than 100 atm. Using this pressure value and the standard hoop stress equation, $P_{\text{burst}} = 2T(w/r)$, the tensile strength T was calculated to be greater than 300 MPa. For reference, the highest reported tensile strength for a solution-cast polyimide is 420 MPa.

The same polyimide shell was sputter coated with a 1- μ m-thick aluminum layer to provide a permeation barrier layer, and the identical pressurization test was repeated. The shell neither buckled during pressurization nor burst during rapid depressurization. The calculated strength and modulus of polyimide shells are compiled in Table 71.VI.

5. Residual Stress Induced in Polyimide Films

As discussed previously, the imidizing cycle is the critical processing phase that affects the survivability and the resulting properties of the shells. The chemical changes that occur during the curing cycle have been described, and the effect these changes have on the mechanical properties of the film are now discussed.

As the temperature increases through the imidizing cycle, polyamic acid reacts (loses water) to form a polyimide. Unreacted monomers present in the film are sublimed, and the molecular weight and rigid structure of the polyimide increase. These mass, morphology, and structural changes alter the intrinsic stress in the polymer. If the polyamic acid was an unconstrained freestanding shell, the induced stress would be reduced by the shell shrinking or expanding to relieve the induced radial and hoop stresses. However, that stress-relief

mechanism does not occur at lower temperatures where the PAMS mandrel possesses sufficient strength to resist these dimensional changes.

Below 110°C, the mandrel possesses significant material strength that constrains the polyamic acid. Above 110°C, the mandrel softens and the polyamic acid/polyimide can increasingly shrink or expand to relax the induced stress. Between 110°C and 180°C the α -methylstyrene vapor pressure increases within the shell and, together with the slumping PAMS mandrel, produces a nonuniform burst pressure on the polyamic acid/polyimide wall. (All the shells heated at rates greater than 1°C/min through the 110°C to 180°C regime exploded. Depolymerization of a typical PAMS mandrel would produce an estimated burst pressure of 45 atm at 150°C in the absence of permeation.) Above 180°C the polyamic acid continues to imidize and the PAMS mandrel has a decreasing influence on the process. All these effects contribute at varying degrees to the instantaneous stress in the shell wall.

The magnitude of stress in the polyamic acid/polyimide composite film at the important processing temperatures (100, 140, 165, 220, and 300°C) was measured using a 14- μ m polyamic acid film deposited on a silicon witness beam, as shown in Fig. 71.23. These data are corrected for the extrinsic thermal stress. (Flat films were used because no suitable technique was available to measure the stress that evolves in a shell.) Note: The data measured only the stresses that develop in polyimide during a specific cure cycle; additional stresses will exist in the polyimide shell because of the spherical geometry and the presence of the mandrel.

The intrinsic stress in the as-deposited film was ~ 15 MPa compressive. Heating the film to 100°C for 3 h induced a tensile stress. The maximum induced stress was 17 MPa at 165°C, where, coincidentally, the induced tensile stress offset the initial compressive stress. The tensile nature of the stress was expected as the film lost mass but was constrained at the interface by the substrate.

Table 71.VI: Calculated strength and modulus of polyimide shells made by the two techniques.

	Elastic Modulus (GPa)	Ultimate Tensile Strength (MPa)
Plasma-assisted polyimide	5.5	95
Thermionic-assisted polyimide	~ 15 (minimum)	300 (minimum)

Figure 71.23
Temperature dependence of the intrinsic stress at 25°C in a polyamic acid/polyimide film.

As the temperature increased to 300°C, the compressive stress in the film returned to 14 MPa. The origin of this stress may be explained by steric considerations: as the molecular weight of the polyimide increased, the polymer chains became more rigid and the scale length of the polyimide increased. This repelling interference of the polyimide chains was constrained by the chains being intertwined and cross-linked, resulting in a compressive force. This behavior would be expected to increase the porosity (nanoscale) of the material and may also explain why the imidized material is more permeable than the polyamic acid.

The geometry of the shell creates an additional stress not present in an unconstrained film: the inner surface will be under greater compressive strain than will the outer surface, and the resulting stress will be in addition to the intrinsic stress measured above. The stress difference will be small for thin-walled, OMEGA-size shells, and, since the magnitude of the intrinsic stress is expected to be small, the induced residual stresses in the shell wall are not significant. Knowing how the intrinsic stress changes with temperature will be used to optimize the imidizing cycle to obtain thinner-walled shells.

By contrast, NIF targets are thick-walled shells, and it is possible that a sizable stress gradient will exist across the wall. When the shell is filled to its specified pressure (currently 360 atm), the maximum tensile stress will be at the inner surface. And, if a sizable residual stress is also present, the maximum pressure rating of the target will be reduced. These considerations make it important to know the magnitude of the

induced stress. A desirable property of the polyimide shells observed here was the fracture toughness of the material; if the inner surface of a target experienced a localized highly strained region and a crack developed, the crack would help dissipate the stress, thereby avoiding catastrophic shell failure. This property may make transporting NIF cryogenic targets at room temperature more feasible.

Summary

A viable method for fabricating polyimide shells as potential targets for direct-drive ICF cryogenic experiments on OMEGA has been demonstrated. Shells with diameters ranging from 0.7 to 0.9 mm and wall thicknesses from 4 to 13 μm have been fabricated. Initial estimates of the elastic modulus, tensile strength, permeability, stress, composition, surface roughness, and burst/buckle pressures have been obtained. All of the effort discussed here has been devoted to obtaining polyimide shells and characterizing the resulting properties. No attempt has been made to optimize the processing conditions to improve the material properties. Importantly, no fundamental impediments were observed that would prevent polyimide shells being used as ICF targets. However, there are clearly many areas where the quality of the shell must be improved before the targets can be used for cryogenic experiments. That effort is the subject of future work.

The initially proposed mechanical-strength advantages of polyimide material have been experimentally substantiated: The shells are significantly stronger than existing polystyrene and plasma polymer shells. An initial concern about the low permeability of polyimide has been dispelled. In fact, the permeability is so high that future study is required to confirm that the permeability is an inherent material property and that the data are not distorted by unobserved interconnected porosity. Also, the impermeability of the shell at 20 K must be demonstrated.

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Sideward Stimulated Raman Scattering of a Short Laser Pulse in a Plasma Channel

Stimulated Raman scattering (SRS)¹ is the decay of an incident, or pump, light wave (0) into a frequency-downshifted, or Stokes, light wave (1) and an electron-plasma wave (2). The conservation of energy and momentum in this decay is signified by the frequency and wave-vector matching conditions

$$\omega_0 = \omega_1 + \omega_2, \quad \mathbf{k}_0 = \mathbf{k}_1 + \mathbf{k}_2, \quad (1)$$

in which $(\omega_0, \mathbf{k}_0)$ and $(\omega_1, \mathbf{k}_1)$ satisfy the dispersion relation $\omega = (\omega_e^2 + c^2 k^2)^{1/2}$, where ω_e is the electron-plasma frequency,¹ and $(\omega_2, \mathbf{k}_2)$ satisfies the dispersion relation $\omega = \omega_e$. The wave-vector matching condition is illustrated in Fig. 71.24(a).

Figure 71.24
Geometry of sideward SRS. (a) Wave-vector diagram for the upward (+) decay. The wave-vector diagram for the downward (-) decay is similar. (b) Region of the plasma illuminated by the laser pulse.

The SRS of a short laser pulse is important in the contexts of particle acceleration² and inertial confinement fusion.^{3,4} In previous studies of the spatiotemporal evolution of SRS,⁵⁻¹³ the Stokes waves were allowed to pass freely through the pulse boundaries. However, the radial ponderomotive force associated with the pulse can expel plasma from the neighborhood of the pulse axis, in which case SRS occurs in a plasma channel.¹⁴ A channel can also be preformed by a second pulse.¹⁵ The reflection of light by the channel walls couples Stokes waves propagating symmetrically relative to the pulse axis, as shown in Fig. 71.24(b), and enhances the growth of oblique SRS.

In this study we investigate the effects of reflections on sideward SRS in a plasma channel. Sideward SRS is a useful paradigm of oblique SRS because its two-dimensional geometry is representative of oblique SRS, but simple enough to preclude some of the mathematical complications associated with oblique SRS.¹³ To analyze SRS in an optical fiber, one decomposes the radial structure of the waves into the eigenmodes of the fiber.¹⁶ We avoid the difficulties associated with the determination of the (evolving) channel profile and the associated radial eigenmodes by using an empirical parameter r to characterize reflections from the channel walls.¹⁷

Mathematical Analysis

In the weak-coupling regime, the initial evolution of sideward SRS in an underdense plasma is governed by the linearized equations

$$(\partial_t \pm c \partial_y) A_{\pm} = \gamma_0 N_{\pm}, \quad \partial_t N_{\pm} = \gamma_0 A_{\pm}, \quad (2)$$

where $A_{\pm}$ represents the vector potentials associated with the Stokes waves, $N_{\pm}$ represents the density fluctuations associated with the electron-plasma waves, and $\gamma_0 = \omega_2 c |\mathbf{k}_2| A_0 / 4 (\omega_0 \omega_2)^{1/2}$ is the temporal growth rate of SRS in an infinite plasma.¹ It is advantageous to define the characteristic variables $\tau = (ct - x)/l$ and $\eta = y/l$, in terms of which Eqs. (2) become

$$(\partial_\tau \pm \partial_\eta) A_\pm = \gamma N_\pm, \quad \partial_\tau N_\pm = \gamma A_\pm, \quad (3)$$

where $\gamma = \gamma_0 l/c$.¹⁸ These equations are to be solved subject to the initial conditions

$$A_\pm(0, \eta) = 0, \quad N_\pm(0, \eta) = 1, \quad (4)$$

which are representative of a pulse convecting into fresh plasma, in which a constant level of density fluctuations is available to seed the instability, and the boundary conditions

$$A_+(\tau, 0) = r A_-(r, 0), \quad A_-(\tau, 1) = r A_+(\tau, 1). \quad (5)$$

It follows from Eqs. (3)–(5) that $A_-(\tau, \eta) = A_+(\tau, 1 - \eta)$ and $N_-(\tau, \eta) = N_+(\tau, 1 - \eta)$. Consequently, only the solutions for A_+ and N_+ will be stated explicitly.

One can solve Eqs. (3)–(5) by using Laplace transforms. For the special case in which $r = 0$ the wave amplitudes are given by

$$A_+(\tau, \eta) = \begin{cases} \sinh(\gamma\tau), \\ \sum_{n=0}^{\infty} F_{2n+1}(\tau, \eta), \end{cases} \quad N_+(\tau, \eta) = \begin{cases} \cosh(\gamma\tau), \\ \sum_{n=0}^{\infty} F_{2n}(\tau, \eta), \end{cases} \quad (6)$$

where

$$F_n(\tau, \eta) = [\eta/(\tau - \eta)]^{n/2} I_n \left\{ 2\gamma [\eta(\tau - \eta)]^{1/2} \right\} \quad (7)$$

and I_n is the modified Bessel function of order n . The first forms of Eqs. (6) are valid for $\tau \leq \eta$ and the second forms are valid for $\tau > \eta$. At any point (x, y) in the plasma, the pulse arrives and initiates the instability at $t = x/c$, after which the Stokes amplitude experiences exponential growth with growth rate γ_0 . The information that the side boundary at $y = 0$ is present reaches that point at $t = (x + y)/c$, after which the Stokes wave experiences Bessel growth. Since $I_n(z) \approx \exp(z)/(2\pi z)^{1/2}$ when $z \gg 1$, the Stokes amplitude grows in proportion to $\exp[2\gamma_0(yt/c)^{1/2}]$ when $t \gg (x + y)/c$. At these late times, $t \gg (yt/c)^{1/2}$. Thus, the sideward convection of the Stokes wave converts fast exponential growth to slow Bessel growth. A snapshot of the Stokes amplitude is displayed in Fig. 71.25(a) for the case in which $ct/l = 2$ and $\gamma_0 l/c = 2$.¹⁹ The transition from one-dimensional exponential growth to two-dimensional Bessel growth is evident.

For the special case in which $r = 1$, the wave amplitudes are given by

$$A_+ = \sinh(\gamma\tau), \quad N_+ = \cosh(\gamma\tau). \quad (8)$$

Thus, the total reflection of the Stokes light maintains the fast exponential growth of sideward SRS, even after the information that the pulse boundaries are present has reached the pulse interior. A snapshot of the Stokes amplitude is displayed in Fig. 71.25(b) for the case in which $ct/l = 2$ and $\gamma_0 l/c = 2$. It is evident that total reflection allows the wave evolution to remain one-dimensional and enhances the wave growth. What must be determined is how large a reflectivity is needed to have a significant effect.

For the general case in which $0 < r < 1$, the wave amplitudes satisfy the reflection principle

$$A_+(\tau, \eta|r) = \sum_{n=0}^{\infty} r^n A_+(\tau, \eta + n|0), \quad (9)$$

$$N_+(\tau, \eta|r) = \sum_{n=0}^{\infty} r^n N_+(\tau, \eta + n|0).$$

Figure 71.25

Contour plots of the logarithm of the Stokes intensity for cases in which $ct/l = 2$ and $\gamma_0 l/c = 2$. White corresponds to high intensity and black corresponds to low intensity. (a) $r = 0$; (b) $r = 1$.

For $0 \leq \eta \leq 1$ the fundamental solutions $A_+(\tau, \eta|0)$ and $N_+(\tau, \eta|0)$ are given by Eqs. (6). For $\eta > 1$ the fundamental solutions are zero until $\tau = \eta - 1$. Subsequently,

$$A_+(\tau, \eta|0) = \begin{cases} \sinh(\gamma\tau) - \sum_{n=0}^{\infty} F_{2n+1}(\tau, \eta-1), \\ \sum_{n=0}^{\infty} F_{2n+1}(\tau, \eta) - \sum_{n=0}^{\infty} F_{2n+1}(\tau, \eta-1), \end{cases} \quad (10)$$

$$N_+(\tau, \eta|0) = \begin{cases} \cosh(\gamma\tau) - \sum_{n=0}^{\infty} F_{2n}(\tau, \eta-1), \\ \sum_{n=0}^{\infty} F_{2n}(\tau, \eta) - \sum_{n=0}^{\infty} F_{2n}(\tau, \eta-1), \end{cases}$$

where the first forms are valid for $\eta - 1 < \tau \leq \eta$ and the second forms are valid for $\tau > \eta$. The modification of solutions (6) described by Eqs. (10) is associated with a pulse of infinite width and the initial condition $N_+(0, \eta) = H(1 - \eta)$, where H is the Heaviside step function. At early times, the wave amplitudes grow exponentially in time with growth rate γ_0 . At intermediate and late times the solutions involve doubly infinite sums of Bessel functions that represent the fundamental contributions to the wave amplitudes and the contributions caused by reflections at the pulse boundaries. A snapshot of the Stokes amplitude described by the exact Eqs. (6), (9), and (10) is displayed in Fig. 71.26(a) for the case in which $r = 0.5$, $ct/l = 2$, and $\gamma_0 l/c = 2$. By comparing Figs. 71.25(a) and 71.26(a), one can see that reflections do not affect the wave evolution in the region $y \geq ct - x$. Reflections affect the wave evolution significantly in the region $y < ct - x$. Despite the complexity of solutions (9), by analyzing the Laplace-transformed solutions from which they originated one can show that, at late times,

$$A_+(\tau, \eta|r) \approx \frac{\gamma(1-r)\exp(v\eta + \lambda\tau)}{v(2\lambda + v)}, \quad (11)$$

$$N_+(\tau, \eta|r) \approx \frac{\gamma^2(1-r)\exp(v\eta + \lambda\tau)}{\lambda v(2\lambda + v)},$$

where

$$\lambda = -v/2 + (\gamma^2 + v^2/4)^{1/2} \quad (12)$$

is the growth rate and

$$v = -\log r \quad (13)$$

characterizes the convective loss of Stokes light through the pulse boundaries. A snapshot of the Stokes amplitude de-

scribed by the approximate Eqs. (11)–(13) is displayed in Fig. 71.26(b). Although the approximate solutions (11) are inaccurate near the leading edge of the pulse, as one should expect, they are extremely accurate in the body of the pulse. The growth rate (12) is plotted as a function of r in Fig. 71.27 for the case in which $\gamma_0 l/c = 2$. For this case $r = 0.05$ is

Figure 71.26

Contour plots of the logarithm of the Stokes intensity for the case in which $ct/l = 2$, $\gamma_0 l/c = 2$, and $r = 0.5$. White corresponds to high intensity and black corresponds to low intensity. (a) Exact solution described by Eqs. (6), (9), and (10); (b) approximate solution described by Eqs. (11)–(13).

Figure 71.27

Normalized growth rate [Eq. (12)] plotted as a function of r for the case in which $\gamma_0 l/c = 2$.

sufficient to maintain exponential growth with $\lambda = 0.5\gamma$. In general, $r = \exp(-3\gamma/2)$ is sufficient to maintain this rate of exponential growth. Thus, *even weak reflections enhance the growth of sideward SRS significantly.*

Physical Model

Because of the significance of Eqs. (11)–(13), it is important to understand the physical principles on which they are based. The eigenmodes of a fiber are formed by internal reflection, and their existence is not predicated on the existence of SRS.¹⁶ These facts prompt the solution of Eqs. (3) and (5), with $\gamma = 0$. Subject to the initial conditions

$$A_{\pm}(0, \eta) = 1, \quad (14)$$

the solution of these equations is

$$A_{+}(\tau, \eta) = \sum_{n=0}^{\infty} r^n [H(\tau + 1 - \eta - n) - H(\tau - \eta - n)]. \quad (15)$$

Solution (15) is displayed in Fig. 71.28, together with the function

$$A_{+}(\tau, \eta) = [(1-r)/v] \exp(v\eta - v\tau), \quad (16)$$

which is the fundamental term in the residue series expansion of solution (15). Because of the tendency of the Stokes wave to propagate out of the plasma, the Stokes amplitude at the right boundary is reduced by a factor of $1/r = \exp(v)$ when $\tau = n$. The continuous function $\exp(-v\tau)$ captures this essential feature of the temporal wave evolution. When $\tau \neq n$, the right part of the Stokes wave has experienced one less reflection than the left part. Consequently, the Stokes amplitude at the right boundary is larger than the Stokes amplitude at the left boundary by a factor of $1/r$. The continuous function $\exp(v\eta)$ captures this essential feature of the spatial wave evolution. If one substitutes the ansatz $A_{+}(\tau, \eta) = a_{+}(\tau) \exp(v\eta)$ in Eqs. (3) and (5), with $\gamma = 0$, one finds that the wave amplitude decays exponentially in time, with decay rate v , in agreement with Eq. (16). Similarly, if one substitutes the ansatz $A_{+}(\tau, \eta) = a_{+}(\tau) \exp(v\eta)$ and $N_{+} = n_{+}(\tau) \exp(v\eta)$ in Eqs. (3) and (5), with $\gamma \neq 0$, one finds that the wave amplitudes grow exponentially in time, with growth rate $\lambda = -v/2 + (\gamma^2 + v^2/4)^{1/2}$, in agreement with Eq. (12). Thus, solutions (11) are consistent with the formation of a spatial eigenmode that grows in the presence of the pulse.

Figure 71.28

Evolution of the Stokes wave for the case in which $\gamma_0 l/c = 2$. The solid lines denote the exact solution (15) and the broken lines denote the approximate solution (16). (a) Stokes amplitude at $y/l = 1$ plotted as a function of ct/l ; (b) Stokes amplitude at $ct/l = 0.5$ plotted as a function of y/l .

Summary

In summary, we studied the spatiotemporal evolution of sideward SRS in the weak-coupling regime. In a uniform plasma the sideward convection of the Stokes wave inhibits the growth of sideward SRS. In a plasma channel the partial reflection of Stokes light by the channel walls allows a spatial eigenmode to form in the body of the laser pulse. After this eigenmode has formed, it grows exponentially in time for the subsequent duration of the pulse. The growth rate (12) depends directly on the reflectivity of the channel walls, which, in turn, depends directly on their height. Even weak reflections enhance the growth of sideward SRS significantly. We also made a preliminary analysis of sideward SRS in the strong-coupling regime. Although strong coupling changes the rate at which the instability grows, it does not change the qualitative features of instability growth described herein. In particular, the partial reflection of Stokes light by channel walls allows a spatial

eigenmode to form in the body of the pulse, and the ansatz described after Eq. (16) can be used to determine the rate at which it grows. For channels created by the pulse, the height of the channel walls depends directly on the pulse intensity. However, since pulses of high intensity evacuate the channel completely, sideward SRS is most important for pulses of moderate intensity. In future work we will apply the physical insights gained in this study of sideward SRS to the study of near-forward and near-backward SRS. Reflections are stronger for these instabilities than for sideward SRS because the Stokes waves approach the channel walls at grazing, rather than normal, incidence. A significant amount of near-forward SRS was observed in recent particle-in-cell simulations of pulse propagation in an underdense plasma.^{20,21}

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18. When the plasma corrections to the group speeds of the light waves are important, the first of Eqs. (2) becomes $(\partial_t \pm v_1 \partial_y) A_{\pm} = \gamma_0 N_{\pm}$, where v_1 is the group speed of the Stokes waves and $\gamma_0 = \omega_2 c |k_2| A_0 / 4 [(\omega_0 - \omega_2) \omega_2]^{1/2}$, and the characteristic variables become $\tau = (v_0 t - x)/l$ and $v_0 y / v_1 l$. Equations (3)–(5) are unchanged.
19. For sideward SRS in an underdense plasma, $\gamma_0 l / c \approx 2.2(n_e / n_c)^{1/4} (l / \lambda_0)(v_0 / c)$, where λ_0 is the laser wavelength and v_0 is the quiver speed of electrons oscillating in the laser field. One possible set of parameters for which $\gamma_0 l / c \approx 2$ is $n_e / n_c = 0.01$, $l / \lambda_0 = 10$, and $v_0 / c = 0.3$.
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Picosecond Photoresponse in Polycrystalline Silicon

The OMEGA laser driver line currently incorporates silicon (Si) photoconductive switches as part of the electro-optic pulse-shaping system. These switches conceptually serve two functions: (1) to generate an electrical square pulse with a sharp leading and trailing edge, when illuminated by a short optical pulse;¹ (2) to appropriately direct a broadband (0.1- to 10-GHz) shaped electrical pulse to the Mach-Zehnder integrated-optic modulator. The operating characteristics of the switches are optimized by minimizing the rise and fall times of the generated electrical square pulse and by maximizing the bandwidth of the directed broadband electrical shaped pulse.² Because the performance requirements of photoconductive switches in OMEGA pulse shaping are stringent and sometimes conflicting, a thorough exploration of the silicon photoconductive process is important.

In this article we address the generation of laser-pulse-induced electrical transients with short rise times, using silicon photoconductive switches. Related aspects of photoconductive switching, such as response at different laser wavelengths and Si integrated-circuit compatibility, are considered. We also explore the optimization of electrical pulse generation by probing the intrinsic limits of the Si and polycrystalline Si (polySi) material. In the 1980s Auston³ and others studied processing methods to reduce the Si intrinsic photoresponse time far below the microsecond limit set by its free-carrier lifetime. Although very fast responses were generated, the temporal resolution of the measurement equipment at that time was not sufficient to directly resolve the rise and fall times of the resulting photoresponse. To circumvent this measurement bandwidth limitation, optoelectronic correlation methods were developed.⁴ Correlation takes advantage of the switch itself by using another, similarly fast photoconductive switch in a pulser/sampler configuration. The measured response from the sampler is then a correlation of the two switch responses. This procedure, however, does not allow direct reconstruction of the switch response because, even if the switch geometries are identical, the measured response is not equivalent to a true autocorrelation since the electric fields across the gaps have different time evolutions. This is important because many

parameters of the photoconductive switches, such as carrier mobility and trapping, vary with the applied electric field strength. Therefore, detailed knowledge of the Si photoresponse and how it depends on the switch bias and illumination for different Si process conditions could not be determined.

We have measured the Si switch response at 800 nm and at the typical communication wavelength of 1.55 μm , using a 34-GHz sampling oscilloscope to quickly and easily measure properties dependent on charge (current integral) effects, such as quantum efficiency and saturation. We have also used an electro-optic (EO) sampling system capable of measuring submillivolt and subpicosecond responses.⁵ This allowed direct observation of the Si photoresponse rise and fall times.

Sample Preparation and Characterization

The tested samples were 2.3- μm -thick, 2×2 -mm, low-pressure chemical-vapor-deposited (LPCVD) films of polySi grown at a substrate temperature of 600°C in 200 mTorr of silane. A 1- μm -thick multilayer of metal was evaporated onto the polySi in an interdigitated fashion. The top layer was Au for solderability, while the bottom layer was Al for good adhesion to Si and to promote ohmic-like contacts. The samples were not annealed, and no implantation or etching was performed. Instead of growing the film on a silicon substrate with an oxide insulating layer, we used a fused silica substrate to facilitate switch illumination. This optimized the oscilloscope test fixture bandwidth because the sample could be directly soldered in a flip-chip manner across a microstrip gap, eliminating bandwidth-limiting discontinuities such as inductive wire-bond leads. Illuminating the sample from the back side (through the silica) also leads to reflection of the incident laser illumination from the top metal contacts, increasing rather than decreasing the number of photons absorbed, as front-side illumination would have done.

For characterization purposes, the polySi surface was etched preferentially at the grain boundaries, and a scanning electron microscope (SEM) image taken of the surface (see Fig. 71.29) revealed a polygrain size of approximately 30 nm. A sample

Figure 71.29

An SEM image of the polySi surface taken after preferential etching at grain boundaries reveals grain sizes of the order of 30 nm.

was also cleaved, and a 4-nm conductive silver layer was evaporated onto the exposed face, allowing a high-resolution SEM photograph of the switch cross section to be taken (see Fig. 71.30). The image shows the silica substrate beneath the 2.3-μm layer of polySi; on the top is a multilayer of metal that allows a robust, repeatable low-temperature-solder contact between the polySi and the microstrip transmission line. Close inspection of the grown polySi layer reveals some evidence of columnarity in the growth direction of the Si grains, which is not unexpected; however, the measurements were not significantly affected since the current flow was primarily in the direction normal to the silicon growth.

To independently confirm the polySi grain size and obtain preliminary data on absorption depth, a Perkin-Elmer Lambda 9 spectrophotometer was used to measure the transmission of the polySi-on-silica sample (see Fig. 71.31). The data, corrected for thin-film etalon fringes, show a broadened absorption band edge and a small but measurable amount of absorption in the energy bandgap, which is consistent with the small-grain microcrystalline film morphology observed in the SEM.

Figure 71.30

An SEM cross-section image showing the metal multilayer on the 2.3-μm polySi layer.

Figure 71.31

Transmission spectra of the polySi sample (solid line), corrected for thin-film etalon effects (dashed line).

Measurement and Experimental Results

A 34-GHz (10-ps intrinsic rise time) sampling oscilloscope measurement setup, shown in Fig. 71.32, allowed convenient measurement of relative quantum efficiencies for various laser wavelengths, since our EO sampling system operated up to only near-infrared wavelengths and did not extend into the fiber optic communication wavelengths. The oscilloscope measurements also permitted convenient testing of the switch signal's dependence on voltage bias (Fig. 71.33) and illumination power (Fig. 71.34). These two plots, taken at a laser wavelength of 800 nm, show that the switch response is linear

with both the voltage bias and incident laser fluence, indicating that the switch was not saturated. Additionally, the pulse-width dependence (Fig. 71.35) on bias voltage was nearly nonexistent, indicating that the photoresponse time was not due to carrier sweep-out but limited by either the free-carrier lifetime or the RC time constant of the switch. Since, as we will demonstrate later, EO sampling measurements show that the material response is in the 3-ps range, we can conclude that the observed photoresponse of our polySi switch is limited by the switch interdigitated geometry and the switch fixture RC time constant.

Figure 71.32
Sampling oscilloscope measurement setup. Inset shows the interdigitated switch geometry. Switch dimension is 2 mm × 2 mm; finger separation as well as finger width is 100 μm.

Z2196

Figure 71.33

The photoresponse at 800-nm illumination wavelength measured for different switch bias levels. The inset plots a linear dependence on gap voltage bias.

Z2195

Figure 71.34

The photoresponse at 800-nm illumination wavelength measured for different pulse average-power levels. The inset shows linear response versus illumination fluence.

Figure 71.35

The photoresponse FWHM at 800-nm wavelength and 40-nJ/cm² fluence versus the switch bias.

An optical parametric amplifier (OPA) pumped by an 810-nm regenerative amplifier and a white-light continuum system at 1-kHz repetition rate were used to measure the response at 1.55 μm . Since our switch was not processed to improve its subband absorption characteristics, the quantum efficiency of the device in the infrared was less than 10^{-4} . We were able, however, to observe a response at 1.55 μm (Fig. 71.36), which was again a switch-limited signal of 36-ps full-width at half-maximum (FWHM). This result indicates that with suitable efficiency improvement, a high-speed, monolithic, all-silicon, optoelectronic integrated circuit is feasible, using, for example, porous silicon⁶ as the transmitter and a type of polySi material for the detector.

Figure 71.36

The photoresponse at 1.55- μm illumination wavelength. The quantum efficiency was less than 10^{-4} .

Our switch configuration was not designed for EO measurements.⁵ Nevertheless, we managed to get preliminary results by connecting the switch as a meander-type slot line. The signal was generated and measured entirely on the face of the switch structure. It propagated in a meander line fashion along the gap from one end of the switch to the other. The transient measured before entering the first bend is shown in Fig. 71.37. We note a subpicosecond (system-limited) rise time, followed by a decay with an initial fall time of approximately 3 ps.

Figure 71.37

Electro-optically measured signal from a polySi switch connected as a meander-type slot line.

Conclusions

A standard fabrication technique was implemented to design small-grain polySi photoconductive switches fully compatible with Si VLSI processing. By means of EO sampling, a material photoresponse of 3-ps FWHM was observed. The switch responded to 800 nm and 1.55- μm femtosecond laser illumination with switch-geometry-limited photocurrent pulses shorter than 40-ps FWHM. The far-infrared response time was generated in nonimplanted silicon with ohmic-like metal-semiconductor contacts, indicating the response was limited by the relaxation time of the extended-state (free) carriers into localized (nonmobile) states. Our preliminary measurements of response time and efficiency can be optimized using simple process changes such as annealing at moderate temperatures, sputter-etching for surface damage, or Fe and Au deep-level defect doping, while still allowing the switch to be integrated with Si IC's.

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We believe that our measurements comparing flash-lamp to laser-diode pumping, as well as our thermal relaxation measurements, are the first measurements of this kind.

Experimental Setup

To measure the thermal distortions in our flash-lamp- and laser-diode-pumped Nd:YLF laser rods, the rods were separately placed in one arm of a Mach-Zehnder interferometer as shown in Fig. 71.38. The laser source for the interferometer was a Nd:YLF cw-mode-locked laser operating at a 1053-nm wavelength (this laser was chosen for convenience). The detection system (CCD camera) integrated over many milliseconds and, hence, could not resolve individual pulses in the mode-locked pulse train. The path lengths of the two interferometer arms were carefully adjusted to be approximately equal in order to produce high-contrast fringes. Wedge was introduced in the interferometer by adjusting the mirrors to provide approximately 40 fringes of tilt. To analyze the fringes, the rod was imaged onto a CCD camera with a lens placed outside the interferometer. The images of the high-contrast fringes from the camera were digitized with a computer equipped with a frame-grabber board. The camera was spectrally filtered with appropriate blocking filters to pass only the 1053-nm radiation of the interferometer laser source. The gain at the 1053-nm wavelength provided by the rod being tested did not significantly degrade the fringe contrast. Fluorescence from the pumped rod was effectively blocked with an aperture placed in the rod arm of the interferometer approximately 50 cm from the rod. The aperture diameter was adjusted to limit the angular field of view of the camera-collection optics while not introducing additional aberrations. The laser used for the interferometer was linearly polarized, and the rod was placed between two half-wave plates (HWP). The laser rods were fabricated with the crystalline a axis along the length of the rod and were placed in the interferometer with the a axis parallel to the interferometer optical axis. The crystalline c axis for

both the laser-diode and flash-lamp systems was therefore perpendicular to the interferometer optical axis and located in the plane of the rod end face. The interferometer laser polarization through the rod could be adjusted perpendicular to the c axis (σ) or parallel to the c axis (π) of the rod by adjusting the half-wave plate before the rod. The half-wave plate after the rod was used to realign the laser polarization state parallel to the interferometer reference arm for maximum fringe contrast on the camera. Fringes were captured by the computer and analyzed.

The digitized fringe patterns were transformed into wavefronts using a spatial-synchronous phase-detection technique.⁸⁻¹⁰ The high-frequency interferogram was captured by the frame grabber, stored in memory, and then Fourier analyzed to obtain the amplitude and phase of the two-dimensional wavefront exiting the laser rod. Fringe data were analyzed over only a portion of the rod's clear aperture (approximately 90% of the rod diameter) to obtain the wavefront distortions. To accomplish this, a circular software aperture centered on the rod was placed over the rod image. Only the fringe data within this aperture were used in the analysis. To calibrate the fraction of the rod aperture used in the analysis, a reticle was placed at the equivalent rod position in the opposite interferometer arm, imaged to the camera with the lens outside the interferometer, and overlapped in software with the image of the rod and the software aperture used for fringe analysis. A background wavefront was taken for each rod with no pumping (but with the simmer on for the flash-lamp data) and subtracted from all other wavefronts. The background-subtracted wavefronts of the pumped rods were then analyzed to determine the number of waves of defocus, astigmatism, spherical, and coma produced by pumping. The overall accuracy of this fringe analysis technique is believed to be $<\lambda/50$ (<0.02 waves). To calibrate this interferometric-measurement system, the focal length of a 1-m lens was measured with an accuracy of greater than 99%.

Figure 71.38

Experimental setup for measuring the pump-induced wavefront distortions in Nd:YLF laser rods.

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The thermal-time constant of our laser-diode-pumped Nd:YLF rod was also measured with the above equipment. The fringes were recorded on tape with a video camera recorder (VCR) as the rod pumping was turned off (and the liquid cooling left on). A frame sequencer was used to number the individual frames from the camera. By replaying the video tape through the computer, sequential data frames could be captured and analyzed. In this way, the aberrations could be measured at 60 frames (data points) per second.

The flash-lamp system had a 4-mm-diam $\times$ 54-mm-long Nd:YLF laser rod pumped with a single 4-mm bore $\times$ 24-mm-arc-length xenon flash lamp. The flash lamp was simmered with approximately a 400-V drop across the lamp and 50 mA of simmer current. The pulse-forming network for the lamp used a 150- μ f capacitor charged to 385 V (i.e., 11 J per pulse). The lamp's current pulse was approximately Gaussian with a pulse width of 250 μ s (FWHM). A small-signal gain of 1.8 at 1053 nm was measured in the Nd:YLF rod with these conditions (sufficient to operate the regenerative amplifier for our application). The system was operated to 60 pulses per second. Fringes within a 3.4-mm-diam software aperture centered on the rod were used to analyze this rod.

The laser-diode-pumped system, including the diode arrays, was designed and built by the McDonnell-Douglas Electronic Systems Company (no longer a part of McDonnell-Douglas) and is similar in performance (although operating at a different wavelength) to the one reported in Ref. 11. The system had a 5-mm-diam $\times$ 54-mm-long Nd:YLF laser rod pumped with three banks of laser diodes along the length of the rod. Each bank consisted of 16 diode bars arranged symmetrically around the circumference of the rod for a total of 48 diode bars arranged around the rod. The length of the rod pumped with this diode arrangement was 34 mm. The laser diodes had their polarization states parallel to the a axis of the crystal and hence perpendicular to the c axis. The diode emission wavelengths were centered around the 797-nm absorption peak of Nd:YLF. The absorption coefficient of Nd:YLF for this case is approximately 3 cm⁻¹. The pulse-forming network produced a 250- μ s-wide, 50-A square current pulse with approximately 100 V applied across the diode bars connected in series. With these conditions each diode bar emitted approximately 45 W of peak optical power (2.16-kW peak array optical power) for an optical-output-to-electrical-input diode-array efficiency of greater than 40%. A small-signal gain of 1.8 at 1053 nm was measured in the rod under these conditions. If we assume a stimulated emission cross section of 1.2×10^{-19} cm² for this

transition,¹² the stored energy in the rod was 0.18 J, resulting in an optical-pumping efficiency (stored energy to diode-array output-optical energy) of approximately 33%. This system was operated to 200 pulses per second. Fringes within a 4.4-mm-diam software aperture centered on the rod were used to analyze this rod.

Interpretation of Interferometric Data

When the Nd:YLF laser rod is unpumped, the wavefront distortion measured in transmission through the rod by the interferometer gives the amount of static aberration in the rod due to internal crystal strain. This wavefront is stored and is the background wavefront that is subtracted from all other measured wavefronts. The background-subtracted wavefronts measured with optical pumping, therefore, provide detailed information of the thermally induced refractive-index changes in the pumped laser rods. The distorted wavefronts exiting the thermally loaded laser rods are quantified in lowest order (using the third-order aberration theory¹³) by the linear superposition of a set of primary aberrations of defocus, astigmatism, coma, and spherical, which are defined below.

The thermally distorted wavefront exiting the rod can be expressed as a surface of constant phase:

$$\Phi_{\text{total}}(\rho, \theta) = \Phi_{\text{def}}(\rho) + \Phi_{\text{ast}}(\rho, \theta) + \Phi_{\text{coma}}(\rho, \theta) + \Phi_{\text{sph}}(\rho), \quad (1)$$

where the surface of constant phase for each of the primary aberrations is given by

$$\Phi_{\text{def}} = \lambda \Delta_{\text{def}} \rho^2, \quad (2a)$$

$$\Phi_{\text{ast}} = -\frac{\lambda \Delta_{\text{ast}}}{2} \rho^2 \cos(2\theta), \quad (2b)$$

$$\Phi_{\text{coma}} = -\lambda \Delta_{\text{coma}} (\rho^3 - \rho) \cos(\theta), \quad (2c)$$

$$\Phi_{\text{sph}} = \lambda \Delta_{\text{sph}} (\rho^4 - \rho^2), \quad (2d)$$

with $\rho = r/r_0$ the normalized radius r within the rod aperture, r_0 is the radius of the rod clear aperture used in the analysis ($r_0 = 1.7$ mm for the flash-lamp-pumped rod and $r_0 = 2.2$ mm for the laser-diode-pumped rod), θ is the polar angle in the rod

aperture measured from the c axis of the crystal, λ is the interferometer wavelength (1053 nm), and the Δ_i 's are the number of waves of the primary aberration measured (i.e., the weighting factors in the linear superposition).

We show experimentally that the dominant thermal aberrations are defocus and astigmatism (discussed in the next section). Both defocus and astigmatism have a quadratic dependence on the aperture coordinate r , and from these a focal length can be defined for two orthogonal planes. If we neglect the aberrations of coma and spherical, then we can define a focal length for the two orthogonal planes that contain the a axis (optical axis) of the crystal. (The plane containing the a and c crystal axes with $\theta = 0$ is referred to as $\parallel c$, and the plane containing the a and b crystal axes with $\theta = \pi/2$ is referred to as $\perp c$.) The wave fronts (surfaces of constant phase) associated with defocus and astigmatism are shown in Fig. 71.39. Pure defocus [Fig. 71.39(a)] is seen to be a spherical wavefront with equal curvatures in all directions (i.e., $0 \leq \theta < 2\pi$). Pure astigmatism [Fig. 71.39(b)] is seen to be a wavefront with curvatures of equal magnitude but opposite sign in two orthogonal directions (i.e., $\theta = 0$ and $\theta = \pi/2$). The thermal-focal length in a given plane is easily calculated from the number of waves of defocus and astigmatism and is given by (using a small angle approximation)

$$f_{\pm} = \frac{r_0^2}{2\lambda(\Delta_{\text{def}} \pm \Delta_{\text{ast}}/2)}, \quad (3)$$

where f_+ is the focal length in the plane perpendicular to the crystal c axis ($\theta = \pi/2$ or $\perp c$) and f_- is the focal length in the plane parallel to the crystal c axis ($\theta = 0$ or $\parallel c$).

Astigmatism and coma are odd aberrations and have an angle of symmetry associated with them. The wavefront analysis program calculates this symmetry angle for these aberrations, and our measurements show that these symmetry axes are aligned perpendicular and parallel to the c axis of the crystal as implied by Eqs. (2)—an expected result but previously unmeasured. The definitions and conventions discussed here were used to analyze our data.

Thermal-Distortion Measurements

We assume that the laser-diode- and flash-lamp-pumped rods are uniformly pumped so that heat is generated uniformly throughout the rod volume. Since the time between pump pulses is short (0.2 s for the lowest pulse-repetition rate used: 5 Hz) compared to the thermal-relaxation time of the rod (1.5 s; see Fig. 71.45 discussed later), pumping can be considered continuous and the temperature distribution in the rod will reach steady state after a brief warm-up period. The rods were liquid cooled, and in steady state the surface temperature of the rod will be constant for a particular pump power. A long cylindrical laser rod pumped with the above assumptions will produce a quadratic radial temperature and strain distribution that will lead to a quadratic radial index-of-refraction profile in the rod.¹² Since Nd:YLF is a uniaxial crystal, the quadratic radial profile will be different in two orthogonal directions.

Figure 71.40 shows the measured aberrations as a function of the pump repetition rate for the case of the laser-diode-pumped laser rod using π -polarized light in the rod (gain at $\lambda = 1047$ nm). Figure 71.41 shows the same measurements using σ -polarized light in the rod (gain at $\lambda = 1053$ nm). In these figures, the data points for defocus and astigmatism are fit to a straight line, while the data points for coma and spherical are

Figure 71.39

Surfaces of constant phase (wavefront shown with thin lines) for (a) pure defocus and (b) pure astigmatism as defined and used in our analysis. Here λ is the wavelength of the interferometer source (1053 nm), Δ_{def} is the number of waves of defocus measured, and Δ_{ast} is the number of waves of astigmatism measured.

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connected by dashed lines. It can be seen from these data that although a small amount of coma and spherical aberration was measured, defocus and astigmatism are the dominant aberrations. Coma and spherical will henceforth be neglected from our analysis. Thermal-focal lengths can be calculated from this data using Eq. (3) with $r_0 = 2.2$ mm (the radius of the rod

aperture used to analyze the fringe data). Figure 71.42 shows the thermally induced dioptric powers (defined as one divided by the thermal-focal length in meters) for the case of the laser-diode-pumped laser rod using π -polarized light in the rod (gain at $\lambda = 1047$ nm), and Fig. 71.43 shows the same data using σ -polarized light in the rod (gain at $\lambda = 1053$ nm). Again the data are fit to straight lines. In these figures, two dioptric powers are given corresponding to the planes parallel and perpendicular to the crystal c axis as described above.

E8487a

Figure 71.40

Waves of aberration measured with π -polarized light (corresponding to gain at 1047 nm) as a function of the pump repetition rate for our laser-diode-pumped Nd:YLF laser rod pumped to a small-signal gain of 1.8 at 1053 nm.

E8487b

Figure 71.41

Waves of aberration measured with σ -polarized light (corresponding to gain at 1053 nm) as a function of the pump repetition rate for our laser-diode-pumped Nd:YLF laser rod pumped to a small-signal gain of 1.8 at 1053 nm.

E8488a

Figure 71.42

Dioptric power (inverse of the thermal-focal length in meters) calculated from the data in Fig. 71.40 as a function of the pump repetition rate. The dioptric power for the plane containing the optic axis (crystal a axis) and the crystal c axis is labeled $\parallel c$ and for the plane containing the optic axis and the crystal b axis is labeled $\perp c$.

Figure 71.43

Dioptric power (inverse of the thermal-focal length in meters) calculated from the data in Fig. 71.41 as a function of the pump repetition rate. The dioptric power is labeled $\parallel c$ for the plane containing the optic axis (crystal a axis) and the crystal c axis and $\perp c$ for the plane containing the optic axis and the crystal b axis.

The dioptric powers in Figs. 71.42–71.44 can be written for the two orthogonal planes as¹²

$$D_{\perp,\parallel} = k_{\perp,\parallel} P_a / A, \quad (4)$$

where $D_{\perp,\parallel}$ are the dioptric powers in the planes perpendicular and parallel to the c axis, P_a is the total heat dissipated by the rod (equivalent in steady state to the total heat absorbed by the rod), A is the rod cross-sectional area, and the $k_{\perp,\parallel}$ are the dioptric-power coefficients defined in terms of the material parameters of Nd:YLF.

The heat dissipated by the rod was not measured; however, to compare our dioptric-power measurements to those stated in the literature, we will make similar assumptions: We assume that the total heat dissipated by the rod is proportional to the average electrical power dissipated in the flash lamps, where the proportionality constant involves many coupling efficiencies. We redefine the dioptric-power coefficients in Eq. (4) to include this proportionality constant. We can then calculate these newly defined dioptric-power coefficients using our dioptric-power measurements and the average electrical power dissipated in our xenon flash lamp. When our xenon

Figure 71.44

Comparison of the dioptric powers (inverse of the thermal-focal length in meters) as a function of the pump repetition rate measured for xenon-flash-lamp-pumped- and laser-diode-pumped Nd:YLF laser rods pumped to the same small-signal gain. The data have been fit to straight lines with the xenon-flash-lamp-pumped data shown as dashed lines and the laser-diode-pumped data (taken from Fig. 71.42) shown as solid lines.

flash lamp is operated at 60 Hz with 11 J per pulse, the average electrical power dissipated in the lamp is 660 W. This calibrates the axes in Figs. 71.42–71.44 to lamp average electrical power, and we can then state dioptric-power coefficients accordingly. Table 71.VII summarizes our dioptric-power coefficients $k_{\perp,\parallel}$, defined in this way, using the appropriate values for the rod areas.

It is coincidental that our dioptric-power measurements for our laser-diode-pumped rod are very similar to those reported by Cerullo *et al.*⁷ for their krypton-flash-lamp-pumped rod. It would be inappropriate, however, to base any conclusions on this since among other things the rods have different diameters and the coupling efficiencies are very different for the two cases. In fact, a true comparison of these data should include a detailed analysis of the cooling and coupling efficiencies mentioned above for each case, since the heat dissipated by the rod is the appropriate quantity to use in Eq. (4). This analysis would go far beyond the intent of this article.

The distortions in the xenon-flash-lamp-pumped rod were also measured at a higher gain. The capacitor-bank voltage was increased from 385 V to 460 V, corresponding to a 43% increase in pump-pulse energy from 11 J per pulse to

Table 71.VII: The coefficients of the thermal-dioptic power as defined in Eq. (4) and discussed in the text.

λ (nm)	Laser diode $k_{\parallel}$ (mm/kW)	Laser diode $k_{\perp}$ (mm/kW)	Xenon flash lamp $k_{\parallel}$ (mm/kW)	Xenon flash lamp $k_{\perp}$ (mm/kW)
1047	-0.767×10^{-3}	-2.76×10^{-3}	-0.565×10^{-3}	-2.38×10^{-3}
1053	-0.566×10^{-3}	$+0.153 \times 10^{-3}$	no data	no data

16 J per pulse. We found that the measured thermal distortions scaled linearly with this corresponding increase in average pump power as expected from Eq. (4). However, if the pump-pulse energy were increased significantly beyond this, the lamp's spectral output would change and different results would be expected.

The thermal-time constant of the distortions in our laser-diode-pumped rod was measured with the camera and VCR frame sequencer as described above. The data are shown in Fig. 71.45. The thermal distortions were measured using π -polarized light (gain at $\lambda = 1047$ nm) in our laser-diode rod pumped at a 200-Hz repetition rate as in Fig. 71.40. Pumping was removed at time equals zero in Fig. 71.45 (with the cooling liquid left on), and the thermal distortions were measured thereafter at known intervals (determined from the frame sequence numbers of the fringe data). The data points for defocus and astigmatism were fitted to exponential curves

(shown as solid lines in Fig. 71.45), and an exponential decay time constant (time for the aberrations to decay to $1/e$ of their value when pumping was removed) of approximately 1.5 s was obtained. These measurements were also made with the 1053-nm transition, and no significant change in the thermal-time constant could be measured.

Summary

We have made detailed interferometric measurements of the thermal distortions in xenon-flash-lamp- and laser-diode-pumped Nd:YLF laser rods. In both cases, defocus and astigmatism were the dominant thermal distortions. The thermal distortions in our flash-lamp-pumped rod were greater than the thermal distortions in our laser-diode-pumped rod when pumped to the same small-signal gain. Finally, we measured the thermal-relaxation time of our laser-diode-pumped Nd:YLF rod to be approximately 1.5 s.

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Figure 71.45

Thermal relaxation of the aberrations for π -polarized light (corresponding to gain at 1047 nm) in the laser-diode-pumped Nd:YLF laser rod. The rod was operated at 200-Hz pump repetition rate, and pumping was removed at time equals zero. The thermal-relaxation time is 1.5 s.

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The OMEGA Target-Positioning System

System Overview

Historically, several different styles of target positioners have been developed and deployed at LLE. The initial target positioner on the OMEGA 24-beam system used a 16-target carousel, which was located inside the experimental target chamber. This approach required that the entire target chamber be vented to load the target carousel; it further constrained the experimental operations by limiting target selection to one of the 16 on the carousel. This last constraint often limited the flexibility of the experimental program to change the target parameters during the shot sequence based on the results from the initial shots in the sequence. In addition, this carousel positioner did not allow rotational alignment of the target, which is required for flat-target experiments.

In the mid-1980s, a new style of target positioner was developed to support initial x-ray laser shots. It featured six axes of adjustment as well as a gas-feed capability; however, it also required that the entire experimental chamber be vented to load a single target. In addition, the six axes were configured in a manner that required complex compound motions for target alignment.

A third positioner developed for the OMEGA 24-beam system used a shuttle approach to load individual targets without requiring the experimental chamber to be vented. This positioner featured four axes of alignment motion without the complex compound motions of earlier positioners (see Fig. 71.46). Although this positioner was developed as an auxiliary positioner for backlighter targets, it soon became OMEGA's primary target positioner because it was reliable and provided quick turnaround between shots, while it maintained the maximum flexibility for the experimental program by allowing any available target to be used during a shot sequence. This positioner became the basis for the design of the target positioner for the OMEGA 60-beam system.

The OMEGA target-positioning system (TPS) has been developed to provide an accurate method for positioning a

wide variety of targets at the required location at or near the center of the experimental target chamber within the 1-h shot cycle of the OMEGA system. This target positioner provides four alignment motions, which include three translation motions (X, Y, Z) and one rotation motion about the axis of the

Figure 71.46 Shuttle-type target-positioning system: (a) target extended, (b) target retracted.

positioner (ω), with an option to add the remaining two rotation motions (θ , ϕ). (See Fig. 71.47 for the definitions of these motions.) The motion (T) used to transport the shuttle to the target chamber center (TCC) is independent of the fine alignment motions.

Figure 71.47
Definition of the OMEGA target-positioner axes.

The definition of TCC is provided by two nearly orthogonal target-viewing systems (TVS), which are rigidly attached to the chamber. The angle between the viewing systems is 101.5° , and the target positioner is 33.2° from being normal to the plane defined by the two viewing systems (see Fig. 71.48). This geometry does introduce coupling between the target-positioner motions and the two viewing systems; however, system operators learn to compensate and readily achieve target alignment.

A typical shot sequence begins with the target chamber evacuated to mid- 10^{-6} Torr; the target positioner is retracted and the air lock vented. A target is manually loaded onto the end of the target positioner, and the air lock is closed and evacuated below 1×10^{-4} Torr. The gate valve is opened, and the target positioner is inserted into the TCC. Once the target is located in the viewing systems, final alignment is done to $\pm 5 \mu\text{m}$ of the designated location using the fine positioning stages in the target positioner. Following the laser shot, the target positioner is retracted, the gate valve is closed, and the air lock is vented so that the process can be repeated.

Figure 71.48
OMEGA target chamber viewing system geometry. (View: normal to plane of viewing system.)

Requirements and Concept

A draft specification written in mid-1993 served as the basis for the mechanical design of the OMEGA target positioner. This specification was based on a shuttle-style device (Fig. 71.46), which was thought to provide the best combination of short cycle time and flexibility in target selection during an experimental shot sequence. The seven axes of motion are defined in Fig. 71.47, where the T and Z axes are along the axis of the target positioner. Experience has shown that the configuration of these axes (shown in Fig. 71.49) is important to minimize coupling of the motion of each axis. In addition, the θ and ϕ motions should have pivot axes as close as possible to the TCC to minimize the translation of the target when these rotations are adjusted. (The basic requirements of the target positioner specification are summarized in Table 71.VIII.) In addition, the bulk of the target positioner is to be located at least 12 in. from the TCC to minimize the amount of ablation due to the reaction products and energetic debris from the implosion.

The OMEGA target positioner is shown in Fig. 71.50. The target is attached to the fine-motion stages, which are supported on the transfer assembly (see Fig. 71.51); this unit moves within the guide tube assembly to the TCC. The guide tube assembly is the main structural member of the positioner and provides the vacuum housing for the air lock. This assem-

bly is attached to the chamber using a gate valve that allows the air lock to be independently vented. The target positioner currently deployed on OMEGA does not include the θ and ϕ

motions, although these motions are shown in Fig. 71.51. The sections of the target positioner are discussed in detail in the following sections.

Table 71.VIII: Required performance of the OMEGA target positioner.

Parameter	Specification	Comments
Required alignment motions	X, Y, Z, ω	θ and ϕ may be added
X, Y , and Z motion range	1 in.	Target within a 1-in. radius
X, Y , and Z motion resolution	5 μm (0.0002 in.)	Maximum step size
ω range	360°	Continuous
ω resolution	0.25°	Maximum step size
θ and ϕ ranges	$\pm 5^\circ$	Minimum
θ and ϕ resolutions	0.25°	Maximum step size
T resolution	1 mm (0.040 in.)	Maximum step size
Stability over 60 min	Within resolution	For all axes
Maximum target acceleration	0.2 g	For all axes
Transport time	Within 5 min	From airlock to TCC
Orientation	Any port on chamber

Figure 71.49
Configuration of OMEGA target-positioner axes.

Figure 71.50
OMEGA target-positioner components.

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Figure 71.51

Transfer assembly and fine-motion stages.

Transport Motion

The transport motion transports the target from the target positioner's air lock to the TCC—a distance of 80 in. This motion is provided by the following subsystems: a guide tube assembly, a transfer assembly, and a cable assembly.

1. Guide Tube Assembly

The guide tube assembly consists of five sections: inner guide tube, gate valve, cross, outer guide tube, and vacuum tee (see Fig. 71.50). The guide tube sections are 8-in.-outer-diam stainless steel tube. The inner guide tube is inside the target chamber and includes a standard 12-in. OMEGA flange, which seals to the target chamber. The cross and outer guide tube make up the air lock for the target positioner when the transfer assembly is withdrawn to load a new target. The gate valve separates the inner and outer guide tubes. Attached to the cross are standard ISO-style electrical-feedthrough vacuum flanges for powering the insertion and fine positioning stages (see Fig. 71.52); a quick access door for manually loading targets; and a viewport. The fourth flange on the cross will be used for an auto target loader, which will be added to the positioner in the future. The vacuum tee terminates the outer end of the guide

tube assembly and provides the interfaces to the roughing and vent valves, the turbo molecular high-vacuum pump, and the high-vacuum valve, which isolates this pump (see Fig. 71.53).

The guide tube assembly provides the structural support for the target positioner. This structure is approximately 94 in. long and is designed to provide the required stability of $\pm 5 \mu\text{m}$ ($\pm 0.0002 \text{ in.}$) at the center of the target chamber. Finite element analysis (FEA) was used to calculate the normal modes of this structure and determine how it would respond to vibration from other equipment attached to the target chamber. This modeling determined that an input of $100 \mu\text{g}$ at the target chamber flange would displace the target by $0.4 \mu\text{m}$ at the first mode of 22 Hz. Modeling determined that adding a support to the vacuum tee increased the first mode to 72 Hz and decreased the forced response to $0.2 \mu\text{m}$. However, in practice, this additional support has not been required on the target positioner currently deployed on OMEGA.

The guide tube assembly supports the cylindrical rails and drive rack for the transport assembly. The double rail design shown in Fig. 71.54(a) was chosen because it is simple and provides adequate straightness of motion. The technique for mounting the rails was taken from the six-inch manipulator (SIM) developed at LLNL and involves cross tapping the rails and using screws from the outside of the guide tube to secure the rail to the inside of the tube. The heads of the screws were then welded to the outside of the guide tube to maintain the vacuum integrity of the guide tube. This technique provided a simple fabrication method and yielded a stiff mounting for the rails after the screws are welded. One of the rails is spaced away from the inside surface of the tube [see Fig. 71.54(b)] to provide the proper clearance for the transport mechanism. The drive rack support is mounted to the inside surface of the guide tube using the same technique. The drive racks in the outer guide tube and the cross can be adjusted axially to allow them to be meshed with the drive rack in the inner guide tube. The

G4173

Figure 71.52

Guide tube assembly. (Note: The vacuum tee is not shown.)

G4174

Figure 71.53
Target positioner on the top port of the OMEGA target chamber.

drive rack has a small gap at the joint between the outer guide tube and the cross and a 2-in. gap across the gate valve. The guide rails also have a 2-in. gap across the gate valve. The transfer assembly is designed to bridge these gaps.

2. Transfer Assembly

The transfer assembly is made up of the transfer body and the gear housing and provides 80 in. of motion inside the guide tube [see Fig. 71.55(a)]. It is constructed with ball-bearing rollers that run on the rails in the guide tube. On one side of the transfer body the rollers form three V's; the opposite side of the transfer body has fixed lower rollers to take the spring preload on the drive gears, and spring-loaded upper and side rollers to prevent binding during the translation in the guide tube [see Fig. 71.55(b)]. Three sets of rollers on each side allow the transfer body to bridge the gap in the rails at the gate valve. The drive motion is provided by a small gear motor driving three drive gears through a worm gear reducer [see Fig. 71.55(c)]. The gear housing deployed with the current target positioner uses two drive gears rather than the three gears shown in Fig. 71.55(c). The two drive gears are separated by a distance that allows the unit to bridge the gap at the gate valve. The three drive gears are used to improve the smoothness of the motion across the gap in the drive rack at the gate valve.

G4175

Figure 71.54
(a) End view of the guide tube rails and drive rack; (b) closeup of the guide tube rails and drive rack.

(a)

(b)

(c)

G4176

Figure 71.55

(a) Side view of the transfer assembly; (b) end view of the transfer assembly; (c) closeup of the gear housing.

The drive motor for the transfer motion was initially a brush-style dc motor with integral gear head and magnetic encoder manufactured by MicroMo Electronics. The brushes have had poor lifetime in the high-vacuum environment, so the motors were replaced every 6 to 8 weeks to minimize the possibility of a failure during an experimental shot sequence. The transfer motion motor has now been replaced with a brushless dc motor manufactured by MicroMo [see Fig. 71.55(a)]. The gear head chosen for this application is a planetary style that provides robust performance and withstands the full stall torque of the motor. The gear head is cleaned and lubricated for vacuum service at LLE.

The motor is sized to provide a transit time to the TCC of less than 5 min. The weight of the shuttle is 25 lbs (the

estimated weight of the conceptual design was 27.3 lbs). The required torque at the motor/gear head was calculated from the estimated weight, and the motor's performance specifications (motor speed versus output torque) were used to calculate the shuttle's transit time for a range of torque settings. (These data are shown in Fig. 71.56.) The estimated torque for this application was ≈ 38 oz.-in. The transit time was estimated at 4.1 min and would increase to 4.3 min if twice the torque was required. Thus, the transit time is not a strong function of the torque required to move the shuttle.

The output shaft of the transfer motor is coupled to the drive gears with a 10:1 worm/worm-gear reducer, as shown in Fig. 71.55(c). This reducer provides the proper overall gear ratio for the drive motor and ensures that the transfer assembly

G4177

Figure 71.56

T-axis transit time versus torque required at the drive motor.

does not drive backward (due to gravity) when power is removed from the motor. The gear head on the motor is constructed with pressed-in shaft bearings, so the gear head cannot support the thrust loads generated by the worm/worm-gear reducer. These thrust loads required the worm to be mounted on a shaft, which is supported with adequate thrust-bearing capacity. The drive motor is coupled to the worm shaft by a slotted cup that engages the cross pin in the worm shaft. The reaction torque at the motor is taken by the two pins shown in Fig. 71.55(c). This arrangement allows the drive motor to be quickly removed from the gear housing and replaced in the event of a motor failure.

A belt drive was initially used to couple the two drive gears; however, this was abandoned because the belt could not be tensioned adequately for the load and the belt would slip as the transfer body moved across the gate-valve gap when the target positioner was mounted in a vertical orientation. A bevel-gear drive was also tried and dropped in favor of the compact and robust gear train shown in Fig. 71.55(c).

Motion along the *T* axis is restricted by limit switches that are activated at each end-of-travel [see Figs. 71.54(b) and 71.55(c)]. This motion uses an incremental encoder that must be zeroed at a repeatable location whenever power is lost to the controller. This switch is mounted with the end-of-travel limit switches and is tripped by the ramp for the retraction switch. In this manner, the *T* axis can be reset to home every time the transfer assembly is retracted from the TCC.

The brush-style drive motor initially used for the *T* axis incorporated a magnetic incremental encoder that produced 40 counts per revolution. This has been replaced with a brushless-

style drive motor, which uses an optical encoder that produces 384 counts per revolution. (See Table 71.IX for a summary of the resolution of the target-positioner motions.)

Table 71.IX: Calculated performance of the OMEGA target positioner.

Parameter	Specification	Final Value
<i>T</i> -axis resolution	1 mm	0.0125 mm
<i>T</i> -axis transit time	5 min	≈ 4 min
<i>Z</i> resolution	5 μm	0.1 μm
<i>X</i> and <i>Y</i> resolution	5 μm	0.8 μm
θ and ϕ resolution	0.25°	$0.6 \times 10^{-6}^\circ$
ω resolution	0.25°	0.004°

3. Cable Assembly

The cable assembly consists of two ribbon cables, a tensioning pulley, an interface printed circuit board (PCB) mounted to the transfer body, and two electrical vacuum feedthroughs mounted in the tee section of the guide tube assembly. This assembly provides power to the gear motors used for each of the alignment motions in the target positioner. It also provides feedback from the encoders and limit switches to the controller, which is remote from the target positioner. The tensioning pulley rides on the same rails as the transfer body. Two constant-force springs provide the necessary tension on the pulley to prevent the cables from tangling without producing excessive load on the transfer motion motor. The electrical feedthroughs are shown in Fig. 71.52 and the PCB in Fig. 71.57.

G4178

Figure 71.57

Printed circuit board on the transfer body.

$\pm 5^\circ$ of rotation about the TCC; historically, this was accomplished with goniometric stages. Goniometric stages are difficult to package in the space envelope of the target positioner, so an alternative approach was developed using four-bar linkages. The geometry of these linkages is shown in Fig. 71.61(a).¹ This concept supports the ω axis on the bar labeled c in Fig. 71.61(a), so the rotation angle of the target located at the end of the ω axis is defined as θ . The law of cosines yields the following relationships for the four-bar linkages:

$$\psi = \cos^{-1} \left(\frac{h^2 + a^2 - b^2}{2ha} \right) + \cos^{-1} \left(\frac{h^2 + d^2 - c^2}{2hd} \right), \quad (1)$$

$$\tau = \cos^{-1} \left(\frac{c^2 + d^2 - a^2 - b^2 - 2ab \cos \phi}{2cd} \right), \quad (2)$$

$$h^2 = a^2 + b^2 + 2ab \cos \phi, \quad (3)$$

where ϕ is referred to as the crank angle.

The target's angle of rotation is

$$\theta = \pi - (\tau + \psi). \quad (4)$$

The location of a point P relative to the origin shown in Fig. 71.61(a) is

$$P_x = b \cos(\phi) + r \cos(\theta + \alpha), \quad (5)$$

$$P_y = b \sin(\phi) + r \sin(\theta + \alpha). \quad (6)$$

In the design of the θ and ϕ mechanisms the links shown as b and d in Fig. 71.61(a) are equal in length, and the geometry of the mechanism for the OMEGA target positioner, as shown in Fig. 71.61(b), yields the following relationships:

$$c = u \left(\frac{a}{u + w} \right), \quad (7)$$

$$b = \frac{w}{\cos \left[\tan^{-1} \left(\frac{a/2}{u + w} \right) \right]}, \quad (8)$$

$$\alpha = \tan^{-1} \left[\frac{u}{c/2} \right], \quad (9)$$

$$r = \frac{u}{\sin \alpha}, \quad (10)$$

where the linkage is designed to provide rotation about the remote point P , shown in Fig. 71.61(a). By defining the variables a (width between the fixed pivots), w (height of the mechanism), and u (distance from the mechanism to the

Figure 71.61
(a) Four-bar linkage geometry; (b) θ - and ϕ -stage geometry.

target), one may numerically solve for the angle of the target, θ , as a function of the crank angle ϕ by using Eqs. (7)–(10) and Eqs. (1)–(4). This has been done to understand the motion for various design dimensions, and the final design parameters are given in Table 71.X for both the θ and ϕ stages. The relationship of the rotation angle of the target to the rotation of the crank is nearly linear, as shown in Fig. 71.62. Equations (5) and (6) are used to calculate the movement of the target in the axial and lateral directions over the rotation range of 5° . These results are given in Figs. 71.63(a) and 71.63(b).

The θ and ϕ stages have been built in a nested fashion to minimize the size of the assembly and allow it to fit in the allotted space between the X/Y stages and the ω stage [see Figs. 71.64(a) and 71.64(b)]. Both the θ and ϕ stages are driven by the same dc gear motors as the X and Y stages. These motors are coupled to the stage crank shafts using a worm/worm-gear reducer. The θ motor is packaged within the Z-stage yolk (see Fig. 71.64); the ϕ -motor location is shown in Fig. 71.64(b). The final resolution of the θ and ϕ stages is listed in Table 71.IX.

Table 71.X: Final design parameters for the θ and ϕ axes.

Parameter	θ Stage	ϕ Stage
u (in.)	12.00	11.97
w (in.)	3.50	3.00
a (in.)	1.25	1.25

Figure 71.62
 θ and ϕ rotation versus crank angle.

Controls

1. Operator Interface

The experimental system operator (ESO) controls the TPS/TVS through a graphical user interface (GUI) residing on a computer terminal located in the control room (see Fig. 71.65). This GUI displays push buttons for control of TPS motions including selection of jog modes, direction, and jog-step size. Values representing the actual (implied) absolute position of each axis are useful. In addition, feedback of actual position within range (and/or limit-switch status) is displayed for each axis. The window provides for actuation of shutters, selection of illumination intensity, magnification option, and enhanced video options, etc., for TVS control.

2. Target-Positioner Control Elements

The OMEGA target positioner is comprised of two discrete and operationally independent functional areas (see Fig. 71.66): (1) the auxiliary vacuum controls that control entry into the target chamber while remaining under a high vacuum condition, and (2) the multiple-axis motion controller necessary for precision motion control of up to six individual axes.

Figure 71.63
(a) θ and ϕ axial position error; (b) θ and ϕ lateral position error.

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Target-Positioner Auxiliary Vacuum Controls

The primary function of the auxiliary vacuum controls is to sequence the pumps, solenoid valves, and an 8-in. gate valve such that the positioner, loaded with a target, may be inserted inside the target chamber without perturbing the chamber's high-quality vacuum condition. The process involves network communications with other auxiliary vacuum controllers to arbitrate the sharing of the roughing pump manifold resource. In addition, the auxiliary vacuum controls provide a remote control and monitor of TPS functions both in the control room and locally at the target chamber's personnel platform.

During airlock evacuation the auxiliary vacuum controls monitor a vacuum sensor to automatically sequence from rough pumping to turbo pumping operation. Once the airlock is properly evacuated and insertion has been commanded, the gate valve is opened, enabling the transport axis motion. The controls establish the "home reference position" to the transport motion axis to maintain coarse positioning accuracy to within ± 0.001 in. over an 80-in. range of travel in order to arrive within the video capture range of the TVS.

Target-Positioner Precision Motion Controller

OMEGA's TPS motion controller provides the control elements for the X, Y, Z, θ , ϕ , and ω axes that provide precise alignment and rotation of the target to the center of the

Figure 71.64

(a) θ and ϕ stages; (b) ϕ -stage drive-motor mounting.

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Figure 71.65

TPS graphical user interface.

Figure 71.66
Target-positioner control elements.

chamber. These six independent functions are provided from a commercially available unit purchased from Galil Motion Control, Inc. Motion instructions are issued via a serial data communication link from the experimental executive computer to the Galil motion controller during alignment operations. The Galil instruction set is based on an ASCII format and includes programmable variables for servo parameters such as gain, velocity, acceleration, and filter terms. These and other parameters are retained in static zero power RAM for convenience.

A motion controller is a part of a closed-loop servo system, which also includes a power amplifier, motor, and encoder. A typical closed-loop-position servo system is shown in Fig. 71.67. The controller is the element that initiates motion by generating an analog command signal that, when applied to the power amplifier, results in an electric current flowing in the motor windings. This current creates torque and subsequent mechanical motion. The encoder senses the position change in the motor and feeds this information back to the controller, which tracks the motion progress over time.

Incremental encoders operate by generating output pulses that represent the motor shaft position and include two signals that, for the case of the TPS, generate ten pulses per revolution. The two signals are displaced by 90° in phase, which enables the controller to determine the direction of rotation according to which signal leads or lags the other. Motor displacement can therefore be determined by counting the number of pulses, and the frequency of the resultant pulse train is proportional to motor velocity.

The controller performs the basic intelligent operations of the system:

- interpreting commands sent from the host computer (experimental executive) via a serial communication interface,
- generating the command signal profile,
- decoding position-feedback information into motion displacement and velocity,

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Figure 71.67
Typical closed-loop-position
servo system.

- closing the position loop,
- compensating the loop by filtering to achieve stability and prevent oscillation, and
- interpreting end-of-travel limit switches to prevent possible damage to the motion mechanisms.

Command Profile

The target is suspended at the end of the target-positioner mechanism from a carbon fiber, which is several microns in diameter. The mass of the stalk is kept intentionally low, and the resultant thin structure is extremely fragile. Consequently, the controller must deliver smooth and reliable motion to minimize (lateral) forces applied to the target and its stalk during alignment procedures. Sudden TPS accelerations or decelerations must be limited to levels that ensure the stalk does not break or the target does not become dislodged from the stalk.

A “trapezoidal velocity profile” applies a linear function to the motor that provides continuous acceleration to maximum velocity and continuous deceleration to a full stop [Fig. 71.68(a)]. However, a more complex command profile known as an “S curve” [Fig. 71.68(b)] is sometimes preferred because it imparts a smoother, less-instantaneous acceleration change.

To date, the TPS controller has utilized the trapezoidal velocity profile successfully. However, should future targets or their stalks become more sensitive to acceleration forces, the controller is capable of generating S-curve velocity profiles.

Compound Coordinated Motions

The mechanical design of the ϕ and θ axes couples ϕ rotation with X and Z motions and θ rotation with Y and Z motions as discussed in a previous section. If the ϕ or θ axes are moved to the full range of motion, the target will shift outside the capture range of the TVS. For this reason a motion controller

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Figure 71.68
(a) Trapezoid velocity profile; (b) S-curve velocity profile.

was selected that is capable of implementing electronic gearing where one or more axes are slaved to follow a master axis at a fixed ratio, thereby providing a motion vector.

Wiring Isolator

Early operational experience with the OMEGA target positioner confirmed that the supporting electronics systems and components were prone to damage from the laser-fusion event. During the event, laser energy is incident upon the target and the tip of the target stalk. This creates a transient high-temperature plasma comprised of charged particles, some of which recombine while others fly in all directions. The electrons, being the lightest, travel at the highest velocity to the target chamber walls. The target-positioner mechanism being the closest conductive mass near the center of the target chamber intercepts a significant portion of these charges, which can couple back into control circuitry of the TPS. The transient nature of this electrical phenomenon occurring with typically nanosecond rise times can damage semiconductor components and perturb information stored in memory devices. These effects require that the wiring and cabling in the target chamber interior be shielded from line-of-sight particles and that the remote electronics be isolated from the TPS using a disconnect switch. Just prior to the shot, relay contacts open in each of the 70+ conductors in the wiring harness to protect the control electronics. After the shot the relay is commanded closed once again, thereby enabling the TPS retrieval process.

Summary

The OMEGA target positioner has been developed to provide a means of positioning a target at the TCC and precisely aligning this target using up to 4° of freedom, X , Y , Z , and ω . A compact two-axis stage was developed to provide θ and ϕ rotation of the target, but this stage is not currently used on OMEGA.

All motions are servo controlled, and a GUI is used in the OMEGA control room to provide convenient control of the target positioner.

The OMEGA target positioner has been used to support and align a wide variety of targets in the experimental chamber including spherical implosion targets, multiple flat-foil instability targets, and hohlraums for indirect-drive experiments.

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OMEGA Experimental Program and Recent Results

The Laboratory for Laser Energetics (LLE) with its OMEGA facility is part of the national laser-fusion effort within the U.S. Department of Energy and has as its main mission direct-drive laser-fusion research in support of the upcoming National Ignition Facility (NIF) at the Lawrence Livermore National Laboratory (LLNL). The upgraded OMEGA laser system has been in operation since May 1995. This versatile 60-beam UV laser facility has proven capable of carrying out a wide variety of experiments related to inertial confinement fusion (ICF). The facility also supports indirect-drive laser-fusion research under the direction of LLNL and the Los Alamos National Laboratory (LANL).

The challenges facing direct-drive ICF are principally connected with drive uniformity (uniformity of laser irradiation on target), control of the hydrodynamic instabilities during the compression phase of the target, and target uniformity, particularly that of the future cryogenic targets required for the success of laser fusion. The experimental program at LLE addresses all of these issues; this article lays out the present status of this research and its path for the future.

The *irradiation uniformity* on target depends on the number of laser beams and their disposition, the intensity distribution within each beam (beam-smoothing scheme), and the energy and instantaneous power of each beam (energy and power balance). While the first two points are unalterable characteristics of the OMEGA laser system (60 beams symmetrically disposed around the target), all other contributors to the irradiation nonuniformity are subject to ongoing research and improvements. The ultimate goals for OMEGA are 4%–5% rms beam-to-beam energy and power balance and <1% rms irradiation uniformity averaged over the target and over 300 to 500 ps.

Understanding and controlling the hydrodynamic instabilities are the main ICF concerns. Of primary concern is the Rayleigh–Taylor (RT) instability in the acceleration and deceleration phases of the spherically imploding ICF target. The problem centers on both growth rates and relevant seed amplitudes from which target perturbations can grow. LLE has

established experimental programs to address both issues. The seed amplitudes can be due to either laser nonuniformities (laser imprinting) or target nonuniformities. The role of drive-generated seed amplitudes (laser imprinting) is studied extensively at LLE. Assessment and improvements of target nonuniformities (particularly those of cryogenic targets) are addressed elsewhere within the national program. RT growth rates have been studied in a collaborative effort between LLNL and LLE for some time. The present LLE experimental program on RT growth rates addresses the problem in both planar and spherical geometries and for a variety of irradiation and laser-beam-smoothing techniques.

LLE is preparing for future cryogenic target irradiation through a program centered at General Atomic in San Diego. Cryogenic targets will be irradiated on OMEGA within two to three years.

OMEGA Laser Characteristics

The OMEGA laser system¹ is a Nd:glass laser system with frequency conversion to the third harmonic with a UV (351-nm) energy capability on target in excess of 30 kJ for pulse durations of ≤ 2 ns. The system's 60 beams (30-cm beam diameter) are arrayed symmetrically around the 3.3-m-diam target chamber and are focused onto the target by 60 near-diffraction-limited $f/6$ lenses.

Laser pulses with a wide variety of predetermined, temporal UV output pulse shapes have been produced by a versatile pulse-shaping system.^{2,3} To date these shapes range from 1- to 3-ns flat-topped pulses to linear ramps (1 to 3 ns), Gaussians (0.2 to 1.2 ns), and LLNL's pulse shape PS22 (2-ns pulse designed for indirect-drive targets).

The demonstrated energy balance of the OMEGA system is ~2% rms for the IR part of the system, ~3% rms in the UV after the conversion crystals, and ~4% to 5% rms on target after the UV transport optics (two turning mirrors, a focusing lens, and a blast shield per beamline). The system's energy reproducibility is <1% rms excluding clearly identifiable flashlamp or

pulse-forming network (PFN) malfunctions. The measurement accuracy of the beam energies (including residual fundamental and second-harmonic output from the conversion crystals) is $\sim 0.8\%$ rms.

The laser system has three complete driver lines with corresponding pulse-shaping systems. Any one or any combination of two of these driver lines can be used to drive the laser system. The laser system supports the propagation of two different pulse shapes, one for the main experiment (driver) and one for a backlighter. In this case one driver line feeds 40 beamlines, while the other feeds the remaining 20. The timing between these two groups of beams can be chosen freely.

Beam smoothing—an essential component for the success of direct-drive laser fusion—is implemented on OMEGA using smoothing by spectral dispersion (SSD) pioneered at LLE.⁴ LLE's approach combines continuous-phase-plate technology with polarization rotators and two frequency modulators with corresponding spectral dispersion in two dimensions (2-D SSD). Present capabilities are limited to 2- to 3-Å total bandwidth for the two modulators (3- and 3.3-GHz modulation frequency). In addition, a limited number of phase plates and polarization rotators are available primarily for planar-target experiments. By the end of 1997 a full complement of phase plates will be available for spherical irradiation experiments. A complete complement of polarization rotators is projected for 1 to 2 years from now. At that time increased bandwidth (~ 10 Å) and frequency modulation (FM) at higher frequencies (~ 10 GHz) will also be available, allowing an on-target frequency bandwidth of ~ 1 THz in the UV.

The pointing capability of the OMEGA laser system has been shown to be ~ 10 μm rms on target. This performance has been repeatedly verified with x-ray imaging using pinhole cameras and x-ray microscopes. The stability of the system is sufficient to maintain the above pointing accuracy without adjustments for at least a period of one day. Furthermore, recent experiments have shown that beams can be placed in basically arbitrary locations within ~ 1 cm from the target chamber center with essentially the same accuracy (~ 20 μm without corrections; better accuracy is achievable with one iterative fine correction). This capability has been demonstrated on recent indirect-drive hohlraums of cylindrical (two laser entrance holes) and spherical shape ("tetrahedral" hohlraums with four laser entrance holes) with or without additional backlighter targets.

Present OMEGA Experiments

The present OMEGA experiments address irradiation uniformity, implosion physics, hydrodynamic instabilities, laser imprinting, and laser-plasma-interaction physics. In addition LLNL and LANL carry out experiments related to indirect-drive laser fusion.

1. Irradiation Uniformity

The irradiation uniformity on target can be analyzed using spherical harmonic decomposition. One thus finds that the unalterable OMEGA irradiation configuration (60 beams symmetrically arrayed around the target) and the on-target beam envelopes contribute primarily to modes of order $\ell < 20$. Energy and power balance as well as pointing on target contribute predominately to the amplitudes of modes with $\ell \leq 4$. The intensity distributions of the individual beams, on the other hand, contribute to high-order modes with ℓ ranging well beyond $\ell = 100$. The effort to control and reduce the irradiation uniformity on target therefore requires many different approaches.

a. Energy and power balance. Energy and power balance of $\sim 4\%$ – 5% rms between beams puts severe demands on the laser system. Very accurate energy- and pulse-shape-measurement capabilities are a prerequisite, as are stringent controls over the gains and losses of like amplification stages within the system (they must be kept within $\sim 2\%$ rms). In addition, computer-based strategies must be developed for energy redistribution and fine control of some key amplifiers.

The energy-measurement requirements have been met by OMEGA. The capability to sufficiently control the individual beam energies has also been demonstrated; however, some persistent problems with PFN's presently prevent consistent achievement of the desired and achievable energy balance of $\leq 4\%$ rms on target. This goal will only be achieved after some critical components that currently perform marginally are replaced.

Power balance on OMEGA has so far been addressed primarily conceptually. Some significant improvements in pulse-shape-measurement capability will result when two high-quality multiplexed UV streak cameras are installed on OMEGA within the next two to three months. They will measure the pulse shapes of 20 of the 60 OMEGA beams. The expected measurement accuracy is $\sim 1\%$ rms at the peak of the pulse for all 20 channels with an optical fiducial channel providing

timing accuracy and on-shot sweep speed calibration. (In addition, a new streak camera tube, designed at LLE and built by Phillips in France, is expected to provide a higher dynamic range and more extended multiplexing capabilities.)

An essential component of LLE's power balance strategy relies on the fact that power balance is assured if energy balance is maintained for full-energy shots as well as for shots with reduced output energy. Of course, the laser system must be operated under identical conditions (e.g., identical bank voltages) with only the input laser energy reduced appropriately.

b. Beam-intensity distribution and smoothing mechanisms.

The intensity distribution of a typical high-power laser beam on target is shown in Fig. 72.1(a). The target outline is indicated by the circle. The introduction of a distributed (quasi-random) phase plate (DPP) with continuous phase relief generates an intensity distribution [Fig. 72.1(b)] whose envelope and speckle pattern are statistically well defined. In addition, LLE's phase-plate approach allows for control of both the envelop function and the speckle distribution. SSD combines FM with spectral dispersion to smooth the high-frequency speckle of the DPP's leading to the very smooth intensity distribution shown in Fig. 72.1(c). This figure is a temporal average over the full pulse duration of 1 ns. Instantaneously, the intensity distribution is still as shown in Fig. 72.1(b).

Further enhancement of the uniformity (smoothness) of the single-beam-intensity distribution on target is obtained by using polarization rotators.⁵ These optics are birefringent wedges that are placed in the beam near the focusing lens and change the polarization across the beam in a continuous manner. The effect of this plate is equivalent to splitting the beam into two orthogonal polarization components whose indepen-

dent speckle patterns on target are slightly displaced and do not interfere with each other. This leads to an instantaneous reduction in beam nonuniformity by a factor of $\sqrt{2}$. LLE is currently developing liquid crystal polarization rotators for all 60 beams of OMEGA.

2. Implosion Physics

Implosion experiments have been conducted on the upgraded OMEGA since it was commissioned in May 1994. The improved symmetry of irradiation and the 30-kJ energy capability of this laser system immediately resulted in record neutron yields in excess of 10^{14} for ~30-kJ UV irradiation of a DT-filled glass shell of ~1-mm diameter.

A series of recent implosion experiments were designed to investigate the scaling of the measured neutron yields with expected core distortions due to the combined effects of the RT instability in the ablation and deceleration phases. For this purpose, targets with different fill pressures and various shell compositions (CH and Cl-doped CH shells) were used. A corresponding plot is shown in Fig. 72.2, where the relative neutron yield is plotted as a function of the expected distortion fraction of the fully imploded target. The distortion fraction was obtained from ancillary calculations using current models for the RT growth factor in the acceleration (ablation) phase, a heuristic feedthrough factor for the perturbations from the outer to the inner pusher surfaces, and similar RT growth factor estimates for the deceleration phase. This distortion fraction is the radial core distortion divided by the final radius of the imploded core. The smooth scaling of these data for a variety of targets and irradiation conditions suggests that this definition of the distortion fraction of the final core is a useful one, combining RT instability and target convergence concepts.

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Figure 72.1

The intensity distribution of a single beam in the target plane for (a) a typical high-power laser beam without any smoothing, (b) the same beam as modified (aberrated) by a phase plate with a continuous-phase-relief pattern, and (c) 2-D SSD (phase plate and two FM modulators with wavelength dispersion in two perpendicular directions).

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3. Hydrodynamic Instabilities

The hydrodynamic instabilities encountered in direct-drive ICF are generally dominated by the RT instability. This instability is active during the target implosion in both the acceleration phase at the ablation interface and the deceleration phase at the fuel–pusher interface before the final stagnation. The instability grows from target imperfections and/or mass perturbations produced by irradiation nonuniformities (laser imprinting). This laser imprinting occurs primarily at early times when there is little or no plasma to smooth irradiation nonuniformities. Since some degree of irradiation nonuniformity is unavoidable, a thorough understanding of the RT growth rates is vital, particularly with regard to effects that may reduce the growth rates.

Figure 72.2

The yield over calculated 1-D yield (YOC) scales reproducibly with the distortion fraction, which is defined as the distortion of the final core relative to its radius. This distortion fraction is obtained from ancillary calculations using models for the RT growth factors in the ablation (acceleration) phase, a heuristic feedthrough factor, and the growth factor of the deceleration phase.

Several experimental programs at LLE investigate various aspects of the RT instability: two planar-foil experiments study laser imprinting and RT growth rates respectively, and several spherical target experiments address RT growth in both the ablation and deceleration phases.

a. Laser imprinting (planar targets). Laser imprinting has been studied at LLE using planar foils and multiple-beam irradiation (typically five to six overlapping beams). The experimental variables included various beam-smoothing techniques and foam-coated targets. Face-on radiography (85-ps exposure time) was used to measure the modulation depth as a function of time (Fig. 72.3). Initial mass perturbations caused by laser imprinting are amplified by the subsequent RT growth and are observed later in time. In other words, the mass perturbations that are initially invisible are rendered visible by the RT amplifier. Changing the smoothness of the laser irradiation (from no DPP's to DPP with 1-D SSD, etc.) one clearly discerns a corresponding decrease in modulation depth at late times. Full 2-D SSD with distributed polarization rotators (DPR's) leads to greatly reduced modulation, indicating significantly reduced laser imprinting. Fourier analysis of the images in Fig. 72.3 also shows clear quantitative reduction in the power spectrum for spatial wavelengths in the range of 20 to 100 μm .⁶

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Figure 72.3

Laser imprinting using planar foils and five overlapping beams (3 ns flat-top) clearly shows the effect of different smoothing techniques applied to the laser beam. The modulation face-on radiographs taken in the light of an x-ray backlighter around 2 keV are less pronounced at early times than later (2.1 ns versus 2.4 ns) when RT growth of the laser imprint becomes much more noticeable.

Foam-covered targets have been proposed to decrease laser imprinting as well as reduce RT growth rates.⁷ Imprint experiments using such foam-covered targets do indeed exhibit significantly reduced laser imprinting, but the analysis of these experiments is rendered more complex due to possible changes in target isentrope, which can reduce the RT growth rates. To assess the utility of foam-covered targets all effects on target dynamics must be considered.

b. RT growth (planar target). Detailed studies of the RT growth rates in direct-drive experiments have been carried out for some time in collaboration with LLNL.⁸ These experiments were initiated on the Nova laser at LLNL and were subsequently moved to OMEGA.⁹ The generic setup for these experiments utilized flat-foil targets with known mass perturbations that were irradiated with five or six overlapping drive beams. These drive beams were outfitted with DPP's and other variants of beam smoothing. The targets were radiographed with an x-ray backlighter and an x-ray framing camera. In addition, acceleration measurements of the foil were taken with a streak camera and side-on x-ray backlighting.

Quantitative interpretation of the RT growth rates requires accurate knowledge of the laser coupling to, and the resulting

acceleration of, this target. Furthermore, since all RT growth rate interpretation requires code simulations, a prerequisite for success is good agreement between the time history of the accelerated target position and code predictions. Such agreement is shown in Fig. 72.4 for a bare (smooth) 20- μm CH target foil accelerated by a 3-ns ramp laser pulse. Similar experiments using foam-covered targets yielded equally good results and indicated unchanged coupling of the laser to the target.⁹

The RT growth of single-frequency sinusoidal perturbations impressed on the target has been measured using initial perturbation amplitudes of $\sim 0.5\text{-}\mu\text{m}$ with 31- μm and 60- μm wavelengths. Initially these perturbations are invisible in the framing images of the backlit foil, but at later times the perturbations become clearly visible and measurable. Plotting the optical depth versus time (Fig. 72.5) shows the expected signature of RT growth. Since the optical-depth fluctuations are proportional to areal density (ρr) changes along the line of sight, one interprets modulation fluctuations as perturbation amplitudes assuming constant density ρ . In order to extract quantitative information from this plot, numerical simulations are carried out using the 2-D hydrodynamic code *ORCHID*. Good agreement between these simulations and the experimental data in Fig. 72.5 is obtained provided the simulations use measured laser pulse shapes and include the experimentally measured noise in the x-ray framing camera snapshots. If

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Figure 72.4

Trajectory of accelerated 20- μm CH foil targets used for RT experiments. The experimental data were obtained using an x-ray streak camera with side-on backlighting of the planar-foil target. The solid line corresponds to numerical predictions using the experimental laser pulse shapes.

E8402

Figure 72.5

Rayleigh-Taylor growth as measured by the optical depth ($\sim \rho r$ of the target) for a 20- μm -thick CH foil with an impressed sinusoidal corrugation of 31- μm wavelength and 0.5- μm amplitude. The perturbation amplitude onset and saturation are well modeled by *ORCHID* simulations, provided the experimental noise in the framing-camera images is included in the analysis (dashed line).

the experimental noise were neglected, one might infer reduced (and erroneous) RT growth rates. Furthermore, proper interpretation of RT experiments requires allowance for nonconstant target acceleration and the concomitant change in RT growth rates.

c. Ablative RT instability in spherical targets. Burnthrough experiments using imploding spherical targets with high-Z signature layers have been carried out at LLE for approximately five years.¹⁰ These experiments have clearly shown that the measured time of appearance of the x-ray lines due to the signature layer is more indicative of the RT instability in the ablation phase than of the propagation of the ablation front. Recent experiments on OMEGA have confirmed the earlier observations and refined the interpretation in terms of equivalent initial mass perturbations, which may be due to laser imprinting or target nonuniformities. Changing beam smoothing permits quantitative estimates for the laser imprinting in terms of equivalent target nonuniformities.

d. Deceleration-phase RT instability (mixing). The RT instability during the final stages of the implosion (the deceleration phase) has been investigated at LLE using various signature layers close to the inner boundaries of the target.^{11,12} These layers are either Ti-doped plastic layers with a 10- to 20- μm CH ablation layer on the outside or 2- μm CD layers similarly overcoated with an ablation layer. Between the signature layer and the inner target surface are CH layers of variable thickness. The target fill is either H_2 -Ar or D_2 -Ar for the two types of targets, respectively. The performance of these targets may be analyzed by using their neutron yield as a function of the innermost layer thickness or by x-ray spectroscopic means. These experiments are still in progress, and results indicate that this will be a powerful experimental technique.

4. Laser-Plasma-Interaction Physics

A new series of laser-plasma-interaction physics experiments have recently been initiated at LLE. The first results pertain to scattering near the incident laser wavelength where stimulated Brillouin scattering (SBS) may be expected. Under current OMEGA irradiation conditions for spherical targets, 5% to 10% of the incident laser light is isotropically scattered due to refraction of the 60 OMEGA beams. The evolving plasma corona of the imploding targets imparts considerable bandwidth to this scattered light, which, however, bears no relationship to SBS. On the other hand, clear signatures of SBS have been observed in implosion experiments when one beam was focused to $\sim 10^{15} \text{ W/cm}^2$ ($R_{\text{SBS}} \approx 12\%$).

The presence of a vigorously excited two-plasmon-decay instability is clearly seen in all implosion experiments on OMEGA. This instability has an experimental threshold in the low 10^{13} W/cm^2 and is difficult to suppress with any beam-smoothing techniques.¹³ The importance of this instability lies in its potential to generate energetic electrons. We are presently preparing to quantify any energetic electron production that may accompany this instability.

In the spectral region where one expects to observe stimulated Raman scattering (480 to 700 nm), we have observed scattering in a number of implosion experiments. However, indications are that this instability is close to its threshold under typical OMEGA operating conditions (a few 10^{14} W/cm^2).

Long-scale-length experiments have begun using exploding foil or solid targets with multiple-beam irradiation and staggered timing of the various beam clusters. These experiments are guided by 2-D hydrodynamic (*SAGE*) simulations that predict ~ 1 -mm density-gradient lengths at $n_e/8$ at electron temperatures in excess of 2 keV.

Summary

The 60-beam OMEGA laser system has now been in operation for three years. It has been used extensively for a wide variety of experiments relating to direct-drive laser fusion, from high-yield implosion experiments to planar and spherical Rayleigh-Taylor experiments and laser-imprinting experiments as well as laser-plasma-interaction experiments. OMEGA is a versatile and effective target irradiation facility that will support the National Ignition Facility now under construction at LLNL. Experimental results obtained to date pertain to both direct-drive and indirect-drive laser fusion.

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Imaging the Cold, Compressed Shell in Laser Implosions Using the $K\alpha$ Fluorescence of a Titanium Dopant

The cold, compressed shell in laser-driven implosions can normally be imaged only by backlighting because its x-ray emission is very weak. However, a high- Z -doped shell can be imaged using the $K\alpha$ line radiation, which fluoresces due to excitation by the intense radiation from the hot core. We show results from titanium-doped target implosions where the one-dimensional spatial profile of the Ti $K\alpha$ radiation yields the dimensions of the cold shell around the time of peak compression. This result, coupled with the shell areal density determined by the K -edge absorption, yields an estimate of the shell density. Since $K\alpha$ emission depends only on the cold shell density (not its temperature), its image directly reflects the shell uniformity. Additionally, the measured absorption above the K -edge and the energy in the $K\alpha$ line, which are directly related, are shown to indeed be consistent.

Laser-imploded targets consist of a hot core surrounded by a cold, compressed shell. Diagnosing the shell is vital to understanding target implosion, and several x-ray-based methods have been developed to address this requirement. For example, it has been shown^{1,2} that the peak in the emergent continuum radiation (caused by shell absorption of the soft radiation) can be used to determine the shell areal density. More recently^{3,4} it has been shown that doping part of the shell with a high- Z element provides additional signatures of shell temperature and areal density, namely, absorption lines and K -edge absorption [including extended x-ray absorption fine structure (EXAFS)]. All these methods depend on the core radiation; thus, a time-integrated measurement yields approximately the conditions at peak compression, when that radiation is strongest.

Additional valuable information can be provided by imaging the cold shell. Since a cold shell emits few x rays, its imaging is normally done by backlighting.⁵⁻⁷ However, we show in this article that when the cold shell is doped, it can be imaged using $K\alpha$ fluorescence excited by the core radiation. The fluorescence of $K\alpha$ radiation can occur only in cold material and is significant only when the pumping core radiation is intense; thus, the $K\alpha$ emission profile delineates the

cold shell surrounding the hot core around the time of peak compression. The term "cold" in this context will be defined below. This imaging method is much simpler than backlighting and has an additional advantage: whereas in backlighting an intense core radiation competes with the signal, here, to the contrary, it enhances it. $K\alpha$ emission in past experiments has been associated with excitation by suprathreshold electrons⁸ or by coronal radiation.⁹ However, the data in the present experiment clearly shows that the excitation source for the $K\alpha$ line is the core radiation.

$K\alpha$ fluorescence can be imaged by monochromatic, two-dimensional imaging; a method involving a Kirkpatrick-Baez (KB) microscope coupled to a diffracting crystal was recently demonstrated by Marshall and Oertel.¹⁰ Here we report on one-dimensional polychromatic imaging obtained with a crystal spectrograph fitted with a spatially resolving slit. One-dimensional imaging cannot show the symmetry of the cold shell, but it can determine its average radial extent. Coupled with the measurement of the shell's areal density by any of the above methods, this result can lead to an estimate of the average shell density around peak compression. Also, the measured spatial extent of $K\alpha$ radiation provides an additional constraint on two- and three-dimensional simulations of target implosion. The purpose of this work is to demonstrate that $K\alpha$ images can be obtained with targets of low-level doping and that a cold shell is indeed found to surround the hot, compressed core.

Experimental Results

Figure 72.6 shows a schematic of the target and measurement configurations. The polymer target is initially empty; at peak compression, part of the overcoat has been ablated and the inner part of the mandrel forms the core. The doped layer contains titanium at 0.5% by atom number. An axial view by the x-ray spectrometer shows the emission of continuum from the core, as well as K -edge absorption and $K\alpha$ emission within the doped layer. *An off-axis view should show only the $K\alpha$ emission.* The overcoat is much thicker than the laser-ablated layer, ensuring that none of the titanium-doped layer is ablated away. Likewise, the titanium layer is sufficiently far from the

Figure 72.6

Schematic representation of the target and measurement configurations. An axial view by the x-ray spectrometer shows the core emission of continuum as well as K -edge absorption and $K\alpha$ emission within the doped layer. An off-axis view shows only the $K\alpha$ fluorescence emission.

inner-shell interface to ensure no emission of titanium lines in the absence of mixing.⁴ The titanium layer is thus cold for the most part and should give rise to $K\alpha$ fluorescence.

We analyze the results from a single target shot, No. 9095, on the OMEGA laser system.¹¹ The target was a 6.1- μm -thick, empty CD mandrel of 931- μm outer diameter, coated with a 8.7- μm -thick CH shell that was doped with titanium at 0.5% concentration by atomic number (or at a density of 0.037 g/cm³). Finally, the target was overcoated with an 18.7- μm -thick CH layer. The 29.6-kJ laser pulse had a 1-ns flat top (to within $\pm 5\%$) with 100-ps rise and fall times. Smoothing by spectral dispersion (SSD)¹² was used in this experiment, but not distributed phase plates (DPP's).¹³

The x-ray spectrum in the range of ~ 4 to 9 keV was recorded by a Ge(111) crystal spectrometer fitted with a 100- μm -wide spatially resolving slit. The crystal and film calibration methods used in the data reduction are explained in Ref. 4. Spectrally dispersed images were also obtained with a KB x-ray microscope of resolution $\leq 5 \mu\text{m}$, fitted with a dispersing grating, similar to the arrangement described by Marshall *et al.*² The KB microscope mirrors are coated with iridium rather than gold.¹⁰ The dispersion of the grating was 0.97 Å/mm (insufficient for resolving spectral lines), and the image magnification was 12.5.

Figure 72.7 shows part of the spectrum (in the axial view) containing the main features of interest in this work: the $K\alpha$ line and the K -edge absorption. The absorption of core emission seen above the K -edge leads to $K\alpha$ excitation. By quanti-

tatively relating the two we can verify that the absorption of core radiation accounts for all the energy in the $K\alpha$ line. The rapid changes in absorption above the K -edge are related to EXAFS.⁴ Also, the narrow spike of absorption at the K -edge has been observed in synchrotron experiments on titanium¹⁴ and other metals¹⁵ and is attributed to a high density of d -orbital states above the Fermi level. The Ti He- α line (along with its dielectronic satellites) is a signature of mixing.⁴ Earlier experiments with thinner, Ti-doped shells⁴ showed absorption lines due to $1s$ - $2p$ transitions in Ti ions with various L -shell vacancies (in addition to K -edge absorption); the absence of these lines in Fig. 72.7 indicates that the doped layer is uniformly cold ($T \approx 100$ eV).

Figure 72.7

Detailed view of the significant spectral features in the spectrum. The Ti $K\alpha$ line is emitted by the cold titanium-doped layer due to excitation by the core radiation. Above the K -edge the continuum level is seen to drop because of absorption by the same layer emitting the $K\alpha$ line. The Ti He- α line (along with its dielectronic satellites) is a signature of mixing.⁴

1. Temperature and Areal Density Determinations

The drop in intensity above the K -edge of cold titanium at 4.966 keV enables us to estimate the temperature of the absorbing titanium layer. To that end we show in Fig. 72.8 the change in the Ti K -edge energy when successive outer electrons are removed.¹⁶ When the 12 M -shell electrons in Ti are removed, the edge shifts progressively from 4.966 keV up to ~ 5.3 keV; when the eight L -shell electrons are removed, the edge shifts progressively from ~ 5.3 to ~ 6.2 keV. Thus, a distribution of charge states would cause the intensity above the edge to decline gradually, rather than abruptly. The solid-state features near the K -edge (which may mask such a decline) extend over ~ 150 eV; we can therefore state that up to ~ 5 electrons could have been ionized. Using a collisional-

radiative steady-state calculation this can be shown to correspond to a temperature in the doped layer of less than ~ 40 eV. This is consistent with the above-mentioned absence of absorption lines, which would have required³ a higher temperature ($T \sim 200$ to 500 eV).

We now derive the areal density of the titanium shell using the absorption above the K -edge. The solid-state features near the K -edge obscure the drop in intensity at the edge. Consequently, in fitting the experimental spectrum with standard cold opacity data (which do not include these K -edge features),

we rely on a wide range of the spectrum away from the K -edge. We show in Fig. 72.9 the total measured spectrum, in axial view, after subtracting the background due to coronal emission (see Fig. 72.10). Starting with the observed spectrum $I(E)$ we derive the spectrum $I_c(E)$ emitted by the core before being attenuated by the cold titanium layer. This is given simply by

$$I_c(E) = I(E)/\exp[-\rho\Delta r \tau(E)], \quad (1)$$

where $\rho\Delta r$ is the areal density of the titanium alone and $\tau(E)$ is the absorption (or opacity) of cold titanium. We adjust $\rho\Delta r$

Figure 72.8

The shift of the Ti K -edge and Ti $K\alpha$ energies with the removal of successive electrons by ionization.

Figure 72.9

Axial view of the calibrated spectrum showing the core continuum emission (after subtracting the background due to coronal emission). The areal density of the titanium-doped layer is obtained by dividing the observed spectrum by the opacity of cold titanium: a layer of $\rho\Delta r = (16 \pm 1.6)$ mg/cm² results in an exponential curve that joins smoothly the spectrum below the K -edge. The slope yields the core temperature.⁴

Figure 72.10

The one-dimensional image of target emission, in the target plane, averaged over a photon-energy interval of 0.1 keV on each side of the $K\alpha$ line. The core emission is clearly distinct from that of the laser-heated coronal region. Note that the $100\text{-}\mu\text{m}$ -wide slit has in this image a magnification larger by a factor $(1+1/M) \sim 1.55$ than the target magnification M . The emission profile at the wavelength of the $K\alpha$ line is also shown.

until the resulting core-emitted spectrum smoothly joins the spectrum below the K -edge. The best fit is shown in Fig. 72.9 by the curve marked "corrected for Ti absorption." The corresponding areal density of the titanium alone is (0.56 ± 0.05) mg/cm²; hence, the areal density of the doped layer is (16 ± 1.6) mg/cm². The error in the areal density is determined by the noise in the spectral intensity. Using this value we can estimate the total areal density of the compressed shell, using the value of total ablated target mass. Charge-collector data showed that a layer of thickness ~ 11 μ m (of the initial target overcoat) has been ablated away, leaving ~ 23 μ m unablated. It has been shown⁴ that the $\rho\Delta r$ ratio of the unablated shell to that of an embedded layer within the shell increases only slightly during the implosion (more so, the closer the embedded layer is to the inner surface of the shell); for our case this ratio increases by up to $\sim 30\%$ for increasing compression ratios. Furthermore, the ratio can be approximately calculated by assuming the shell to implode incompressibly (the increase in areal density is mostly due to the radial convergence, not the compression); thus, a lower limit of this ratio is given by its value in the initial target. For the present target this value is 2.6, leading to an estimate of the total shell areal density of ≥ 42 mg/cm². The slope of $I_c(E)$ indicates a core electron temperature of ~ 1.1 keV.

2. The Relationship between $K\alpha$ Emission and K -Edge Absorption

The absorption of core radiation seen above the K -edge can be used to calculate the energy in the $K\alpha$ line, which can then be compared with its measured value. The relationship between the two is given by the equation

$$\varepsilon(K\alpha) = 0.5 \exp[-\tau(CH)] E(K\alpha) \omega_K(Ti) \int_{E_K}^{\infty} [I(\text{measured}) - I(\text{corrected})] dE/E. \quad (2)$$

The integral calculates the total number of photons absorbed above the K -edge and is obtained directly from the difference between the two curves in Fig. 72.9. ω_K is the relative probability for an ion with a K -shell vacancy to decay by emitting a $K\alpha$ photon (the fluorescence yield). For neutral Ti the measured fluorescence yield is $\omega_K = 0.21$.¹⁷ When the outer electrons are successively removed, the fluorescence yield increases gradually.¹⁸ As discussed above, the titanium atoms in the cold shell could have up to five electrons ionized. We calculated ω_K for these ionic species using the model described by McGuire.¹⁹ by removing five M -shell electrons, ω_K increases gradually from 0.21 to 0.31. Since the number (five) of electrons ionized

is an upper limit, we use the value for neutral Ti (namely, 0.21) and obtain a lower bound on the expected $K\alpha$ energy. Figure 72.8 shows that the $K\alpha$ line emitted by these ions would shift by less than ~ 5 eV. The observed $K\alpha$ linewidth is in fact much larger than 5 eV (see below) and thus possibly includes the emission from all these ions. $E(K\alpha)$ in Eq. (2) is the photon energy of the $K\alpha$ line (~ 4.5 keV), and the factor 0.5 accounts for the fact that only half the $K\alpha$ photons travel outward. The correction factor in front of the integral accounts for absorption of $K\alpha$ photons in the polymer shell (the $K\alpha$ line cannot be resonantly absorbed since if the temperature is high enough to produce the required L -shell vacancy, it will surely be high enough to deplete the M shell). Using the above estimated areal density of the compressed shell, we obtain $\exp[-\tau(CH)] \sim 0.7$. Applying Eq. (2) to Fig. 72.9, we obtain for the predicted $K\alpha$ energy $\varepsilon(K\alpha) \sim (15 \pm 3)$ mJ. The error results from the noise in the spectrum and uncertainty in $\tau(CH)$. Measuring $\varepsilon(K\alpha)$ directly from Fig. 72.9 (integrating over the line profile), we obtain $\varepsilon(K\alpha) \sim (13 \pm 2)$ mJ. Since the measured value is lower than the predicted value (the latter being an underestimate), we conclude that the $K\alpha$ line is indeed excited primarily by core radiation. If, for example, coronal excitation played a significant role in exciting the $K\alpha$ line, the measured $\varepsilon(K\alpha)$ would exceed the value predicted by Eq. (2). This point is further supported below. The low measured value of $\varepsilon(K\alpha)$ also suggests that the shell temperature is well below 40 eV; otherwise, the increase of ω_K with ionization would increase the value of $\varepsilon(K\alpha)$ as determined by Eq. (2).

3. One-Dimensional Continuum and $K\alpha$ Target Images

One-dimensional images of the target emission in the range of ~ 4 to 9 keV were obtained by placing a slit in front of the crystal spectrometer. Because of the relatively large width of the slit (100 μ m), the true core images will be derived by differentiating the measured slit images. Since the core emission is mostly a continuum, we select a narrow band of the continuum, close to the energy of the $K\alpha$ line, to characterize the core size. Figure 72.10 shows the calibrated one-dimensional image of target emission, averaged over a photon-energy interval of 0.2 keV (0.1 keV on each side of the $K\alpha$ line). The image is shown in the target plane, where the 100- μ m-wide slit is magnified by a factor $(1+1/M) \sim 1.55$ in terms of the target magnification M . The core emission is clearly distinct from that of the laser-heated coronal region. Note that because the slit is wider than the core (see below) but narrower than the target, the coronal intensity is per slit width, whereas the central emission is integrated over the core volume. The emission profile at the wavelength of the $K\alpha$ line (averaged over a 10-eV interval around the line center) is also

shown and is discussed later. Figure 72.11 shows the net core emission after subtraction of the coronal emission from the continuum profile in Fig. 72.10. To remove the effect of the slit, the derivative of the slit image is taken near the two slit edges (Schlieren). Figure 72.11 shows the result of the differentiation; as seen, the slit is wide enough with respect to the core size for the two images to be completely separated. In Fig. 72.11, neither the slit image nor its derivative has been smoothed (except for the averaging of the former over a 0.2-keV spectral interval). The full-widths at half-maximum (FWHM) of the two resulting images in Fig. 72.11 are marked in the figure; the two values, 46 μm and 44 μm , agree closely. We further compare the core emission profile obtained in Fig. 72.11 with that obtained by the dispersed images of the 5- μm -resolution KB microscope.² Although the shape of the time-integrated emission in the zeroth-order image indicates a nonuniform

implosion, this is nearly eliminated by the convolution of space and spectrum in the diffracted emission. It is possible, however, to estimate the size of the core emission region by taking a lineout through the diffracted emission. In Fig. 72.12 we show one section of the core profile derived in Fig. 72.11 (which was obtained from an image averaged around the $K\alpha$ line) and a lineout at 4.3 keV through the dispersed KB-microscope image. The difference in units is immaterial for this comparison. The average FWHM of the core size from Figs. 72.11 and 72.12 is thus $\sim 45 \mu\text{m}$.

We now study the spatial profile of the $K\alpha$ line. Figure 72.13 shows the spectrum emitted at various distances from the axial view, after subtraction of the background due to the coronal emission (for clarity, the lowest spectrum was shifted down from its zero-level position). Continuous core emission, as

Figure 72.11

The continuum core emission profile (from Fig. 72.10) after subtraction of the coronal emission. The derivative of the image near the two slit edges (Schlieren) yields the core emission profile. The FWHM's of the two images are marked in the figure.

Figure 72.12

Comparison of the continuum core emission profiles obtained by two instruments: (a) KB microscope images dispersed by a transmission grating, lineout at 4.3 keV, and (b) crystal spectrograph with a slit, averaged over a 0.2-keV interval around the $K\alpha$ line (from Fig. 72.11). The difference in units is immaterial for the comparison.

well as absorption and He- α lines, is observed only close to the axial view. On the other hand, the $K\alpha$ line, emitted by the cold titanium-doped layer, is seen to come from a much wider area around the hot core. Since the local $K\alpha$ emission depends only on the cold shell density (not its temperature), its two-dimensional image can reflect directly the shell uniformity. The He- α line is indicative of mixing,⁴ which is outside the scope of this work. Suffice it to say that the spatial profile of this line as seen in Fig. 72.13 clearly shows that the mixing occurs close to the core, i.e., during the shell deceleration rather than during the acceleration phase. To obtain the net $K\alpha$ spatial profile we could use the difference between the two curves in Fig. 72.10. However, a more reliable method is to examine a large number of spectra like those in Fig. 72.13 and for each to integrate over the spectral profile of the $K\alpha$ line to obtain its total energy. The results of the latter method are shown as points in Fig. 72.14 and compared with the core emission profile as obtained from

the continuum radiation. Clearly, the $K\alpha$ -emitting cold shell is considerably larger than the core; it has a FWHM of $\sim 250 \pm 20 \mu\text{m}$. The bars show typical errors in determining the $K\alpha$ energy from the spectrum. However, the scatter of the points is greater than the error bars, indicating significant shell nonuniformity, and even asymmetry with respect to the center. Because of these limitations we can use only the $K\alpha$ profile to determine the average outer extent of the cold shell. Studying the inner extent of the cold shell would require it to be highly spherical and uniform. In that case, the Abel inversion would be applied (after deconvolution of the slit broadening); a flat-top profile would correspond to a hollow shell of radial emission.²⁰ Alternatively, a target where the doping includes the mandrel layer can be chosen (the part that forms the hot core does not emit $K\alpha$ radiation). In that case, the difference in width between the two curves in Fig. 72.14 would give directly the thickness of the cold doped layer.

Figure 72.13

The spectra emitted at various distances from the axial view (after subtraction of the background due to coronal emission). Continuous core emission, as well as absorption and He- α lines, is emitted only close to the axial view. However, the $K\alpha$ line, emitted by the cold titanium-doped layer, is seen to come from a much wider area around the hot core. For clarity, the lowest spectrum has been shifted down from the zero-level position.

Figure 72.14

The spatial profile of the energy in the Ti $K\alpha$ line as measured by the spectrograph slit. The data points were obtained from lineouts such as in Fig. 72.13 and were fitted by a Gaussian curve. The spatial profile of the core emission (from Fig. 72.11) is shown for comparison. The cold, $K\alpha$ -emitting region is seen to be considerably wider than the core.

The outer extent of the cold shell is derived from the measured profile in Fig. 72.14, after unfolding the broadening by the slit. Figure 72.14 shows that the data points can be approximately represented by a Gaussian function, which we use merely to facilitate the deconvolution. The net profile of target emission $P_T(x)$ can be obtained from the measured profile $P_M(x)$ through the relationship

$$dP_M(x)/dx = P_T(x+d) - P_T(x-d), \quad (3)$$

where the slit width is $2d$. To solve this equation one starts at one end of the image where $P_T(x-d)$ is negligible and $P_T(x+d) \sim dP_M(x)/dx$, then moves successively through the rest of the image. The resulting FWHM of the deconvoluted image is found to be $\sim 235 \pm 20 \mu\text{m}$. The fitted curve in Fig. 72.14 is required merely for obtaining a smooth derivative in Eq. (3), but we could simply estimate a FWHM based on the experimental points and then subtract the slit width in quadrature, namely $(250^2 - 100^2)^{1/2} \sim 230 \mu\text{m}$; the two results are seen to be very similar. The diameter calculated here relates to the doped layer, including the cold mandrel it encompasses, but excluding the remainder of the compressed shell that is outside the doped layer (these three layers have comparable thicknesses).

It should be noted that the radial extent of $K\alpha$ emission in this experiment is not limited by radiation-pump depletion. This is evident in Fig. 72.9, where K -shell absorption reduces the pump intensity by less than $\sim 30\%$. Also, in the case of pump depletion the curve corrected for titanium absorption would not be a smooth exponential curve. The centrally peaked $K\alpha$ spatial profile provides an additional indication that it is excited by core radiation: a coronal source would produce a profile that decreases toward the center.

We can now obtain a rough estimate of the density of the compressed shell as follows: subtracting the core size ($45 \mu\text{m}$) from the diameter of the doped layer ($230 \mu\text{m}$) we obtain $92 \mu\text{m}$ for the thickness of the combined mandrel and doped layers. We deduced above the $\rho\Delta r$ of the total compressed shell from that of the doped layer. Similarly, we estimate here the $\rho\Delta r$ of the combined mandrel and doped layer, $\rho\Delta r(md)$, from that of the doped layer alone, $\rho\Delta r(d)$, found above to be 16 mg/cm^2 . Using the ratio $\rho\Delta r(md)/\rho\Delta r(d)$ in the initial target, namely, 1.7, we find for the compressed shell $\rho\Delta r(md) = 27 \text{ mg/cm}^2$. Coupled with the $92\text{-}\mu\text{m}$ thickness found above for the combined mandrel and doped layer we finally obtain $\rho \sim 2.9 \text{ g/cm}^3$ for the average density of the compressed shell, or a compression by a factor of ~ 2.7 . Because of the approxima-

tions involved, this estimate of the shell density has a large uncertainty ($\sim 50\%$). This uncertainty is large because the $K\alpha$ image is time integrated; thus it is unknown to what extent the sloping curve is due to a density gradient or to smearing in time. In a recent experiment,⁴ a pure titanium layer embedded within a thinner shell ($20 \mu\text{m}$) was shown (using an EXAFS spectrum) to compress by about a factor of 1.4 ± 0.15 . As explained above, an improved estimate can be obtained by extending the doping to the inner target surface; in that case, the inner radius of the cold shell would simply be given by the measured radius of the hot core. The relatively low shell compression measured here is the result of the low laser-irradiation uniformity; as mentioned above, distributed phase plates¹³ were not used in this experiment.

The relatively large emission region of the $K\alpha$ line should be reflected in a broader spectral profile (i.e., in the dispersion direction). Indeed, Fig. 72.7 shows that the $K\alpha$ line has a larger linewidth ($\sim 16 \text{ eV}$) than that of the $\text{He-}\alpha$ line ($\sim 8 \text{ eV}$) mainly because the latter is emitted in the hot core, whereas the former is emitted in the cold shell surrounding the core; the possible broadening of the $K\alpha$ line because of the shift with successive ionizations was shown above to be less than $\sim 5 \text{ eV}$. The source-size broadening for these lines is given by $\Delta E(\text{eV}) \sim 0.037 D (\mu\text{m})$ in terms of the emission region size D . However, this size cannot be reliably derived from this relation because of the additional broadening, which must be subtracted out before applying this relationship (this includes broadening due to the crystal, $\sim 2.5 \text{ eV}$,²¹ and the natural linewidth, which for the $K\alpha$ is $\sim 4.5 \text{ eV}$).²²

4. Core Density Derived from the Continuum Intensity

An additional diagnostic is provided by the absolute intensity of the continuum radiation from the core. Knowing the core temperature and the volume and duration of the core emission, we can derive the compressed-core density and compare it with the shell density derived above. The absolute intensity can be reliably determined since the crystal calibration agrees well with theoretical calculations.⁴ The continuum core emission is due to bremsstrahlung (free-free) and recombination (free-bound) radiations of the CH core. At the core temperature ($\sim 1 \text{ keV}$) the hydrogen and carbon ions are fully ionized and the emission at an energy E into the total solid angle, per unit volume, time, and photon-energy interval, can be shown to be²³

$$\epsilon(10^{16} \text{ keV/keV ns } \mu\text{m}^3) = 1.7 \times 10^{-4} \exp(-E/T) [\rho^2/T^{1/2}] [1 + (2E_i/T) \exp(E_i/T)]. \quad (4)$$

Here, E_i is the ionization energy of C^{+5} (0.49 keV) and T is the temperature in keV. The second term in brackets is the ratio of free-bound to free-free radiation intensities. We choose in Eq. (4) $E = 4.5$ keV (i.e., the continuum underlying the $K\alpha$ line). The Gaunt correction factors are very close to 1 when $E/E_i \gg 1$ (here $E/E_i = 9.2$). The imaging streak camera showed that the time history of core emission has a FWHM of ~ 0.25 ns. With the measured core temperature of 1.1 keV, the time-integrated emission ε_i is

$$\varepsilon_i (10^{16} \text{ keV/keV } \mu\text{m}^3) = 1.6 \times 10^{-6} \rho^2. \quad (5)$$

The measured time- and space-integrated emission at 4.5 keV is $\sim 1 \times 10^{16}$ keV/keV, after subtracting the coronal emission (see Fig. 72.9). This value has to be corrected for the attenuation through the cold shell. From the shell areal density measured above we find this attenuation to be by a factor of ~ 2.5 ; thus, the net emission is $\sim 2.5 \times 10^{16}$ keV/keV. The profile of the one-dimensional image (Fig. 72.12) can be approximated by a Gaussian of FWHM of $\Delta = 45 \mu\text{m}$. Assuming the core emission is isotropic, the peak emission per unit volume can be shown to be related to the volume-integrated emission through

$$\text{peak emission} = 8(\ln 2/\pi)^{3/2} (\text{total emission}/\Delta^3). \quad (6)$$

Combining Eqs. (5) and (6) we find a core density of $\sim 4 \pm 1 \text{ g/cm}^3$. The areal density corresponding to the FWHM of the continuum spatial profile is $\sim 17 \pm 4 \text{ mg/cm}^2$.

5. Summary of Derived Shot Parameters

Table 72.I summarizes the diagnostic results for shot 9095 analyzed in this work, and the methods used in their derivation. For comparison, the table also lists the corresponding results for shot 8207, which was the subject of analysis described in Ref. 4. The two shots are different in several important respects but yield comparable results: shot 8207 had a thinner shell ($20 \mu\text{m}$ as compared with $33.5 \mu\text{m}$) and a lower laser energy (20 kJ as compared with 30 kJ). Also, shot 8207 had a thin layer ($0.3 \mu\text{m}$) of pure titanium as compared with a thick layer ($8.7 \mu\text{m}$) of 0.5%-doped CH. The $K\alpha$ analysis in this work was not applied to shot 8207 because the $K\alpha$ line was outside the measured spectral range (see Fig. 4 in Ref. 4). Also, the EXAFS analysis of Ref. 4 was not applied to shot 9095 in this work because EXAFS data for a doped polymer (rather than a pure metal) is not known. The results given in Table 72.I show that spectral data can provide very useful and wide-ranging information of target behavior.

Conclusions

In summary, a diagnostic method is proposed for imaging the cold, compressed shell of imploding targets using high- Z doping and imaging the $K\alpha$ fluorescence. Two-dimensional

Table 72.I: Summary of diagnostic results and methods used for shots 8207 and 9095.

	Shot 8207	Shot 9095
Core temperature	0.94 keV Continuum slope	1.1 keV Continuum slope
Shell $\rho\Delta r$	$>17 \text{ mg/cm}^2$ K -edge absorption, absorption lines	$>42 \text{ mg/cm}^2$ K -edge absorption
Shell compression ratio (ρ/ρ_0)	1.4 ± 0.15 EXAFS	2.7 ± 1.3 $K\alpha$ spatial profile
Doped-shell temperature range	$\sim 10 \text{ eV}$ to $\sim 300 \text{ eV}$ K -edge absorption, absorption lines	$<40 \text{ eV}$ K -edge absorption
Core density Core areal density	$7 \pm 2 \text{ g/cm}^3$ $32 \pm 8 \text{ mg/cm}^2$ Continuum intensity and core size	$4 \pm 1 \text{ g/cm}^3$ $17 \pm 4 \text{ mg/cm}^2$ Continuum intensity and core size
Mixing	He- α line emission	He- α line emission

imaging and time gating would greatly enhance the usefulness of the method. However, even one-dimensional, time-integrated images can yield the average outer extent of the cold shell around the time of peak compression. An advantage of one-dimensional, polychromatic imaging is that the spectrum of the $K\alpha$ line (including the underlying continuum) can be measured with high spectral resolution and the net $K\alpha$ emission can be reliably determined. The present results show that adequate $K\alpha$ intensity can be obtained with a low level of doping, which minimizes target-behavior modification. The $K\alpha$ emission signature studied in this work should prove very useful in tracking the improved target performance with improved laser uniformity. Additionally, the measured absolute intensity of continuum x rays from the core can be reliably used to estimate its density and areal density.

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Near-Forward Stimulated Brillouin Scattering

In stimulated Brillouin scattering¹ (SBS) an incident (pump) light wave is scattered by the electron-density fluctuations associated with an ion-acoustic (sound) wave. This process generates a frequency-upshifted (anti-Stokes) light wave and a frequency-downshifted (Stokes) light wave. For most scattering angles the anti-Stokes wave is driven nonresonantly and can be neglected *a priori*. The resulting instability is referred to as three-wave SBS. However, for near-forward scattering the anti-Stokes wave can be driven near-resonantly and must be retained in the instability analysis. When the role of the anti-Stokes wave is significant, the resulting instability is referred to as four-wave SBS.

In this article the spatiotemporal evolution of near-forward SBS is studied in detail. The conditions under which three- and four-wave SBS occur are quantified, and expressions for the saturation times and steady-state gain exponents are derived for both types of instability.

Governing Equations

SBS is governed by the Maxwell wave equation

$$(\partial_t^2 + \omega_e^2 - c^2 \nabla^2) \mathbf{A}_h = -\omega_e^2 n_l \mathbf{A}_h, \quad (1)$$

together with the sound wave equation

$$(\partial_t^2 - c_s^2 \nabla^2) n_l = \frac{1}{2} c_s^2 \nabla^2 \langle \mathbf{A}_h \cdot \mathbf{A}_h \rangle. \quad (2)$$

The electromagnetic potential $\mathbf{A}_h = (v_h/c_s)(m_e/m_i)^{1/2}$ is the high-frequency electron velocity divided by a characteristic speed that is of the order of the electron thermal speed, n_l is the low-frequency electron-density variation divided by the equilibrium electron density, ω_e is the electron plasma frequency, c_s is the sound speed, m_e and m_i are the electron and ion masses, respectively, and the brackets $\langle \rangle$ signify that only low-frequency terms are retained. The evolution of the light waves is modified by the nonlinear electron current, which is proportional to $n_l \mathbf{A}_h$, and the sound wave is driven by the pondero-

motive force associated with the light waves, which is proportional to $\nabla \langle \mathbf{A}_h \cdot \mathbf{A}_h \rangle$.

The geometry of near-forward SBS is illustrated in Fig. 72.15. The pump wave propagates in the direction of $\mathbf{k}_0$, which defines the z axis. By substituting the Ansätze

$$\begin{aligned} \mathbf{A}_h(t, \mathbf{r}) = & \hat{\mathbf{x}} \{ A_0 + A_+ \exp[i(\mathbf{k}_s \cdot \mathbf{r} - \omega_s t)] \\ & + A_- \exp[-i(\mathbf{k}_s \cdot \mathbf{r} - \omega_s t)] \} \\ & \exp[i(\mathbf{k}_0 \cdot \mathbf{r} - \omega_0 t)] + c.c., \end{aligned} \quad (3)$$

$$n_l(t, \mathbf{r}) = N \exp[i(\mathbf{k}_s \cdot \mathbf{r} - \omega_s t)] + c.c. \quad (4)$$

in Eq. (1), where $\mathbf{k}_s \cdot \mathbf{k}_0 \approx 0$ and $\omega_s = c_s k_s$, and using the envelope and paraxial approximations, one can show that

$$[\partial_z + i(k_s^2/2k_0 - \omega_s/v_0)] A_+ = -i(\omega_e^2/2\omega_0 v_0) N A_0, \quad (5)$$

$$[\partial_z - i(k_s^2/2k_0 + \omega_s/v_0)] A_-^* = i(\omega_e^2/2\omega_0 v_0) N A_0^*, \quad (6)$$

where $v_0 = c^2 k_0/\omega_0$ is the group speed of the pump wave. The time derivatives were omitted from Eqs. (5) and (6) because the time taken for the light waves to cross the plasma is much shorter than the time taken by the sound wave to

Figure 72.15
Geometry of near-forward SBS. The pump, anti-Stokes, Stokes, and sound waves are denoted by the subscripts 0, +, -, and s, respectively.

respond to the ponderomotive force.² The term $k_s^2/2k_0$ is the reduction in k_z associated with off-axis propagation and the terms $\pm\omega_s/v_0$ are the changes in k_z associated with the anti-Stokes and Stokes frequency shifts, respectively. For all but the smallest angles the latter terms can be neglected. By substituting Ansätze (3) and (4) into Eq. (2), and making the weak-coupling approximation, one can show that

$$\partial_t N = -i(\omega_s/2)(A_0^* A_+ + A_0 A_-^*). \quad (7)$$

By making the substitutions $\omega_0^{1/2} A_+ \rightarrow A_+$, $\omega_0^{1/2} A_-^* \rightarrow A_-$, $\omega_e N/\omega_s^{1/2} \rightarrow N$, and $z/v_0 \rightarrow z$, and adopting the convention that A_0 is real, one can rewrite Eqs. (5)–(7) as

$$(\partial_z \pm i\kappa)A_{\pm} = \mp i\gamma N, \quad (8)$$

$$(\partial_t + \nu)N = -i\gamma(A_+ + A_-) + S(t, x), \quad (9)$$

where

$$\kappa = c^2 k_s^2 / 2\omega_0, \quad \gamma = \omega_e \omega_s A_0 / 2(\omega_0 \omega_s)^{1/2}, \quad (10)$$

ν is a phenomenological term that accounts for the Landau damping of the sound wave, and S is a phenomenological term that maintains the electron-density fluctuations associated with the sound wave at their thermal level in the absence of instability. In addition to the instability described in Eqs. (8) and (9), there is a mirror-image instability, in which the directions of the anti-Stokes and Stokes waves are interchanged, that evolves independently.

Analysis

The Green functions (impulse responses) associated with Eqs. (8) and (9) are defined by the equations

$$(\partial_z \pm i\kappa)G_{\pm} = \mp i\gamma G, \quad (11)$$

$$(\partial_t + \nu)G = -i\gamma(G_+ + G_-) + \delta(t)\delta(z). \quad (12)$$

It follows from Eqs. (11) and (12), and the theory of characteristics, that $G_{\pm}(t, z)$ and $G(t, z)$ can only be nonzero for $t \geq 0$ and $z \geq 0$. The solutions of Eqs. (8) and (9) are

$$A_{\pm}(t, z) = \int_{-\infty}^{\infty} \int_{-\infty}^{\infty} G_{\pm}(t-t', z-z') S(t', z') dt' dz', \quad (13)$$

$$N(t, z) = \int_{-\infty}^{\infty} \int_{-\infty}^{\infty} G(t-t', z-z') S(t', z') dt' dz'. \quad (14)$$

It follows from Eqs. (13) that $G_{\pm}(t-t', z-z')$ describes the effects on the light waves at the point (t, z) of an impulse applied to the sound wave at the point (t', z') . The role of $G(t-t', z-z')$ is similar.

One can solve Eqs. (11) and (12) by using a Laplace transform in time, for which

$$\bar{G}(s, z) = \int_0^{\infty} G(t, z) \exp(-st) dt. \quad (15)$$

The transformed Green functions satisfy the equations

$$(d_z \pm i\kappa)\bar{G}_{\pm} = \mp i\gamma\bar{G}, \quad (16)$$

$$(s + \nu)\bar{G} = -i\gamma(\bar{G}_+ + \bar{G}_-) + \delta(z). \quad (17)$$

By substituting Eq. (17) in Eq. (16) one finds that

$$\frac{d}{dz} \begin{bmatrix} \bar{G}_+ \\ \bar{G}_- \end{bmatrix} = \begin{bmatrix} -\alpha & -\beta \\ \alpha & \beta \end{bmatrix} \begin{bmatrix} \bar{G}_+ \\ \bar{G}_- \end{bmatrix} + \begin{bmatrix} -\mu\delta(z) \\ \mu\delta(z) \end{bmatrix}, \quad (18)$$

where

$$\alpha = \beta + i\kappa, \quad \beta = \gamma^2/(s + \nu), \quad \mu = i\gamma/(s + \nu). \quad (19)$$

If one denotes the eigenvalues of the propagation matrix by $\pm\lambda$, where

$$\lambda = (\alpha^2 - \beta^2)^{1/2}, \quad (20)$$

one can write the solution of Eq. (18) in the form

$$\begin{bmatrix} \bar{G}_+ \\ \bar{G}_- \end{bmatrix} = \begin{bmatrix} -\rho \\ 1 \end{bmatrix} \frac{\mu \exp(\lambda z)}{1 + \rho} - \begin{bmatrix} 1 \\ -\rho \end{bmatrix} \frac{\mu \exp(-\lambda z)}{1 + \rho}, \quad (21)$$

where

$$\rho = \beta/(\lambda + \alpha). \quad (22)$$

It follows from Eqs. (17) and (21) that

$$\bar{G} = -\frac{\mu^2(1-\rho)\exp(\lambda z)}{(1+\rho)} + \frac{\mu^2(1-\rho)\exp(-\lambda z)}{(1+\rho)} + \frac{\delta(z)}{s+\nu}. \quad (23)$$

By using the fact that $(1-\rho)/(1+\rho) = (\alpha-\beta)/\lambda$, one can rewrite the Green functions in the form

$$\bar{G}_{\pm} = \mu[\mp \cosh(\lambda z) + (\alpha-\beta) \sinh(\lambda z)/\lambda], \quad (24)$$

$$\bar{G} = -2\mu^2(\alpha-\beta) \sinh(\lambda z)/\lambda + \delta(z)/(s+\nu). \quad (25)$$

Since the Green functions are even functions of λ , they have no branch points in the complex s plane. The inversion integral associated with the Laplace transform (15) is

$$G(t, z) = \frac{1}{2\pi i} \int_B \bar{G}(s, z) \exp(st) ds, \quad (26)$$

where B denotes the Bromwich contour. To the best of our knowledge, the inversion integrals associated with Eqs. (21) and (23) cannot be done exactly.

1. Three-Wave SBS

The anti-Stokes wave is driven nonresonantly when $\kappa \gg |\beta|$. In this three-wave regime $\lambda \approx \alpha$ and $\rho \approx i\beta/2\kappa$. Since $\text{Re}(\lambda) > 0$ for small s , the $\exp(\lambda z)$ terms in Eqs. (21) and (23) correspond to spatial growth, whereas the $\exp(-\lambda z)$ terms correspond to spatial decay. The growing terms all contain the factor $\exp(i\kappa z)$, which is the spatial dependence required to drive the Stokes response resonantly. The transformed Green functions

$$\bar{G}_+ \approx -\mu[\rho \exp(\alpha z) + \exp(-\alpha z)], \quad (27)$$

$$\bar{G}_- \approx \mu \exp(\alpha z), \quad (28)$$

$$\bar{G} \approx -\mu^2[\exp(\alpha z) - \exp(-\alpha z)] + \delta(z)/(s+\nu). \quad (29)$$

The second terms on the right sides of Eqs. (27) and (29) are required for small values of z but can be neglected for large values of z , as one can verify by evaluating the inversion integrals.

For large values of t one can use the method of steepest descent³ to evaluate the inversion integrals. The Stokes response

$$G_-(t, z) \sim \frac{i\gamma \exp[i\kappa z + 2\gamma(tz)^{1/2} - \nu t]}{2(\gamma\pi)^{1/2}(tz)^{1/4}}, \quad (30)$$

and the anti-Stokes and sound responses

$$G_+(t, z) \sim i(\gamma/2\kappa)(t/z)^{1/2} G_-(t, z), \quad (31)$$

$$G(t, z) \sim -i(t/z)^{1/2} G_-(t, z), \quad (32)$$

respectively. Equations (30) and (32) are the usual results for three-wave SBS, in which the anti-Stokes response is neglected *a priori*.⁴ For future reference, notice that the coefficient of the second term in the exponent of the Stokes response is real: In three-wave SBS the Stokes frequency shift equals the sound frequency. Let

$$t_3 = \gamma^2 z / \nu^2, \quad g_3 = \gamma^2 z / \nu \quad (33)$$

and

$$\tau = \kappa^2 z / \gamma^2. \quad (34)$$

The Stokes and sound responses grow until $t \approx t_3$, at which time their gain exponent is g_3 .^{4,5} Subsequently, they decay in a time that is comparable to the growth time. It is clear from Eq. (31) that as the Stokes and sound responses grow, they produce a weaker anti-Stokes response that grows along with them. At intermediate times $G_+/G_- \sim \gamma/\kappa$. The smallness parameter γ/κ , which is the ratio of the temporal growth rate of three-wave SBS to the Stokes frequency shift associated with off-axis propagation, also arises when one studies the temporal growth of near-forward SBS.⁶ When $t \sim \tau$ the amplitudes of the anti-Stokes and Stokes responses are comparable. Since the anti-Stokes process consumes phonons, the anti-Stokes response moderates the growth of SBS significantly at this time and Eqs. (30)–(32) cease to be valid. This assertion is justified mathematically in Appendix A.

2. Four-Wave SBS

The anti-Stokes wave is driven near-resonantly when $\kappa \ll |\beta|$. In this four-wave regime $\lambda \approx (2i\kappa\beta)^{1/2}$ and $\rho \approx 1 - (2i\kappa/\beta)^{1/2}$. Since $\text{Re}(\lambda) > 0$ for small s , the $\exp(\lambda z)$

terms in Eqs. (21) and (23) correspond to spatial growth, whereas the $\exp(-\lambda z)$ terms correspond to spatial decay. The transformed Green functions

$$\bar{G}_{\pm} \approx \mp \mu [\exp(\lambda z) + \exp(-\lambda z)]/2, \quad (35)$$

$$\bar{G} \approx -\mu^2(1-\rho)[\exp(\lambda z) - \exp(-\lambda z)]/2 + \delta(z)/(s+v). \quad (36)$$

The second terms on the right sides of Eqs. (35) and (36) are required for small values of z but can be neglected for large values of z , as one can verify by evaluating the inversion integrals.

For large values of t one can use the method of steepest descent to evaluate the inversion integrals. The anti-Stokes and Stokes responses

$$G_{\pm}(t, z) \sim \mp \frac{\gamma e^{i5\pi/12} \exp\left[3e^{i\pi/6}(\gamma^2 \kappa z^2 t/2)^{1/3} - vt\right]}{(12\pi)^{1/2} (\gamma^2 \kappa z^2 t/2)^{1/6}}, \quad (37)$$

respectively, and the sound response

$$G(t, z) \sim e^{-i\pi/3} (4\kappa t/\gamma z)^{1/3} G_{-}(t, z). \quad (38)$$

It is clear from Eqs. (37) that the anti-Stokes and Stokes responses are comparable. Let

$$t_4 \approx 3^{3/4} \gamma \kappa^{1/2} z / 4v^{3/2}, \quad g_4 \approx 3^{3/4} \gamma \kappa^{1/2} z / 2v^{1/2}. \quad (39)$$

The impulse responses grow until $t \approx t_4$, at which time their gain exponent is g_4 .^{5,7} Subsequently, they decay in a time that is comparable to the growth time. When the Stokes response is maximal, the temporal rate of change of its exponent is $iv/\sqrt{3}$: In four-wave SBS the Stokes frequency shift differs from the sound frequency by an amount that is proportional to the damping rate of the sound wave. It is verified in Appendix B that Eqs. (37) and (38) are valid for $t \gg \tau$.

3. Evolution of the Impulse Responses

The spatiotemporal evolution of the impulse responses depends on the parameter $\kappa v/\gamma^2$.⁵ Suppose that $\kappa v/\gamma^2 \gg 1$; then $t_3 \ll \tau \ll t_4$. It follows from the first of these inequalities that the impulse responses grow and decay according to the three-wave equations [Eqs. (30)–(32)]. Their maximal gain

exponent is g_3 and the four-wave equations [Eqs. (37) and (38)] are never relevant. Conversely, suppose that $\kappa v/\gamma^2 \ll 1$; then $\tau \ll t_4 \ll t_3$. It follows from these inequalities that the impulse responses begin to grow according to the three-wave equations, then continue to grow and decay according to the four-wave equations. Their maximal gain exponent is g_4 . Consistent with the assertion that the anti-Stokes response moderates the growth of near-forward SBS, in this parameter regime g_4 is smaller than g_3 by a factor of $(\gamma^2/\kappa v)^{1/2}$. Notice that the transition from three-wave to four-wave growth always occurs in the absence of damping.

Discussion

The anti-Stokes and Stokes waves evolve according to Eqs. (13). The properties of the Green functions were discussed in the previous section. To complete our analysis we make the common assumption that the source term is a random function with the statistical properties

$$\langle S(t', z') \rangle = 0, \quad (40)$$

$$\langle S(t', z') S^*(t'', z'') \rangle = \sigma \delta(t' - t'') \delta(z' - z''), \quad (41)$$

where $\langle \rangle$ now denotes an ensemble average.^{8–10} The source strength σ is determined by the requirement that the density fluctuations associated with the sound wave have their thermal values in the absence of instability.^{8–10} It follows from Eqs. (13) and (41) that

$$\langle |A_{\pm}(t, z)|^2 \rangle = \sigma \int_0^t \int_0^z |G_{\pm}(t-t', z-z')|^2 dt' dz'. \quad (42)$$

Because the integrands in Eqs. (42) are non-negative, the contributions to the wave intensities from each source point increase monotonically and saturate in times that are comparable to the Green-function growth times described in the **Analysis** section. The Green-function gain exponents and growth times are proportional to $z - z'$ in both the three-wave and four-wave regimes. Thus, the contributions from adjacent source points ($z' \approx z$) are small and saturate quickly, whereas the contributions from distant source points ($z' \approx 0$) are large and saturate slowly: The saturation times and asymptotic values of the integrals in Eqs. (42) are dominated by the contributions from distant source points.

Suppose that $\kappa v/\gamma^2 \gg 1$; then the saturation time $t_s \sim t_3$.¹¹ One can evaluate the integrals in Eqs. (42) by using the method

of steepest descent.³ In steady state the Stokes intensity has the form

$$\langle |A_-(z)|^2 \rangle \sim \frac{\sigma \exp(2g_3)}{4\pi^{1/2}(2g_3)^{1/2}}. \quad (43)$$

Consistent with the discussion of the previous paragraph, the gain exponent for the Stokes intensity is $2g_3$. Equation (43) is the analog for near-forward SBS of the result of Boyd *et al.*⁸ for backward SBS. Conversely, suppose that $\kappa v/\gamma^2 \ll 1$; then the saturation time $t_s \sim t_4$.¹¹ In steady state the Stokes intensity has the form

$$\langle |A_-(z)|^2 \rangle \sim \frac{(\gamma^2/\kappa v)^{1/2} \sigma \exp(2g_4)}{2^{5/2} 3^{3/4} \pi^{1/2} (2g_4)^{1/2}}. \quad (44)$$

Consistent with the discussion of the previous paragraph, the gain exponent for the Stokes intensity is $2g_4$.

The preceding analysis is based on the weak-coupling approximation, which is not valid for short times, very small scattering angles, or very high pump intensities. A general analysis of Eqs. (1) and (2), which allows the coupling to be strong or weak, is described in Appendix C. When $\kappa\omega_s/\gamma^2 \gg 1$, SBS begins to grow as a strongly coupled three-wave instability, then continues to grow as a weakly coupled three-wave instability.⁵ The preceding results describe the later growth phase. In particular, the value of $\kappa v/\gamma^2$ determines whether SBS saturates as a weakly coupled three- or four-wave instability. When $\kappa\omega_s/\gamma^2 \ll 1$, SBS begins to grow as a strongly coupled three-wave instability, then continues to grow as a strongly coupled four-wave instability and saturates as a weakly coupled four-wave instability.⁵ The preceding four-wave results describe the saturation of SBS.

The transient (spatiotemporal) phase of SBS was observed recently by Lal *et al.*¹² In their experiment SBS was initiated by optical mixing rather than density fluctuations associated with the sound wave. There is good agreement between the theoretical predictions described herein, modified to include the effects of a low-intensity probe wave, and the experimental results.

Summary

The spatiotemporal evolution of near-forward SBS was studied in detail. Two types of instability can occur. In three-wave SBS only the Stokes and sound waves interact strongly

with each other and the pump wave. However, as the Stokes and sound waves grow, they produce a weaker anti-Stokes wave that grows along with them. In four-wave SBS the anti-Stokes, Stokes, and sound waves all interact strongly with each other and the pump wave. In the weak-coupling regime the spatiotemporal evolution of SBS depends on the scattering angle through the parameter of $\kappa v/\gamma^2$, where γ is the temporal growth rate of three-wave SBS in an infinite plasma, κ is the frequency shift associated with off-axis propagation, and v is the damping rate of the sound wave. For large scattering angles ($\kappa v/\gamma^2 \gg 1$) the instability grows and saturates according to the three-wave equations. The saturation time and steady-state gain exponent are given by Eqs. (33). For small scattering angles ($\kappa v/\gamma^2 \ll 1$) the instability begins to grow according to the three-wave equations, then continues to grow and saturates according to the four-wave equations. The saturation time and steady-state gain exponent are given by Eqs. (39). Since the anti-Stokes process consumes phonons, the presence of a strong anti-Stokes wave reduces the saturation time and steady-state gain exponent significantly. The initial growth of SBS in the strong-coupling regime and the subsequent transition to the weak-coupling regime were also discussed briefly.

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Appendix A: Three-Wave SBS

In the three-wave regime the transformed Green functions can be approximated by Eqs. (27)–(29). The inverse transform of each term in these equations is tabulated,¹³ so it is not difficult to show that

$$G_+(t, z) \approx -\gamma \left\{ (\gamma/2\kappa)(t/z)^{1/2} I_1 \left[2\gamma(tz)^{1/2} \right] \exp(i\kappa z - vt) + iJ_0 \left[2\gamma(tz)^{1/2} \right] \exp(-i\kappa z - vt) \right\} H(t)H(z), \quad (A1)$$

$$G_-(t, z) \approx i\gamma I_0 \left[2\gamma(tz)^{1/2} \right] \exp(i\kappa z - vt) H(t)H(z), \quad (A2)$$

$$G(t, z) \approx \gamma(t/z)^{1/2} I_1 \left[2\gamma(tz)^{1/2} \right] \exp(i\kappa z - vt) H(t)H(z) + H(t)\delta(z)\exp(-vt). \quad (A3)$$

Pulse Shaping on the OMEGA Laser System

Current target designs for laser-driven inertially confined fusion experiments require that the laser drive pulse be tailored to the target, i.e., that the pulse have some temporal shape other than Gaussian. On the 60-beam OMEGA¹ laser system this is accomplished by placing an electro-optic modulator at the input of the system, which is driven by an electrical pulse from a shaped microstrip line. This modulator shapes a 10-ns square optical pulse from a single-mode Nd:YLF laser, which is preamplified in a regenerative amplifier (regen) and then sent through the subsequent OMEGA amplifier chains. Gain saturation in the amplifiers and the presence of the frequency-tripling cells cause a significant distortion of the input pulse at the output; however, by modeling these effects we can generally construct a microstrip line that produces a desired output in one pass.

Pulse-Shaping System

The pulse-shaping system, shown schematically in Fig. 72.16, has been described in detail in Ref. 2 and will be described only briefly here.

1. Optical System

The initial pulse is generated from a 10-ns pulse sliced from a 100-ns single-longitudinal-mode Nd:YLF laser pulse using conventional Pockels cells. This pulse is input to the dual-amplitude, fiber-coupled waveguide integrated-optic modula-

tor. In our system, one of the two modulators is used as a gate and the other for shaping. The gate and shaped waveforms are applied to the modulator radio frequency inputs, while the dc inputs are used to adjust the modulator dc offset. The gate pulse is used to suppress possible pre- and post-pulses from the shaped-pulse-generation process and can aid in the production of pulse shapes with very steep rise or fall times.

2. Electrical Waveform Generation

The electrical waveforms are generated using a system modeled after a design developed at Lawrence Livermore National Laboratory (LLNL).³ The gate pulse is generated using a 50- Ω microstrip with a single Si photoconductive (PC) switch at one end. The microstrip is charged to the half-wave voltage of the modulator. When the Si PC switch is triggered using a short optical pulse, a square pulse of twice the electrical length of the microstrip line is generated.

The shaped pulse is generated using a 50- Ω microstrip charge line with an Si PC switch at each end. One end of the line is connected to a shaped microstrip line, while the other end is connected to the modulator. The switch nearest the shaped microstrip is triggered first, using a short optical pulse. The square pulse from the charge line propagates down the shaped microstrip, generating a shaped reflected pulse. The second switch is triggered after the initial pulse has left but before the

Figure 72.16

The pulse-shaping system consists of the optical modulators, electrical-pulse-generation system, and the SBS pulse-generation system used to trigger the Si photoconductive switches.

reflected pulse has returned, allowing the reflected pulse to propagate to the modulator. Recently, the separation between the charge line and shaped line has been increased by connecting a cable between them. This allows transients caused by the switches to die out before the reflected pulse propagates back through the charge line.

The reflection coefficient required to produce the desired electrical pulse is calculated using a layer-peeling technique.³ From this technique, the width of the shaped microstrip can be calculated using the formulas of Ref. 4. The microstrip is machined to the required profile using a computer-controlled precision milling machine.

3. Trigger-Pulse Generation

The fast-rise-time pulses required to trigger the Si PC switches are provided by focusing the regen-amplified, 1- to 3-ns pulse from a Nd:YLF oscillator into a liquid cell containing carbon tetrachloride (CCl_4). This produces a backward-propagating SBS pulse with a rise time of less than 100 ps. The timing of the leading edge of the SBS pulse depends on the amplitude of the incident pulse; hence, the input pulse is actively amplitude stabilized. Using this system, an absolute timing jitter of less than 30 ps is obtained.

Pulse-Shape Calculation

Generating the input pulse required for a desired UV output pulse involves backward propagating the pulse through an accurate model of the frequency-conversion cells and the laser system. Originally this was done by using the laser-propagation code *RAINBOW*⁵ combined with the *MIXER*⁶ model of the frequency converters. The presence of the frequency converters and the system's nonuniform radial gain cause the result obtained by this procedure to be inaccurate. This occurs because backward propagating the desired output pulse through the system produces as many input pulse shapes as radial zones used in the calculation, and forward propagating the average of the computed input pulses does not reproduce the desired output pulse. Iteration was used to improve the results, but the resulting procedure was slow and unwieldy.

By recognizing that the laser provides energy gain, rather than power gain, arbitrary pulse shapes can be forward or backward propagated through the system based on a table lookup of the output of a single *RAINBOW* run. Forward propagation through the system then requires 11 table lookups, one for the laser and ten (one for each radial zone) for the frequency converters. Backward propagation still requires iteration, but since the required calculations are now simply

table lookups, quick, accurate solutions can be obtained by standard root-finding routines.

By incorporating the laser-gain tables, frequency-conversion tables, regen gain model, and the $\sin^2$ modulator transfer function into a spreadsheet, we can now rapidly compute the electrical input pulse required to produce a specified output pulse. The layer-peeling synthesis and microstrip line impedance to width formulas have also been incorporated, allowing direct generation of the Gerber plot file required by the mill used to machine the microstrip line. Thus, microstrip generation is reduced almost to a one-step process. Bandwidth limitations in the system in general require the introduction of one extra step because our output pulses typically require an electrical input pulse with a sharp cusp at the end. The bandwidth of the shaping system is insufficient to produce this cusp, so the pulse is distorted. This problem is avoided by extending the electrical pulse beyond the desired end of the pulse and rounding it off, thereby reducing the required bandwidth. This extra portion is gated off optically using the gate provided by the second modulator and, hence, does not contribute to the resulting optical pulse. Presently, the rounding process is not automated, resulting in a two-stage microstrip generation process. It should be noted that microstrips must be designed for a particular UV output energy. In general, even small deviations from the design energy will result in significant pulse-shape variations.

Early in the development of the pulse-shaping system, it was determined that the modulator should not be driven too close to full transmission since the nonlinearity of the modulator's $\sin^2$ transmission function near peak transmission increases the electrical bandwidth requirement in this region. For this reason, we limit the shaped electrical pulse to 75% of the modulator half-wave voltage. The same bandwidth effect also occurs near minimum transmission, but many of our initial pulse shapes were not particularly sensitive to the limitation.

Pulse shapes that incorporate a lower-intensity lead-in, or "foot," place additional demands on the system. Although the intensity of typical foot pulses is not generally low, the very high initial gain of the system means that the foot is formed by an input intensity where the nonlinearity of the modulator is significant. This makes the pulse shape sensitive to offset voltages. One source of offset is caused by impedance mismatches in the electrical lines of the electrical-pulse-generation system, which introduces an effective baseline offset to the electrical pulse. This offset can be corrected by adjusting the modulator's dc bias-offset voltage. The original modulator

control system, which could automatically adjust the bias offset to zero, did not allow further adjustment of the offset. We have recently added provision to set the dc bias offset and have found that it effectively allows cancellation of baseline offsets in the electrical pulse.

Another source of offset, the importance of which was only recently realized, is the finite contrast of the modulator used for shaping. This introduces a small, but for some pulse shapes, significant optical offset, which cannot be compensated for with a modulator bias adjustment. If the modulator's contrast is known, the offset can be partially compensated for during microstrip design. However, there will be no shape control in the modulator leakage region other than that provided by the gate. For this reason, for demanding pulse shapes, we find it desirable to use the high-contrast modulator, normally used for the gate, for shaping and the lower-contrast modulator as the gate. In this case, it is important to ensure that sufficient prepulse suppression is provided by other optical gates in the system.

The overall transfer function of the pulse-shaping system, laser, and frequency converters places extraordinary demands on accurately producing the input shape. Figure 72.17 shows the calculated output obtained by placing a sinusoidal variation on the input electrical waveform with an amplitude of 0.5% of the peak. Early in the pulse this results in an output error of about 12% (relative to design).

Figure 72.17

A very small sinusoidal variation (0.5% of peak) in the input electrical waveform results in a significant variation in the output pulse shape (heavy line). The effect is most severe early in the pulse. The perturbation is barely noticeable in the optical input (light line).

Pulse-Shaping Measurements

Figure 72.18 shows the measured regen output and the measured and calculated UV output from a nominal 3-ns square pulse. While the integrated energy must be adjusted by about 10%, the calculated and measured UV output shapes are in exceedingly good agreement, indicating that the models of the laser and frequency converters are highly accurate. We have found that apparent disagreements in shape between input and output measurements are invariably due to measurement problems.

While the measured output does not perfectly reproduce the design, taking the transfer function into account, we are producing input shapes within a few percent of design. The remaining errors are largely understood. Some deviations in the output are the result of dielectric variations in the stock used to produce the shaped microstrip. These can be seen in time domain reflectometer (TDR) measurements of the microstrips. We are investigating other materials to try to improve microstrip quality. Other deviations are caused by the Si PC switches and can be reduced by inserting delay between the switches and the shaped microstrip line. Still other deviations occur when the output energy does not match the design energy. Differences between measured output and output predicted from measured input result from bandwidth limitations and the finite resolution of the measurement of the initial portion of the input pulse.

The pulse-shaping system has produced a variety of shapes, some of which are shown in Fig. 72.19. Some of the data were taken prior to recent improvements in the pulse-shaping system and thus show varying levels of quality.

Figure 72.18

Measured regen output (dashed line) and measured (heavy line) and calculated (light line) UV output for beam 19 on shot 8559, a nominal 3-ns square pulse. The measured output was scaled down about 10% to show the matching of the shape.

Figure 72.19

The pulse-shaping system has produced a variety of shapes including [(a) and (b)] 1- and 3-ns flat tops, (c) 1-ns ramp to 2-ns flat, and (d) a ramped pulse with a "foot".

Summary

Although we continue to make refinements, the OMEGA pulse-shaping system can be considered fully operational. We have produced a variety of pulse shapes and have demonstrated that we can accurately model the performance of the laser, conversion crystals, and other transfer functions involved in pulse-shape generation. Although measured outputs do not yet perfectly reproduce the designs, the errors are understood and should be reduced in the near future.

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Angular-Scattering Characteristics of Ferroelectric Liquid Crystal Electro-Optical Devices Operating in the TSM and ESM Modes

Two optical effects in liquid crystals most frequently employed for light modulation applications include (1) electrically variable birefringence, which can be used to produce either phase modulation^{1,2} or, if polarizers are employed, amplitude modulation;³⁻⁵ and (2) field-induced scattering effects.⁶ Although variable birefringence devices have received much greater attention in recent literature,^{1-5,7} the inherent optical losses posed by the requirement for polarization optics make them less desirable for applications in which incident light energy is limited. One such application is in mid-infrared imaging systems that employ uncooled focal-plane-array detectors, which are comprised of a two-dimensional array of IR detectors, sensitive to radiation in the 8- to 14- μm region, that are thermally isolated from their surroundings.⁸ The detector material is selected to display a sharp change in its electrical properties (resistance, pyroelectric polarization, dielectric constant) with temperature. A frequently used example of such a material is the ferroelectric ceramic barium strontium titanate (BST). The individual detector array elements are read out sequentially by means of an electronically addressable array of readout cells connected to each detector. The detector signals are then multiplexed out of the focal plane for signal processing and display of a visible image of the thermal signature of the scene and targets. Since these focal-plane arrays operate near room temperature, only a single-stage, low-power thermoelec-

tric cooler is required to provide thermal stabilization. A simple mechanical chopper is the only moving part in systems using BST detector arrays. These systems are compact, lightweight, and highly reliable with low-power requirements since neither a complex mechanical scanner nor cryogenic cooling is required. Numerous applications in both the military and commercial sectors have been envisioned for this emerging imaging technology.

Development activities in uncooled focal-plane-array technology are currently driven by the need to achieve further reductions in size, weight, and power requirements without sacrificing performance.⁸ One issue that has a relatively large impact on systems employing BST detector arrays is the requirement for a modulation device that enables the detector to see a change in temperature between the target scene and some background reference temperature. The current technology employs a rotating mechanical chopper containing a germanium disk with a series of lenslet arrays ground into its surface at various positions to alternately defocus (diffuse) and transmit the incident IR radiation. A simplified schematic representation of the BST focal-plane-array imaging device is shown in Fig. 72.20. Without the chopper in the system, only moving objects in the scene could be detected since, in the absence of a repeatedly refreshed background signature, the

Figure 72.20

A simplified schematic representation of an uncooled BST focal-plane-array imaging device. The motor-driven germanium chopper disk contains a series of lenslet arrays ground into its surface at various positions to alternately diffuse and transmit incident IR radiation.

detector surface charge would dissipate and cause the scene contrast to reduce to near zero. Replacement of the motor-driven mechanical chopper with a low-power, solid-state modulator that could be scaled to detector size would be a highly desirable technology-development option. Liquid crystal devices suggest themselves as an excellent alternative for this application by virtue of their low power consumption, short-path-length requirements, scalability of size, and excellent transmission characteristics in many regions of the near- and mid-infrared.^{1-7,9-12} Modulation device concepts employing the cholesteric-nematic phase transition,^{9,10} dynamic scattering,¹¹ and polymer-dispersed nematic liquid crystals¹² have been demonstrated; however, the long electro-optic relaxation times observed in these devices (100 ms to several seconds) are several orders of magnitude too slow to satisfy current requirements (<1 ms). A more promising approach uses the scattering produced by rapid field-induced unwinding of the helical structure in thick (10- to 100- μm) ferroelectric liquid crystal (FLC) devices in which the helix axis is oriented parallel to the substrates. This "transient-scattering mode" (TSM) device, first reported by Yoshino and Ozaki,¹³⁻¹⁸ is shown in Fig. 72.21. The initial application of an electric field produces a highly transparent, helix-unwound state; the electric field polarity is then rapidly reversed, and the violent molecular motion that occurs as the ferroelectric LC domains align with the new field direction results in intense light scattering. The transmissive state is then restored as the helix once again

becomes completely unwound and the new direction of spontaneous polarization is established. Rise and decay times ranged from tens to hundreds of microseconds,¹⁴ depending on the magnitude of the drive voltage pulse and the cell thickness. More recently, Marshall *et al.*¹⁹ demonstrated that it was possible to extend the duration of the scattering state to periods ranging from hundreds of microseconds to milliseconds in TSM devices by using specially designed waveforms that consisted of alternating polarity dc pulses superimposed upon a high-frequency ac signal. The ability of these "extended-scattering-mode" (ESM) devices to modulate radiation in both the visible and mid-infrared regions was verified in a simple experiment using an Fourier-transform infrared (FTIR) spectrometer, in which an unoptimized ESM device displayed a 40% modulation depth for IR radiation in the 8- to 12- μm region.¹⁹

In this article, we expand on this previous work to examine the angular distribution of forward-scattered light by the field-induced unwinding of the helical structure in TSM and ESM ferroelectric LC devices. Such information on the distribution of forward-scattered light is important for the design of optical systems employing TSM or ESM devices, in that it gives an understanding of how close to the uncooled focal-plane-array detector the modulation device must be to ensure that the entire image fills the detector's clear aperture. Although it would have been most desirable to conduct angular-scattering measurements in the mid-infrared for greatest relevance to the application of interest, our FTIR spectrometer was, by design, incapable of conducting measurements at any angle other than normal incidence. Because no other mid-IR detector was available to us, we instead conducted these measurements using an existing goniometer setup employing a visible-region light source.

Angular-Scattering Distribution Measurements

Three FLC materials commercially available from Merck, Ltd.—ZLI-4139, ZLI-4003, and SCE-9—were evaluated using the goniometer test setup shown in Fig. 72.22. The beam from a 15-mW, linearly polarized, helium-neon laser was expanded and passed through a quarter-wave plate to produce circular polarization. We chose to use circularly polarized light in this work in order to observe the distribution of forward-scattered light without polarization effects. After propagating through the quarter-wave plate, the beam was collimated and allowed to pass through the FLC cell at normal incidence. A larger beam diameter (≈ 11 mm) than in our previous work (≈ 1 mm) was used to ensure that scattering was averaged over a large portion of the device's clear aperture, thus allowing a

Figure 72.21

The transient-scattering mode in ferroelectric liquid crystals. (a) With no field applied, the helix axis lies parallel to the substrates, producing a weakly scattering texture; (b) application of a dc field causes the helix to unwind, rendering the cell highly transparent; and (c) a rapid reversal of dc field polarity and intense, transient light scattering occurs as the ferroelectric LC domains align with the new field direction.

Figure 72.22

Diagram of the experimental setup used to measure angular scattering for ferroelectric LC devices operating in both the TSM and ESM modes. The integrating sphere can be rotated through an angle of $\pm 60^\circ$, about the normal to the FLC substrate surface.

more realistic evaluation of the device's behavior. The forward-scattered light was collected by an integrating sphere with a UDT PIN-10D photodiode. The integrating sphere was rotated about the normal to the cell substrate plane to determine the angular spread (θ_s) of the incident light out to 20° ; beyond this point the signal-to-noise ratios attainable with this setup were too small to allow reliable data.

The FLC cells were driven with a Stanford Research Systems DS345 digital function generator connected to a Hewlett-Packard 6827A bipolar power supply/amplifier. The amplified signal was routed to a Hewlett-Packard 54520C digital oscilloscope to display the driving waveform. The integrating sphere detector was connected to a Melles Griot 13AMP005 wideband detector amplifier, which in turn is connected to the oscilloscope to display the output signal.

The cells evaluated in this study—SCE-9 ($10\ \mu\text{m}$ and $24\ \mu\text{m}$), ZLI-4139 ($10\ \mu\text{m}$ and $22\ \mu\text{m}$), and ZLI-4003 ($24\ \mu\text{m}$)—were assembled from glass substrates bearing a $500\text{-}\text{\AA}$, transparent, conductive ITO coating with a resistivity of $100\ \Omega/\text{square}$. The substrates were spin coated with a 2-wt%

solution of nylon in formic acid, which, after being baked at 115°C for 1.5 h and then buffed, served as an alignment layer for the FLC. The coated substrates were oriented with their alignment layers parallel to each other, but with opposing rub directions, and then bonded together using Master Bond UV15-7TK1A UV curable epoxy mixed with 10-, 22-, or $24\text{-}\mu\text{m}$ glass fiber spacers to control cell thickness. The cells were then heated above the clearing temperature of the FLC material's isotropic phase and filled by capillary action. To ensure homogeneous alignment of the FLC throughout the cell, we used a Mettler FP82HT Hot Stage to cool the cell very slowly ($0.2^\circ\text{C}/\text{min}$), from the isotropic phase down to the ferroelectric S_c^* phase, while applying a 15- to 30-V, 0.1-Hz sine wave by means of a function generator. This step was most important for thicker cells ($>20\ \mu\text{m}$) and the short helical pitch materials (ZLI-4139 and ZLI-4003) to help ensure good alignment within the FLC devices. Prior to filling, the cell thickness was verified through an interference fringe counting method using a Perkin-Elmer Lambda-9 spectrophotometer.²⁰ These empty cells were shown to vary by less than 7% from the nominal spacer diameter. For consistency, all scattering response measurements were done at ambient temperature (22°C).

TSM Angular Scattering

All TSM angular-scattering measurements were taken using a simple square wave as the driving waveform (Fig. 72.23). Similarly shaped plots were obtained for all of the devices

Figure 72.23
Electro-optical response of ZLI-4003 in the transient-scattering mode. The cell path length employed was 25 μm .

tested when the photodiode output values were plotted over a 20° angular spread, as shown in Figs. 72.24(a) and 72.24(b). Since all evaluated devices displayed symmetrical scattering behavior about normal incidence, only one-half of the scattering envelope about normal incidence is shown in these and all subsequent figures for convenience. Each cell was driven at its optimum frequency and within its ideal voltage range [determined experimentally (at $\theta_s = 0$) for each material/cell combination] to produce the greatest possible modulation depth. As shown in the figures, the bulk of the forward-scattered radiation falls within 6° of the normal and drops off sharply up to approximately 5° . Very little scattering is observed past 6° , which is in direct contrast to the behavior observed by both Tagawa *et al.*²¹ and Kobayashi *et al.*,²² who reported measurable scattering of the input beam at angles of up to 50° . One probable explanation is that the FLC cells used in these studies were one order of magnitude thicker ($\approx 100 \mu\text{m}$) than the cells used in our study and thus would be expected to exhibit much stronger scattering behavior due to the increased path length. We were unable to verify this supposition because the voltage required to drive such long-pathlength TSM devices at their

Figure 72.24
(a) Angular-scattering behavior of 10- μm -path-length SCE-9 and ZLI 4139 FLC cells operated in the TSM mode. The applied voltage per unit path length for each device is shown in the figure.
(b) Angular-scattering behavior of SCE-9 (24- μm), ZLI 4139 (22- μm), and ZLI 4003 (24- μm) FLC cells in the TSM mode. The applied voltage per unit path length for each device is shown in the figure.

G4197

optimal response levels exceeded the output capabilities of our equipment. However, we can qualitatively compare scattering as a function of path length under non-optimal drive-voltage conditions in cases where the applied voltages per unit path length are of a similar magnitude. If we examine the data in Table 72.II, which shows the relationship between the magnitude of angular forward scattering on applied field, we see that for the 10- μm and 22- μm ZLI-4139 cells, both driven at nominally 6.8 V/ μm , the thicker cell shows the greater amount of forward scattering in the 0°–6° range. A similar trend can also be seen from Table 72.II for the 10- μm and 24- μm SCE-9 cells, driven at applied field strengths of 7.5 V/ μm and 7.29 V/ μm , respectively.

The strength of the spontaneous polarization of the FLC material and its helical pitch length were both shown to

enhance temporal response in our earlier work with TSM and ESM devices, which prompted us to look for a similar relationship in these experiments with respect to the intensity of forward-scattered light. In Fig. 72.24(b), the forward-scattering response of 24- μm -thick ZLI-4003, 24- μm -thick SCE-9, and 22- μm -thick ZLI-4139 cells, each driven at its own optimum frequency and voltage requirements, is compared. The SCE-9 cell, in which the FLC material has a spontaneous polarization of 33.6 nC/cm², clearly scatters more strongly than do the ZLI-4003 and ZLI-4139 cells, in which the FLC materials have spontaneous polarization values of -20.2 nC/cm² and 13.8 nC/cm², respectively. The helical pitch length of the FLC material, shown in our earlier work to have a marked effect on response times, appears to make little if any contribution to improving the intensity or angular spread of forward-scattered light. Although our previous work

Table 72.II: Angular-scattering response data obtained from FLC cells of various materials composition and path length operated in the TSM mode.

Material Path Length (μm) Applied Field (V/ μm)	SCE-9 10	SCE-9 10	SCE-9 24	SCE-9 24	ZLI-4139 10	ZLI-4139 10	ZLI-4139 22	ZLI-4139 22	ZLI-4003 24	ZLI-4003 24
Angle (°)	Detector output (V)									
0	1.516	2.312	1.969	2.562	1.562	2.141	1.453	1.719	1.719	2.328
1	1.344	2.140	1.766	2.328	1.422	1.906	1.344	1.610	1.562	2.219
2	1.109	1.719	1.429	2.008	1.125	1.531	1.094	1.312	1.219	1.891
3	0.734	1.289	0.867	1.321	0.688	1.031	0.812	0.953	0.781	1.297
4	0.367	0.555	0.469	0.750	0.359	0.562	0.406	0.500	0.406	0.625
5	0.070	0.148	0.078	0.141	0.062	0.125	0.109	0.156	0.094	0.188
6	0.031	0.023	0.078	0.086	0.078	0.000	0.312	0.312	0.062	0.000
7	0.031	0.039	0.062	0.078	0.109	0.031	0.312	0.312	0.062	0.031
8	0.031	0.039	0.062	0.078	0.094	0.031	0.312	0.312	0.062	0.031
9	0.031	0.031	0.055	0.078	0.062	0.031	0.312	0.312	0.062	0.031
10	0.023	0.023	0.055	0.062	0.062	0.031	0.312	0.312	0.062	0.016
11	0.023	0.023	0.055	0.062	0.062	0.031	0.312	0.312	0.047	0.016
12	0.023	0.023	0.031	0.055	0.062	0.031	0.312	0.312	0.047	0.000
13	0.023	0.023	0.031	0.055	0.062	0.031	0.312	0.312	0.047	0.000
14	0.023	0.023	0.023	0.047	0.062	0.031	0.312	0.312	0.047	0.000
15	0.023	0.023	0.023	0.031	0.062	0.031	0.312	0.312	0.047	0.000
16	0.023	0.023	0.016	0.031	0.031	0.031	0.312	0.312	0.047	0.000
17	0.023	0.023	0.016	0.031	0.016	0.031	0.312	0.312	0.047	0.000
18	0.023	0.023	0.008	0.016	0.016	0.031	0.312	0.312	0.031	0.000
19	0.023	0.023	0.008	0.016	0.000	0.031	0.312	0.312	0.016	0.000
20	0.016	0.023	0.008	0.016	0.000	0.031	0.312	0.312	0.016	0.000

illustrated that materials with the *shortest* helical pitch and largest spontaneous polarization were the most desirable for maximizing the speed of response, here we can see from Fig. 72.24(b) that the material with the *longest* helical pitch and the highest spontaneous polarization (SCE-9) shows the largest amount of forward scattering over the 0° – 6° range. The ZLI 4003 cell, whose FLC material has the *shortest* helical pitch and a large spontaneous polarization, follows closely behind, while the ZLI 4139 device, whose FLC material has a relatively weak spontaneous polarization and a moderate pitch length, shows the poorest performance of the group.

ESM Angular Scattering

The ESM driving waveforms in these experiments were produced using Stanford Research System's Arbitrary Waveform Generator (AWG) software running on a standard PC platform and downloaded via an RS-232 interface to the function generator. These ESM waveforms consisted of alternating-polarity dc pulses superimposed upon a high-frequency ac signal, which made it possible to modulate the incident light so as to produce nearly ideal square wave, optical output (Fig. 72.25). Several specially designed ESM waveforms from our previous studies were examined for their ability to modulate the incident light; those waveforms showing the best performance were used in this study. These waveforms were further optimized at the function generator after downloading by making real-time adjustments to the amplitude and frequency of the overall waveform while driving the device under investigation.

The angular-scattering data collected from ESM measurements, shown in Figs. 72.26(a), 72.26(b), and Table 72.III, show some significant variations from the TSM data in Figs. 72.24(a), 72.24(b), and Table 72.II. The same general trends observed in the TSM experiments with respect to forward-scattered intensity at different applied voltages, path length, spontaneous polarization, and helical pitch length were also observed in the ESM experiments. However, comparison of the TSM and ESM data for ZLI-4139 at $10\ \mu\text{m}$ and $10\ \text{V}/\mu\text{m}$ (Fig. 72.27) shows that the TSM mode produces a considerably higher degree of scattering than does the ESM mode within the 0° – 6° angular range. Similar results were also observed for SCE-9 at $24\ \mu\text{m}$, in the TSM and ESM modes with comparable voltages per unit path length of $7.3\ \text{V}/\mu\text{m}$ (Table 72.II) and $7.1\ \text{V}/\mu\text{m}$ (Table 72.III), respectively. This finding is somewhat contradictory to our earlier investigations with TSM and ESM devices, in which there was little difference observed in the intensity of forward-scattered radiation

between the TSM and ESM modes of operation in the same device. One possible explanation may be that in these experiments the apparatus sampled a tenfold-larger area of the cell's clear aperture than in our previous work. Cells that employed longer material path lengths and/or shorter helical pitch lengths are known to have a much stronger tendency to be multidomain and thus reduce the overall scattering response. We can see some evidence of this by looking at the results in Table 72.III for the $24\text{-}\mu\text{m}$ ZLI-4003 cell, which has the shortest helical pitch length ($\sim 3\ \mu\text{m}$) and also shows the greatest variation in forward-scattered intensity between TSM and ESM modes. Another contributing factor may be that the ESM waveforms employed to obtain these results were those used in our earlier experiments and thus may require further optimization. Although some adjustment of the waveforms is possible after downloading to the function generator, this adjustment is limited to modifying only the frequency and voltage values of the entire waveform. Separate adjustments to the dc and ac components of the waveforms may provide comparable TSM and ESM performance.

Figure 72.25

Electro-optical response of the $25\text{-}\mu\text{m}$ -thick ZLI-4003 cell in the extended-scattering mode. An expanded view of the drive waveform used for the measurement is also included.

Figure 72.26

(a) Angular-scattering behavior of 10-μm-path-length SCE-9 and ZLI 4139 FLC cells operated in the ESM mode. This data was obtained using the driving waveform shown in Fig. 72.25. (b) Angular scattering behavior of SCE-9 (24-μm), ZLI 4139 (22-μm), and ZLI 4003 (24-μm) FLC cells in the ESM mode.

Figure 72.27

Comparison of angular-scattering behavior in the TSM and ESM modes for the 10-μm-path-length ZLI 4139 cell. The same 10 V/μm applied field strength was used in both driving modes.

Table 72.III: Angular-scattering response data obtained from FLC cells of various materials composition and path length operated in the ESM mode.

Material	SCE-9	SCE-9	ZLI-4139	ZLI-4139	ZLI-4003
Path Length (μm)	10	24	10	22	24
Applied Field ($\text{V}/\mu\text{m}$)	10	7.1	10	7.4	8.5
Angle ($^\circ$)	Detector output (V)				
0	1.531	1.422	1.109	2.047	1.100
1	1.461	1.344	1.078	1.921	0.965
2	1.226	1.102	0.773	1.625	0.717
3	0.766	0.726	0.531	1.180	0.497
4	0.289	0.336	0.258	0.719	0.244
5	0.070	0.133	0.109	0.328	0.094
6	0.133	0.055	0.172	0.109	0.044
7	0.102	0.031	0.187	0.055	0.031
8	0.078	0.023	0.164	0.055	0.025
9	0.055	0.023	0.156	0.055	0.025
10	0.031	0.023	0.125	0.055	0.025
11	0.031	0.031	0.109	0.039	0.021
12	0.023	0.031	0.078	0.039	0.025
13	0.023	0.031	0.078	0.039	0.019
14	0.023	0.031	0.070	0.039	0.019
15	0.023	0.031	0.062	0.031	0.019
16	0.016	0.031	0.055	0.031	0.016
17	0.016	0.031	0.047	0.031	0.016
18	0.016	0.031	0.047	0.031	0.019
19	0.016	0.031	0.039	0.031	0.016
20	0.016	0.031	0.039	0.023	0.016

Summary

In this work, we have shown that the magnitude and distribution of forward-scattered light in FLC devices operating in the TSM and ESM modes are most greatly influenced by the device's path length, the strength of the spontaneous polarization of the particular FLC material used in the device, and the operation mode—TSM or ESM. The helical pitch length, which was shown to have a large effect on the temporal response of TSM and ESM devices in our previous work, appears to have little or no effect on the magnitude or distribution of forward-scattered light in either mode of operation. Although there appears to be little difference in the angular distribution of forward-scattered light between the TSM and ESM modes, the TSM mode appears to produce more efficient

forward scattering over the 0° – 6° measurable range of our experimental apparatus than does the ESM mode. These results indicate that additional work will be required before ESM devices can be employed as modulation devices in practical IR imaging systems employing uncooled focal-plane-array detectors. Of greatest importance is to obtain a much better understanding of the physical and optical mechanisms that are the underlying processes that govern the ESM effect. Only with a better working knowledge of these processes will it be possible to design driving waveforms and new FLC materials to better enhance angular scattering while at the same time maintaining the rapid temporal response required for mid-IR modulation applications.

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Magnetorheological Finishing of IR Materials

MRF of Optical Glasses

In magnetorheological finishing (MRF), a magnetic-field-stiffened ribbon of fluid is used to polish out a workpiece. Removal occurs through the shear stress created as the ribbon is dragged into the converging gap between the part and some reference or carrier surface. The zone of contact is restricted to a spot that is smaller than the part's diameter. The removal spot constitutes a subaperture lap that conforms perfectly to the local topography of the part. Deterministic finishing of flats, spheres, and aspheres is accomplished by mounting the part on a rotating spindle and sweeping it through the spot under computer control, such that dwell time determines the amount of material removed. The MR fluid lap is unique because (1) its compliance is adjustable through the magnetic field, (2) it carries heat and debris away from the polishing zone, (3) it does not load up, and (4) it does not lose its shape. A preliminary theory and an extensive review of the MRF process are given elsewhere.^{1,2}

1. "Standard" MR Fluid

To date most MRF work has utilized an MR fluid consisting of magnetic particles of carbonyl iron (CI) in water, with small concentrations of stabilizers added to inhibit oxidation. Although this formulation will polish glass, removal rates are accelerated with the addition of cerium oxide (CeO_2). A "standard" formulation used in many experiments at the University of Rochester's Center for Optics Manufacturing (COM) consists of (given in vol %) 36/CI, 6/ CeO_2 , 55/water, and 3/stabilizers. This high-solids-content fluid exhibits an apparent viscosity³ of $\sim 0.5 \text{ Pa}\cdot\text{s}$ outside of a magnetic field. In a magnetic field of 160 to 240 kA/m (2 to 3 kG), the apparent viscosity exceeds $\sim 10,000 \text{ Pa}\cdot\text{s}$ at a shear rate of 8/s,⁴ making it stiff enough to support loads for finishing. Figure 72.28 gives initial size histograms and median particle sizes for the solid particles.⁵ The CI (left top) @ $4.5 \mu\text{m}$ is in the middle size range for commercially available magnetic powders,⁶ whereas the CeO_2 (left bottom) @ $3.5 \mu\text{m}$ is a coarse, impure, and aggres-

Figure 72.28
Histograms showing initial particle sizes of solids in MR fluid. The SEM shows the solids after one week of use.

sive commercial polishing abrasive.⁷ The scanning electron micrograph (SEM) of the MR fluid shown on the right in Fig. 72.28 was taken after one week of use. The round CI particles are essentially unaltered from their initial median particle size, but the irregularly shaped CeO_2 particles have been broken down to smaller sizes by constant milling action of the CI.

2. MRF Machine Platforms

In MRF, the workpiece is positioned above a moving surface that supports and carries a shaped ribbon of MR fluid into and through the polishing zone. The dc field from an electromagnet, located just below the carrier surface and centered under the workpiece, stiffens the ribbon before it contacts the part. Two machine configurations are shown in Fig. 72.29. The research testbed⁸ in Fig. 72.29(a) consists of a 560-mm-diam rotating aluminum plate with a shallow trough along its rim. With this machine, it is possible to conduct screening experiments on different MR fluids and/or workpiece materials, provided that the test samples used are convex and less than 50 mm in diameter. Figure 72.29(b) shows the vertical wheel configuration used in the new prototype MRF machine.⁹ Here, the wheel rim is a segment of a 150-mm-diam sphere. Convex, flat, or concave parts up to 100 mm in diameter can be finished without changing the setup.

3. Removal Rates, Polishing Spots, and Smoothing of Optical Glasses

Removal rates are obtained by generating spots as illustrated in Fig. 72.30. With the part mounted on a stationary (nonrotating) spindle, a program is used that moves the lens or surface to be polished into the field-stiffened ribbon at a fixed

contact angle of $\sim 5^\circ$. Material removal occurs during the 5- to 10-s contact time between ribbon and part. Peak and volumetric removal rates are calculated from interferometric data obtained on a phase-shifting interferometer¹⁰ after correcting for the initial surface shape. It is important to measure removal rates for surfaces that are initially well polished out (rms surface roughness less than 5 to 10 nm). This avoids reporting erroneously high rates of removal that are seen for spots taken on ground surfaces. Usually, three spots are generated over several minutes, and the average result is reported. Data for fused silica suggest that results are good to $\pm 4\%$ – 5% .

The shape and magnitude of material removed in the polishing spot are used in two ways. Peak removal rates in $\mu\text{m}/\text{min}$ give an indication of polishing efficiency for different materials when studied under identical fixed conditions. Volumetric removal rates are necessary for determining the machine program and finishing time required to process a specific part geometry.

By fixing all process parameters, it is possible to use MRF to study removal rates and conduct finishing experiments on a variety of optical glasses. Process conditions and experimental results are given in Table 72.IV, Fig. 72.31, and Fig. 72.32. The standard MR fluid is very effective at removing mass from silicates, borosilicates, lead silicates, lanthanum borates, and phosphates. Peak removal rates as high as 10 to 12 $\mu\text{m}/\text{min}$ are seen for soft glass compositions. Rates are between 1 and 2 $\mu\text{m}/\text{min}$ for the hardest glasses. An excellent positive linear correlation between volumetric removal rate and $E^{5/4}/K_c H_k^2$ (E , Young's modulus; K_c , fracture toughness; H_k , Knoop hardness) is seen in Fig. 72.31.¹¹

Figure 72.29
(a) Horizontal-trough MRF machine;
(b) vertical-wheel MRF machine.

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Figure 72.30

Removal rates are derived from polishing spots. Interferometry is used to determine the peak and volumetric removal for a part that is positioned without spindle rotation in the MR fluid ribbon at a $\sim 5^\circ$ angle for ~ 5 s.

Figure 72.31

Volumetric removal rates for optical glasses under fixed processing conditions given in Table 72.IV, plotted against a set of mechanical properties as described in the text.

Figure 72.32

Smoothing to remove microroughness^{12,13} on glass surfaces under conditions given in Table 72.IV. Levels of ~ 1 nm rms are achieved in 10 min with the "standard" MR fluid.

Table 72.IV: Removal-rate results with "standard" MR fluid on optical glasses.

Conditions: • Plano-convex plugs, ~40-mm diam $\times$ 70-mm radius; pitch polished • Horizontal-trough machine • Trough velocity, 20 rpm (~0.5 m/s) • Work spindle fixed for spots, 70 rpm for polishing • Ribbon height/gap, 2/1 mm • Magnet current, 8 A		
Glass ^(a)	Peak r rate $\mu\text{m}/\text{min}$	Vol. r rate mm^3/min
KzF6 (S)	7.81	0.55
F7 (S)	11.85	0.91
LHG8 (H)	9.16	1.43
SF7 (S)	8.48	0.65
SF56 (S)	9.24	1.24
Q89 (K)	7.75	1.40
K7 (S)	4.87	0.53
FS (C, HA)	2.08	0.29
BK7 (S,O)	4.03	0.48
SK7 (S)	10.10	0.73
SK14 (S)	4.79	0.46
LaK10 (S)	2.42	0.45
LaK33 (S)	0.99	0.11
TaFD5 (H)	1.87	0.21

(a)(S) Shott, (H) Hoya, (K) Kigre, (C) Corning,
(HA) Heraeus Amersil, (O) Ohara.

The ability of MRF to remove mass from a surface is necessary to eliminate subsurface damage and accomplish figure correction, but is not sufficient: MRF must also smooth away surface microroughness. This is done exceedingly well for glass. Figure 72.32 shows the evolution of surface microroughness with elapsed finishing time for fused silica and six optical glasses, measured with noncontacting optical profilometry.^{12,13} Conditions are as indicated in Table 72.IV. The final rms surface microroughness, independent of glass type, is ~1 nm. The smoothing process is sensitive to the initial condition of the glass surface and the glass hardness. If the initial rms surface microroughness is less than ~30 nm, smoothing occurs in 5 to 10 min. More time is required for rougher surfaces (e.g., the result for BK7 in Fig. 72.32).

MRF of IR Materials

The standard MR fluid composition is effective for finishing optical glasses, glass-ceramics, plastics, and some non-magnetic metals.¹¹ This standard fluid does not perform well on magnetic or IR materials. In this section, we identify IR materials of interest, illustrate some of the difficulties encountered when polishing these materials with the standard MR fluid, and show how new MR fluids and modified processing conditions enable the MRF of IR materials to be accomplished.

1. IR Materials of Interest

Some physical properties of eight IR materials of interest to the COM and its collaborators are given in Table 72.V. Properties of BK7 are included for comparison. Data are rank ordered according to increasing hardness. Most of these materials are considerably softer than optical glasses. One is water soluble; two are significantly harder than any glasses. Several of the single-crystal materials exhibit orientational anisotropies in hardness. The two polycrystalline compounds ZnSe and ZnS differ in grain size by an order of magnitude. Samples of most materials listed in Table 72.V were obtained in the form of plugs, 25 to 40 mm in diameter, and subsequently processed on one side into 70-mm-radius spheres. Sapphire samples were provided as bend bars, and the CVD diamond part was in the form of a 1.1-mm-thick by 17-mm-diam disk.

2. Processing Limitations of Standard MR Fluid

Table 72.VI gives removal rates and smoothing data for several IR materials under conditions similar, but not identical, to those set forth in Table 72.IV. The modified process conditions (2 $\times$ lower trough velocity and 50% larger gap—e.g., smaller spot—between workpiece and trough) reduce the polishing efficiency for the BK7 reference part by ~3 $\times$, without affecting achievable surface finish. Under these somewhat relaxed conditions, MRF of soft IR materials is too aggressive. Initially smooth surfaces become rougher by anywhere from 2 $\times$ to 30 $\times$ after 1- μm -removal runs, and pits and/or scratches appear. Figure 72.33 shows the degradation created in a pitch-polished ZnS surface by MRF.

Conversely, the standard MR fluid has little effect on hard IR materials. The removal rate for sapphire is very small, and no effect is seen for CVD diamond. The correlation between removal rate and material mechanical properties (not shown) is not observed.

3. Improved MR Fluids and Process Conditions

Two strategies are found in the literature²⁷ for polishing IR materials conventionally. One approach uses aqueous-based

Table 72.V: IR material properties.

Material	Source	Structure ^(a)	Water Solubility g/100 g @ 20°C	Knoop Hardness $H_k^{(b)}$ (GPa)	Young's Modulus ²³ E (GPa)	Fracture Toughness K_{IC} (MPa $\cdot\sqrt{m}$)
LiF	Optovac	cubic sc	0.27 ¹⁴	0.85–0.95 ¹⁵	91	0.36 ^(c)
ZnSe	II–VI	cubic pc, ~50 to 80 μm	insoluble	1.08 ¹⁶	67 ¹⁶	0.5 ¹⁶
CaF ₂	Optovac	cubic sc	0.0017 ¹⁴	1.54–1.75 ¹⁵	110	0.33 ^(c)
AMTIR-1	Amorphous materials	glass, Ge ₃₃ As ₁₂ Se ₅₅	insoluble	1.67 ¹⁷	22 ¹⁷	0.34 ^(d)
ZnS	II–VI, Morton	cubic pc, 2 to 8 μm	insoluble	2.11 ¹⁶	75 ¹⁶	0.8 ¹⁶
MgF ₂	Optovac	tetrag. sc, c-cut	insoluble ¹⁴	2.87–4.07 ¹⁸	138	0.9 ^(c)
BK7	Schott, Ohara	glass	insoluble	5.84 ¹⁹	81 ¹⁹	0.85 ²⁰
Al ₂ O ₃	U.S. Army	hexag. sc, c-cut	insoluble	13–21 ²¹	405	2.2 ^(c)
CVD diamond	NAWC	cubic pc	insoluble	56–102 ²²	1100 ²⁴	5–6 ²⁵

(a)sc = single crystal; pc = polycrystalline

(c)Weakest plane

(b)100 gf

(d)Unpublished, measured at COM

Table 72.VI: Removal-rate and smoothing results with standard MR fluid on pitch-polished IR materials.

Conditions: • Plano-convex plugs, ~25- to 40-mm-diam $\times$ 70-mm radius • Horizontal-trough machine • Trough velocity, 10 rpm (~0.25 m/s) • Work spindle fixed for spots, 38 rpm for polishing • Ribbon height/gap, 2/1.5 mm • Magnet current, 8 A			
Material	Peak r rate ($\mu\text{m}/\text{min}$)	Volume r rate (mm^3/min)	Roughness ²⁶ initial $\rightarrow$ final 1 μm removed rms (nm)
LiF	1.04	0.041	3 $\rightarrow$ 5 (p) ^(a)
ZnSe	0.34	0.012	4 $\rightarrow$ 75 (p, scr)
CaF ₂	0.26	0.029	2 $\rightarrow$ 16 (p)
AMTIR-1	4.21	0.36	2 $\rightarrow$ 4 (scr)
ZnS	0.17	0.006	4 $\rightarrow$ 132 (p)
MgF ₂	0.17	0.017	2 $\rightarrow$ 10 (p, scr)
BK7	1.78	0.18	1 $\rightarrow$ 1
Al ₂ O ₃	0.036	0.002 ^(b)	na
CVD diamond	na	na	na

(a)p = pits; scr = scratches.

(b)Trough velocity, 20 rpm; spindle velocity, 75 rpm.

slurries containing alumina (coarse/ $0.3\ \mu\text{m}$ to fine/ $0.05\ \mu\text{m}$) on a pitch or wax lap at low (50-gf) loads.^{27,28} Another approach recommends diamonds (coarse/ $3.0\ \mu\text{m}$ to fine/ $1.0\ \mu\text{m}$) in nonaqueous (or partially aqueous) glycol-based suspensions.²⁹ It is recommended that, in order to prevent pits or scratches, removal rates should be low in the final finishing stages for soft materials.

Alumina or diamond abrasives are easily substituted for CeO_2 in the MR fluid. Table 72.VII gives the properties of some submicron-size alumina and diamond polishing agents used in this work. Two interesting materials are the

Figure 72.33

2.6- μm -deep pits on a 0.25-mm $\times$ 0.35-mm area of ZnS after removing 1 μm of material with MRF under conditions in Table 72.VI.²⁶

“nanoalumina” and “nanodiamond” products, which are in the 0.04- μm size range. Figure 72.34 gives a nomograph showing, for a given particle size (assumed monosized), how many particles would be contained in 1 ml of MR fluid prepared from a given volume percent of abrasive. The standard MR fluid has $\sim 10^9$ CeO_2 particles per milliliter. There are nearly 10^{14} nanoabrasive particles dispersed in the MR fluid when it is prepared with as little as 0.1 vol % of these materials.

a. Removal rates with alumina-based MR fluids. The MR fluid is adjusted for IR materials by substituting 0.3 μm alumina at ~ 5 vol % for CeO_2 , and by further relaxing process

Figure 72.34

Nomograph for estimating the number of abrasive particles in 1 ml of MR fluid, based upon volume percent and particle size (assumed monosized).

Table 72.VII: Polishing abrasives for IR materials.

Abrasive	Source ³⁰	Structure	Form of Supply	Particle size ^(a) (μm)	
				Primary ^(b)	Aggregate/Agglomerate
9220 nanoalumina	Norton Co.	alpha platey	19%–23% solids	0.04	0.07/—
“B” alumina	Praxair	gamma blocky	dry	0.09	0.16/—
“A” alumina	Praxair	alpha platey	dry	~ 0.3 ^(c)	na
DIANAN [®] nanodiamonds	Mark Technologies	—	dry	0.03	—/0.75

(a) Numerical average of 20 primary particles by JEOL 200 EX transmission electron microscopy (TEM).

(b) *Primary particles*: single domain, homogeneously ordered, single crystals; *aggregates*: strongly bonded assembly of two or more primary particles; *agglomerates*: weakly bonded ensemble of primary particles, aggregates, or both.³¹

(c) From vendor literature.

polishing aggressiveness with reductions to trough velocity, gap, and magnet current. As shown in Table 72.VIII, volumetric removal rates for the BK7 reference are now 22× and 60× lower than those reported, respectively, in Tables 72.VI and 72.IV. Volumetric removal rates for the IR materials (Table 72.VIII compared to Table 72.VI) have actually increased (LiF, +3×; ZnSe, +1.3×), remained approximately the same (CaF₂, ZnS), or dropped (AMTIR-1, -3×; MgF₂, -8.5×). Figure 72.35 shows that the correlation previously established between removal rates and mechanical properties for MRF of glasses with CeO₂ in the MR fluid is preserved for alumina slurries, and it is now seen to hold for IR materials as well. The advantage to these adjustments is smoothing without scratching or pitting for IR materials.

Table 72.VIII: Removal-rate results for IR materials with alumina-based MR fluids and relaxed process conditions.

Conditions:		
<ul style="list-style-type: none"> • Plano-convex plugs, ~25- to 40-mm-diam × 70-mm radius, pitch polished • Horizontal-trough machine • 0.3-μm alumina abrasive in MR fluid • Trough velocity, 5 rpm (~0.13 m/s) • Work spindle, fixed (spots) • Ribbon height/gap, 1.5 to 2/1 to 1.5 mm • Magnet current, 1 A 		
Material	Peak <i>r</i> rate (μm/min)	Volume <i>r</i> rate (mm ³ /min)
LiF	4.07	0.149
ZnSe	0.57	0.016
CaF ₂	0.40	0.024
AMTIR-1	3.68	0.127
ZnS	0.22	0.007
MgF ₂	0.21	0.002
BK7	0.35	0.008
SF7	0.51	0.013
FS	0.16	0.005

b. IR material finishing experiments with modified MR fluids. Details of finishing experiments for IR materials with alumina- and diamond-based MR fluids are given in Table 72.IX. Indicated across the top and down the left-hand side of the table are the machine processing conditions and abrasives used for each material. The middle of the table gives results of smoothing experiments conducted on ring-tool-

Figure 72.35

Volumetric removal rates for IR materials and optical glasses using an alumina-based MR fluid, plotted against mechanical properties.

generated parts,³² and results from figure-correction experiments are summarized on the right-hand side of the table.

Surface smoothing without the generation of pits or scratches is seen for all materials. Roughness in LiF, CaF₂, AMTIR-1, ZnS, and MgF₂ is reduced from ~2000 nm to 4 to 6 nm with alumina abrasives under relaxed MRF conditions. Examples of smoothing are given for AMTIR-1 and ZnS in Fig. 72.36.

The addition of as little as 0.14-vol % nanodiamonds to the standard MR fluid significantly improves removal rates for sapphire, as was previously reported.³³ Table 72.IX (bottom) shows the results of an experiment to process four bend bars (3 mm × 4 mm × 45 mm long) meant for fracture strength testing.³⁴ Work was conducted with the vertical-wheel machine, using translational motion of the part along the fluid ribbon without spindle rotation. MRF reduces p-v and rms roughness of initially ground surfaces by ~10× after 4 to 6 h. Simultaneously, because of the compliant nature of the spot and its large lateral size relative to the part width, the sharp edges of the part adjacent to the work surface are smoothed. Figure 72.37 shows the setup, a part, and part edges before and after MRF. Most edge chips, cracks, and scratches are removed.

An MR fluid with 0.14-vol % nanodiamonds can be used for room-temperature polishing of CVD diamond. Table 72.IX shows that, in 80 h of elapsed polishing time, MRF with the vertical-wheel machine reduces the p-v roughness of CVD diamond from almost 9 μm to under 1 μm. The rms roughness is reduced by 10× to under 40 nm.

Table 72.IX: Finishing results for IR materials with alumina- or diamond-based MR fluids and relaxed process conditions.

Conditions:					Smoothing		Finishing	
Material	Abrasive/ vol %	Trough Vel./ Spindle Vel. (rpm)	Ribbon ht/ gap ht (mm)	Amount Material Removed (μm)	Roughness initial→final p-v (nm^{26})	Roughness initial→final rms (nm^{26})	Fig. Error ¹⁰ initial→final (90% aperture) (μm)	Amount Material Removed (μm)
LiF	B alumina/2.7	10/53	1.5/1	3	1880→45	50→3		
ZnSe	9220 nano/5.4	20/75	2/1	20	2000→250	180→15		
CaF ₂	A alumina/5.1	15/53	2/1.5	3	650→35	12→4	0.63→0.16 0.63→0.32	0.65 1.0
AMTIR-1	B alumina/2.7	1.9/53	2/1.5	9	2370 ^(b) →30	100→4		
ZnS	A alumina/5.1	20/75	2/1.5	9	2550→56	108→6	0.63→0.16	0.3
MgF ₂	B alumina/2.7	20/75	2/1	9	2340→50	57→6	0.19→0.63	1.0
BK7	A alumina/5.1	15/53	2/1.5	6	2410→20	73→1	0.32→0.13	1.0
Al ₂ O ₃	CeO ₂ Std ^(a) /5.7	175/none	2/1	na	8650→1365	540→56		
CVD diamond	9220 nano ^(a) /2.5	175/100	2/1	11	8735→840	435→39		

(a) With 0.14-vol % DIANAN[®] nanodiamonds.(b) Initial surface, loose abrasive ground with 3- μm alumina.

Figure 72.36

(a) Smoothing of AMTIR-1 with MRF after grinding with 3- μm loose abrasives; (b) smoothing of ZnS with MRF after grinding with 2- to 4- μm fine diamond ring tool on a Rogers & Clarke GP 150.

Two types of artifacts from the MRF smoothing process require further study. One type of artifact is seen in the processing of curved surfaces on single-crystal CaF₂ and MgF₂. For small amounts of material removed (e.g., 0.5 μm) during figure-correction runs, surface figure errors are corrected well. In removing larger amounts of material (e.g., >1 μm), asymmetric features are seen to develop that may relate to orientational anisotropies in mechanical properties for these crystals.^{15,18} An example for CaF₂ is shown in Fig. 72.38. Similar behavior in MgF₂ causes the increase in figure error reported in Table 72.IX.

A second type of artifact is seen in finishing ZnSe. As shown by the smoothing data in Table 72.IX, gross roughness is reduced by MRF, but only to a level that is ~10× higher than that achieved for other crystals. Microscopic examination shows that the surface features causing this microroughness are not pits or scratches but are a manifestation of hills and valleys on the part surface whose lateral feature size is comparable to the ZnSe grain size (see Fig. 72.39).²⁶ The irregular surface texture is called "orange peel" and is related to the chemistry and mechanics of the polishing process.³⁵ More work is required to adjust the MRF process for finishing ZnSe.

Figure 72.37

MRF processing of sapphire bend bars by the translation along the ribbon. The nanodiamond doped MR fluid removes edge chips, cracks, and scratches (ws = work surface).

G4211

Figure 72.38

Evolution of one type of artifact possibly related to anisotropies in mechanical properties in single-crystal CaF_2 , when polishing out a convex surface. Surface figure error data is referenced to best-fit 70-mm radius-of-curvature sphere, 90% aperture.¹⁰

G4212

Figure 72.39

Characteristic texture called "orange peel" in the surface of a ZnSe part after MRF.

4. Aspheric Surface on AMTIR-1

An important goal of MRF is the ability to finish aspheric lens surfaces for use in IR imaging systems. This article concludes with an example of rapid aspheric surface manufacturing using the IR material AMTIR-1.

A 40-mm-diam by 8-mm-thick part was cut to a base radius of 70 mm with a series of three bronze-bonded, diamond ring tools on the Optipro® SX50.³⁶ The part was prepolished for

10 min on a special lap³⁷ to remove deep tool marks. This surface was then processed by MRF in three runs (~24 min) to remove surface microroughness and change the surface from a sphere to an asphere with a conic constant of -0.35 (ellipse). Processing conditions and results are given in Table 72.X. A nanoalumina polishing agent was used in the MR fluid. Figure 72.40 gives the required radial deviation from least-removal sphere³⁸ and stylus profiler scans³⁹ of the part surface before/after MRF processing. The final figure was $0.13 \mu\text{m}$.

Table 72.X: MRF results for AMTIR-1 ellipse.

Conditions:				
<ul style="list-style-type: none"> • Plano-convex plug, 40-mm diam $\times$ 70-mm radius • MR fluid, 2.4-vol % 9220 nanoalumina • Vertical-wheel machine • Wheel velocity, 175 rpm ($\sim 1.4\text{m/s}$) • Work spindle velocity, 55 rpm • Ribbon height/gap, 2/1 mm • Magnet current, 2.5 A 				
Process Step (time, min)	Amount Material Removed (μm)	Final Roughness p-v (nm^{26})	Final Roughness rms (nm^{26})	Final Figure ^(a) Error p-v (μm)
Grind	—	2510	194	—
Prepolish (10)	—	300	5	5.11
Run #1 (14)	4	15	2	0.79
Run #2 (7)	1	13	1.5	0.20
Run #3 (2.5)	0.5	13	1.4	0.13

^(a)Ref: least-removed sphere.

Figure 72.40
Aspheric finishing of AMTIR-1 part.

Conclusion

MRF of optical glasses is accomplished with an MR fluid containing CeO_2 to accelerate removal rates. By substituting nanoalumina or nanodiamond abrasives for CeO_2 and making changes to the MRF process conditions, it is possible to finish IR materials as well. Work will continue to understand process chemistry and the mechanics of removal for these materials to further improve the quality of finish.

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Integrated Circuit Tester Using Interferometric Imaging

A noninvasive optical diagnostic technique for testing the electrical state of an integrated circuit has been developed at LLE. The interferometric probe is capable of distinguishing between the on- and off-state of an N -channel field-effect transistor (FET) with micron-size dimensions. Continued development of advanced semiconductor devices relies on the availability of diagnostic techniques to identify and eliminate design flaws. As semiconductor integrated circuits (IC) become faster and more complex, they are starting to exceed the capabilities of the present generation of diagnostic probes.¹ The technique described here requires no external leads to probe the voltage state; thus, it is relatively noninvasive. A subband-gap optical pulse is used as the probe so we can look at both buried structures and flip-chip devices. The temporal resolution is determined by the duration of the optical pulse, which is in the picosecond regime. With appropriate optics, submicron spatial resolution can be achieved. Finally, if a pulse train from a mode-locked laser is used, the sampling rate can approach the IC's natural clock frequency.

Our interferometric imaging technique is being used to provide time-resolved diagnostics of semiconductor integrated circuits. The device under test (DUT) is placed in one arm of an interferometer and illuminated with a picosecond pulse from a subband-gap infrared laser. Even though the material is transparent, it does interact with the light probe. The laser samples local variations in the index of refraction, giving rise to an interference pattern, which we detect. The variations are caused by a number of physical phenomena including doping² in the material, heating due to the flow of electrical currents,³ and changes in carrier concentrations due to injection.⁴ These variations have both static and dynamic components. The dynamic components are associated with the normal device operation and are the most interesting. To separate the two components, the device is first imaged in a quiescent state (no external voltage applied), and then a second image is taken after the device enters a known voltage state. Differences between the two images determine where the local index of refraction has changed and by how much. Activation of the DUT electronic state is synchronized with the pulsed illumina-

tion source, and the time delay between the initiation of the state and the laser probe pulse allows us to take a series of images that map the time evolution of the interferogram. So long as the DUT state is stable, multiple probe pulses can be averaged to improve the signal to noise. The stability does not require that the state be static. For example, a given state can be initiated in the DUT and probed at a later time Δt . The DUT can then be reset to zero and reinitialized before the next laser pulse arrives at $T_{\text{clock}} + \Delta t$, where T_{clock} is a multiple of the clock period. A mode-locked laser typically operates at 50 MHz to 100 MHz, so it is possible for a device to be tested at or near the normal clock frequency.

Physical Basis of the Measurement

Semiconductors are characterized by a band structure with a nearly full valance band and a nearly empty conduction band. We build devices by modulating the number of charge carriers in these bands. This modulation affects not only the electronic properties but also the optical properties of the semiconductor. The particular optical property we are interested in is the index of refraction. An excess of free carriers ΔN causes a change in the index of n :

$$\Delta n = n \Delta N e^2 / (2 \epsilon \omega^2 m^*), \quad (1)$$

where e is the electronic charge, m^* is the effective mass of the carriers, ϵ is the dielectric constant, and ω is the angular frequency of the incident radiation.⁴ Using numbers for silicon and 1- μm radiation we get

$$\Delta n \sim 10^{-21} \text{ cm}^{-3} \Delta N. \quad (2)$$

A submicron FET can have an excess free-carrier concentration of 10^{20} to 10^{18} cm^{-3} extending to a depth l of 0.25 μm into the bulk.⁵ A light pulse passing through this region (double pass) will experience a change in optical path ($l \Delta n$) of 0.0005 to 0.05 waves. The former value is probably beyond the resolution of our system, but 1/20 wave should be observable.

Free carriers are not the only contributors to the change in the index of refraction: temperature can also change the index of refraction. Typically the index of refraction will change by one part in 10^{-4} to 10^{-5} for each degree Kelvin.³ The temperature can change because as current flows in the semiconductor, I^2R heating occurs. Two situations can develop: If the device is embedded in a silicon substrate, the high thermal conductivity will diffuse the energy over a large volume. Thus temperature rises of a few degrees can be distributed over as much as $500\text{ }\mu\text{m}$, which is the typical substrate thickness. The small change in the index of refraction will be offset by the large path length. These combine to induce changes in the optical path lengths of about $1/20$ wave. Alternatively, a buried SiO_2 layer can insulate the device from the substrate leading to 100°C temperature changes over a few wavelengths, again producing few-percent changes in the optical path length.⁶ We estimate that the resolution of our imaging system is 0.02 waves limited by our 8-bit CID camera. This is based on an analysis of standard deviation of several identical measurements. This magnitude of change is easily seen in the silicon FET circuits we have examined.

We did not have a sufficiently detailed description of the device structure to calculate the combined effects of all the physical phenomena in a real device. That calculation would require a 3-D map of all the dopant and barrier layers, a detailed knowledge of the index of refraction of each of these structures, and the dependencies on temperature and free-carrier concentration. In addition, a dynamic model that simulates device operation and converts the output into optical-index-of-refraction data is required. Rather than developing such a first-principles method of evaluating our data, we have

chosen a simpler, more empirical approach. A set of transistors was placed in the interferometer and set to a number of different voltage states. An interferogram was acquired at each state, and the differences were compared to see if the different voltage states were resolvable. This same criteria would apply if this diagnostic were fielded in a manufacturing setting, i.e., it would be used to determine if a particular transistor was working but would probably not be used to determine, for example, the exact carrier density in the gate channel of an FET.

The Tyman-Green interferometer⁷ (see Fig. 72.41), which is used to determine the index-of-refraction changes, allows the whole substrate to be examined at once since it works with a large, collimated input beam. The interferometer uses a short-pulsed laser instead of a cw beam, which allows us to take time-resolved "snap-shot" interferograms. The interferometer is constructed as follows: The beam from a pulsed Nd:YAG laser (120 ps) is up-collimated in a telescope to a diameter greater than the DUT. The laser passes through a compensated absorbing doublet, which converts the Gaussian spatial profile into a flat-top. This beam strikes a 50/50 beamsplitter. The DUT acts as a mirror in one of the two interferometer arms. The second arm is terminated in a flat mirror. Both beams then retrace their respective paths back to the beamsplitter, where they recombine, i.e., the electric fields add. The combined beam passes through a pinhole to eliminate stray reflections from optical surfaces and grating reflections from the circuit elements, then through imaging optics, and finally onto a charge-injection-device (CID), 512×512 camera where the interference pattern is detected.

Figure 72.41

In the experimental setup a pulsed laser illuminates a Tyman-Green interferometer. The recombined beam was imaged onto a CID camera with $4\text{-}\mu\text{m}$ resolution.

The interferometer mixes the electric fields from the reference, E_{ref} , and device, E_{DUT} , arms to produce a total irradiance I . Due to differences in the optical path length the sum can be written as

$$E_{\text{ref}} + E_{\text{DUT}} e^{i\Psi}, \quad (3)$$

where Ψ is a phase angle. The detected quantity is actually the total irradiance obtained by multiplying the total electric field by its complex conjugate. We note that the intensity is equal to the square of the electric field times a constant, $I = \epsilon c E^2/2$, where ϵ is the permittivity and c is the speed of light. Thus the total I irradiance can be written as⁸

$$I = I_{\text{ref}} + I_{\text{DUT}} + 2(I_{\text{ref}} I_{\text{DUT}})^{1/2} \cos \Psi, \quad (4)$$

where I_{ref} is the reference-arm irradiance and I_{DUT} is the test-arm irradiance. The irradiance formula has three physically measurable quantities that are used to determine the phase angle Ψ , I_{ref} , I_{DUT} , and I . Several images must be acquired to accomplish this determination: First, I_{DUT} and I_{ref} are measured by alternatively blocking each of the two arms of the interferometer. Next, the interferogram I is acquired with both arms unblocked while the DUT is in a zero-voltage or inactive state. This enables the determination of $\Psi(V=0)$. Finally, an image is acquired with the DUT in a known voltage state, which allows $\Psi(V)$ to be determined. The interesting physical quantity is $\Delta\Psi = \Psi(0) - \Psi(V)$ because it can be related to a length times a change in the index of refraction. The DUT may have different combination external voltages applied to it to determine $\Psi(V)$ for different states.

Optics must be of high quality to make this device work; however, the most difficult optic to deal with is the DUT, which is typically a piece of silicon with semiconductor components fabricated on it. We want to probe the interior of this material, so the laser beam must penetrate the surface. The back surface of the device is the choice for the laser illumination because the interior structures are not obscured by surface metalization. However, the back surface of an IC is not typically optically polished and not parallel with respect to the front surface. Our test chips were polished but not parallel. Silicon has a 3.7 index of refraction, leading to a 32% reflection as the IR enters the material.⁹ This reflection can combine with the interferogram and reduce the visibility of the fringes. The antiparallelism of the front and back surfaces produces beams that are not quite collinear. If the beams become sufficiently separated in the detector leg of the interferometer, a pinhole can be used to

block the stray reflection. This pinhole must be selected carefully. If the pinhole is too small, it will attenuate the high spatial frequencies in the image, thereby reducing the resolution. The back side of the chip can be antireflection coated to minimize the reflection. We have demonstrated that an antireflection coating consisting of single quarter-wavelength-thick layer of HaO_2 can reduce the intensity of the reflection, but it is not clear that the device can survive the deposition processes.

The light source was a Nd:YAG laser operated at 30 Hz and was synchronized with the scan rate of our 512×512 CID camera. The laser pulse arrived during the blanking period of the frame to prevent extraneous noise in the image. The second mirror is the front surface of an uncoated glass wedge. The intensity of the reflection roughly approximated the average reflection from the DUT. From Eq. (4) it can be seen that the maximum fringe visibility will occur when the reference and test arms are exactly balanced. It was impossible to exactly balance the arms because the DUT is a nonuniform mirror. At some points the IR light will reflect off of surface metalization with a near-unity reflectivity; at other points it will encounter the bare silicon interface with 32% reflectivity; and at some points absorption in the IC will result in no reflected light.

Even with high-quality optics, it was impossible to completely eliminate large-scale fringes in our images due to multiple reflections and misalignment. The fringes were typically confined to a few well-defined frequencies. These were removed by Fourier transforming the image and filtering in the frequency domain.

A test structure provided by Intel consisted of several NMOS FET's with varying channel sizes. There was a common gate as well as a common source and P_{well} electrodes for all of the devices, but each device had a separate drain. The transistors were situated in a line between 100- μm -square drain contact pads with the source and gate being fed in through a bus line on either side of the pads. The back of the IC package was open to allow optical access to the back side of the chip. A circuit was constructed around each of the transistors consisting of a grounded source, a 2-V gate pulse supplied through a 3-M Ω resistor, and a variable-amplitude drain voltage supplied through a 500-k Ω resistor. For these tests, a 200- μs , 0- to 3-V pulse was applied to the common gate with and without the drain voltage present. Images taken under these conditions constitute the generalized set of images I_{DUT} . These images, together with the appropriate reference images, allow us to calculate the phase angles associated with each device state.

Experimental Results

The experimental results consist of the phase angles associated with each state of the transistors. Three different NMOS FET's were examined: a $1344\text{-}\mu\text{m} \times 0.7\text{-}\mu\text{m}$ channel serpentine device, a $1.0\text{-}\mu\text{m} \times 20\text{-}\mu\text{m}$ device, and $20\text{-}\mu\text{m} \times 0.7\text{-}\mu\text{m}$ device. The first number represents the length of the source and drain electrodes; the second number is the width of the gate channel. Each device was illuminated in four different voltage states: no applied bias, gate bias only, drain bias only, and gate and drain together. The gate bias was pulsed for $100\text{ }\mu\text{s}$, and the probe came in the middle of the pulse. The imaging of the device was such that each camera pixel corresponded to $4\text{ }\mu\text{m} \times 4\text{ }\mu\text{m}$; thus, several pixels were active for each device. The active pixels encompassed the gate channel, drain, and source regions. The phase angle presented for each state is the average over all of the active pixels. In addition several images were averaged to produce each measurement. The standard deviation of the value of Ψ as computed from the several images was always less than $1/30$ wave and was typically less than $1/50$

wave. Figures 72.42 and 72.43 show the phase variation for the three devices. The $1344/0.7$ and the $1.0\text{ }\mu\text{m} \times 20\text{ }\mu\text{m}$ show similar behavior with the phase angle increasing (with respect to the no bias case) as the gate is turned on, then decreasing as the drain is turned on. The last device shows different behavior, with the case of drain voltage without a gate being the most different of all the device states. However, without an underlying physical model it is impossible to interpret these results in terms of specific carrier-density changes. What is important is that the different states are resolvable. When transistors turn on, we see variations of 0.1 to 0.15 waves giving signal-to-noise ratios of from 3 to 5. A similar analysis on portions of the device containing no transistors found no correlation with applied voltage. This technique makes it possible to determine if a particular transistor in an IC is actually functioning.

Limitations to this technique are obvious: For example, there was also a $1.0/0.7\text{-}\mu\text{m}$ device on the Intel chip. Given our imaging system, this device should only have affected about

Figure 72.42
The $1344/0.7$ and the $1.0\text{-}\mu\text{m} \times 20\text{-}\mu\text{m}$ devices showed similar behavior. The error bars are displayed next to the 0 bias case.

Figure 72.43
The $20/0.7$ device showed the largest phase variations.

one pixel. We were unable to resolve this device. When we examined pixels likely to be associated with this device, no correlation was found with any applied voltages. Also, some structures on the IC precluded the use of this technique. Typically the metalized bonding pads showed up as dark spots on the IC images. This was unexpected since the metalization should have provided a good mirror. It is probable that the polysilicon used to electrically isolate the pad acts as infrared absorber. Thus another device located under a surface metalization structure was inaccessible because no light from that area reached the camera.

Conclusions

The technique described here offers an exciting new diagnostic for semiconductor integrated circuits. The technique is noninvasive and compatible with high-speed operation of integrated circuits. The picosecond time resolution enables us to either characterize specific logic states or even watch an individual device turn on. This imaging technique is sensitive to all of the index-of-refraction changes that can be associated with IC's, including heating due to current flowing through narrow wires and charge injection into the depletion region of a transistor. As more manufacturers shift to flip-chip technology and its associated inaccessibility of surface leads for device testing, the noninvasive nature of the interferometric imaging technique could enable it to become one of the few external IC testers still capable of operating with the latest chip designs. It should also be able to operate at the IC's rated clock speed.

The interferometric method examined and tested here will provide a very effective method of IC testing. The fringe sensitivity is of the order of $1/30$ wave, which is sufficient to distinguish the voltage state of an FET in an integrated circuit. This work is in a preliminary state, and several improvements could easily be achieved by using an electronic camera with lower noise, using better optics with higher magnification, and applying antireflection coatings to reduce optical losses and eliminate stray beams. These improvements should make it possible to actually measure physical parameters such as the time-resolved carrier density in a switching transistor.

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LLE's Summer High School Research Program

During the summer of 1997, ten students from Rochester-area high schools participated in the Laboratory for Laser Energetics' Summer High School Research Program. The goal of this program is to excite a group of high school students about careers in the areas of science and technology by exposing them to research in a state-of-the-art environment. Too often, students are exposed to "research" only through classroom laboratories, which have proscribed procedures and predictable results. In LLE's program, the students experience all of the trials, tribulations, and rewards of scientific research. By participating in research in a real environment, students become more excited about careers in science and technology.

The students spend most of their time working on individual research projects with members of LLE's technical staff. The projects, which can typically be completed during the eight-week period, originate from all areas of interest in the laboratory: from plasma physics and optics to chemistry and materials science. The projects vary from year to year depending on LLE's requirements. An example of the breadth of projects is shown in Table 72.XI, where the students and their projects from 1997 are listed.

Table 72.XI: High School Students and Projects—Summer 1997.

Student	High School	Supervisor	Project
Devon J. Battaglia	Greece Arcadia	M. Guardalben	Analyzing Aberrations in a Laser System Using a Liquid Crystal Point-Diffraction Interferometer
Daniel Bouk	Greece Athena	S. Craxton	Beam Phasing for Indirect-Drive Experiments on OMEGA
Anthony E. Green	Bishop Kearney	W. Donaldson	Constructing and Characterizing a Diode Laser
Lisa Haber-Thomson	Brighton	I. Walmsley	Femtosecond Autocorrelation Measurements Based on Nonlinear Two-Photon Induced Photocurrent
Yanlin Liu	Brighton	J. Delettrez	Fourier Decomposition of Rayleigh-Taylor Unstable Surfaces
Sarah Mitchell	The Harley School	K. Marshall	Self-Organizing Lyotropic Liquid Crystals Based on Cellulose Derivatives
Matthew T. Pandina	Rush-Henrietta	J. Kelly	Lower-Level Lifetime and its Effects on Square-Pulse Distortion
David Rea	Honeoye Falls-Lima	T. Kessler	High-Efficiency Silver Halide Holographic Diffraction Gratings: Production, Evaluation, and Results
Danielle Schillinger	Our Lady of Mercy	R. Epstein	The Effect of Rayleigh-Taylor Instability of Compressible Fluid
Jeremy Yelle	Honeoye Falls-Lima	J. Knauer	The Calibration of Plasma/X-Ray Calorimeters

The students attend weekly seminars on technical topics associated with LLE's research, including an introduction to lasers in general, to LLE's 60-beam OMEGA laser, and to plasma physics and controlled fusion. The students also receive safety training and are introduced to LLE's resources, especially the computational facilities. Topics have also included introductions to nonlinear optics, instabilities, data analysis, numerical methods, and ethics.

The program culminates with the High School Student Summer Research Symposium at which the students present the results of their research. They also prepare written reports, which are bound into a permanent record of their work. These reports are available by contacting LLE.

Eighty-two high school students have participated in the program since it began in 1989. Each year students are selected from approximately 50 highly qualified applicants.

In 1997, LLE added a new component to its high school outreach activities: an annual award to an Inspirational Science Teacher. This award honors teachers who have inspired our High School Program participants in the areas of science, mathematics, and technology. The students in our program have been exceptional, and we are honoring those teachers who have brought them to this point. The award is presented at the High School Student Summer Research Symposium and includes a \$1000 cash prize. Alumni of our Summer High School Research Program nominate teachers for the award. The 1997 winner is Mr. Raymond Sherbinski from Brighton High School. Mr. Sherbinski, a mathematics teacher, inspired Daniel Nelson, a participant in our 1994 program. In nominating Mr. Sherbinski, Daniel wrote, "I left my year with Mr. Sherbinski with renewed vigor for not only math, but for all learning. Mr. Sherbinski creates a class atmosphere that removes a student's inhibitions toward risking vocal expression of ideas before his or her peers."

Laser Facility Report

During FY97 the OMEGA Laser Facility was in full operation, shooting targets for 41 of the 52 weeks. The remaining time was invested in adding new capabilities as well as performing preventive maintenance. New features added this fiscal year include a fourth-harmonic fiducial laser for timing reference, a limited set of distributed polarization rotators (DPR wedges for up to five beams), and a multibeam streak camera. Incremental improvements in smoothing by spectral dispersion (SSD), timing systems, power conditioning, and beam-balance equipment were accomplished this year to supplement OMEGA's versatility.

During the first quarter the fiducial laser was installed in the target bay. This independent laser source is synchronous with the OMEGA beamline pulses and is utilized by a variety of target and laser diagnostics as a timing reference. The deep UV of the fourth-harmonic pulse is used on a number of x-ray diagnostics to measure relative timing of target events. Flexibility in delivering pulses for diagnostic setup and troubleshooting improves the overall performance of instruments in the single-shot acquisition mode common to experiments on the OMEGA facility.

Uniformity improvements for FY97 included the deployment and testing of DPR wedges on five of the 60 OMEGA beams. Time instantaneous $\sqrt{2}/2$ improvements in single-beam uniformity, as theory predicts for DPR's, was confirmed with UV and x-ray imaging techniques. The five prototype components were fabricated from KDP material and AR coated for high UV transmission. Several of the planar-foil campaigns characterized improved beam imprinting and foil acceleration when these devices were installed.¹ LLE continues to evaluate options for cost-effective production of these devices.

The first of several multibeam streak cameras was installed on one ten-beam cluster of OMEGA. While the device is collecting data on beam-to-beam power balance, the preliminary nature of the data and startup of this complex diagnostic preclude definitive power-balance conclusions. As our experi-

ence with this instrument grows, we expect to obtain high-dynamic-range, high-temporal-resolution characterization of OMEGA UV beams.

SSD system implementation continued. Amplitude modulation at the fundamental frequency of the modulators was significantly reduced through two minor modifications to the SSD configuration. Reducing the beam spot size in the far-field crystals and limiting the beam deflection angle while imaging through front-end components allowed the bandwidth to be nearly doubled (to $1.25 \times 1.75 \text{ \AA}$) without inducing appreciable amplitude modulation (5%). In addition, a phase-locking technique was implemented to optimize smoothing and ensure repeatability of smoothing parameters.

Progress in uniformity of laser amplifier performance resulted from adding the capability to diagnose small-signal gain. Equipment and procedures were developed that allow accurate measurement of the gain of each of OMEGA's 213 amplifiers in a next-day operation. The data from these experiments demonstrates consistency within the six gain stages and clearly identifies individual units that do not lie within the normal distribution. The few units not performing at the average level for a given stage generally are found to have electrical circuit component failures, which can be readily rectified prior to proceeding with target shots.

The shot summary for OMEGA for FY97 is as follows:

Driver	722
Beamline	419
Target	<u>805</u>
Total	1946

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National Laser Users' Facility News

Management and Operations

Several important changes were made to the National Lasers Users' Facility (NLUF) management and operations during the 1997 fiscal year:

1. Proposal Submission Schedule

In coordination with the Department of Energy the proposal review schedule was advanced by three months. DOE will now issue calls for proposals sometime near the end of the calendar year, and the proposals will be reviewed in May of the following year. This will allow the principal investigators to receive information on the disposition of their proposals in sufficient time to make plans for the subsequent fiscal year.

2. Research Areas

In keeping with the NLUF Steering Committee's recommendation, DOE approved an expansion of research areas to include nearly all aspects of high-energy-density science.

3. NLUF Users' Guide

A new NLUF Users' Guide was produced and distributed. The new guide has enhanced descriptions of the OMEGA laser facility and its capabilities to assist prospective users with the development of their proposals.

4. NLUF Web Page (<http://www.lle.rochester.edu/pr/nluf.html>)

An enhanced NLUF home page was implemented that contains an on-line version of the NLUF Guide. Also available on the LLE web page for NLUF users is the annual OMEGA facility schedule. In addition, a complete diagnostic database is available via the LLE Home Page to approved users.

5. DOE Grants Program

Coordination between the NLUF and the newly formulated DOE Grants Program was established.

6. Users Questionnaire

Based on a recommendation of the NLUF Steering Committee, a users questionnaire was established to solicit input from facility users to continually improve NLUF operations.

7. FY98 Proposals

Eleven proposals, with funding requests totaling \$1,371,593, were submitted for FY98 work. This is the largest number of proposals submitted since 1991.

Progress in Current Programs

Users funded under FY96 and FY97 grants continued to make progress on their research:

Hans Griem, Ray Elton *et al.* (University of Maryland) completed their FY96 work with a series of OMEGA shots taken in September 1996 and submitted a final report. They have continued their FY97 work and have taken two series of runs on OMEGA during this year. One of the goals of their experiments has been to measure the electric field in the laser-produced plasma at early times using satellites-to-forbidden line ratios in the extreme ultraviolet region. The inferred electric fields from these measurements are consistent with estimated values of the laser irradiance. Several publications and reports were produced on these experiments.

John Seely (Naval Research Laboratory) continued to make progress on obtaining time-resolved, hard- and soft-x-ray images of laser-irradiated targets (see Fig. 72.44). Time-resolved, soft-x-ray images of targets were obtained with a microscope and compared to hard-x-ray images obtained with filtered x-ray pinhole cameras. Recently Dr. Seely's work has focused on extending these results to four different energy regimes with the same instrument. Theoretical modeling is planned to compare the images with the expected laser irradiance uniformity.

Joseph J. MacFarlane (University of Wisconsin) has begun soft-x-ray tracer diagnostic experiments on hohlraums on the Nova laser. When this experiment was originally proposed, no other hohlraum experiments had been conducted on the OMEGA laser, so the NLUF Steering Committee recommended that Dr. MacFarlane's experiments be coordinated with the national laboratories and begun on the Nova laser. A proposal from the same group for FY98 has since been submit-

Figure 72.44

Monochromatic soft-x-ray image (257-eV energy) of the emission from the focal spot of six overlapping OMEGA beams without SSD beam smoothing. The emission is uniform on a spatial resolution scale of $10\text{ }\mu\text{m}$. The image was formed by a multilayer microscope and was recorded on two striplines of a gated framing camera 2.0 and 2.5 ns after the rise of the pulse.

ted to NLUF and accepted for indirect-drive experiments on the OMEGA laser.

Charles Hooper (University of Florida) completed during the past year the experiments called for in his FY96 grant and has completed planning for the FY97 work. In experiments conducted in March 1997, a density-dependent shift of spectral lines is indicated. Time-resolved comparisons of *K*-shell Ar and *L*-shell Kr lines were completed, and diminution and subsequent recovery of broadband emission were observed (using time-resolved x-ray spectroscopy) near the stagnation

time of imploded targets. Dr. Hooper's follow-on experiments require the use of full 2-D SSD with phase plates and thus are scheduled to be conducted in FY98 when a full complement of phase plates will be available on OMEGA.

Paul Drake (University of Michigan) began experimental design work for long-scale-length plasma experiments to be conducted on OMEGA. Software is being developed to manage the data to be produced in these experiments, and instrumentation is under development to conduct Thomson-scattering experiments on long-scale-length plasmas.

Qichang Su (Illinois State University) has continued collaboration with LLE scientists on analysis of x-ray spectra from imploded capsules. The work has so far produced four papers on diagnostics for high-density implosions.

Stephen Padalino (SUNY Geneseo) carried out a pilot project to compare the neutron yield [produced in $T(d,n)^4\text{He}$ reactions] as measured by aluminum activation to that measured by copper activation.

FY98 Proposal Review

The NLUF Steering Committee met on 13 May 1997 to review proposals submitted for FY98. The Steering Committee members in attendance were Drs. Michael Cable (Lawrence Livermore National Laboratory), Allan Hauer (Los Alamos National Laboratory), Ramon Leeper (Sandia National Laboratory), Kevin McGuire (Princeton Plasma Physics Laboratory), John Soures (NLUF Manager), and Prof. Tudor Johnston (Institut National de la Recherche Scientifique). Eleven proposals were submitted to DOE for consideration as NLUF research during FY98. The committee carefully reviewed all proposals and approved for funding the eight proposals listed in Table 72.XII.

Table 72.XII: Approved FY98 NLUF Proposals.

Principal Investigator	Affiliation	Title of Proposal
H. Griem	University of Maryland	Early-Time Measurements of Laser-Plasma Conditions in OMEGA-Upgrade ICF Targets.
Q. Su	Illinois State University	Measurement of High Ion Temperature Using the Doppler Width of the Kr He β Line (0.8 Å) from Kr-doped Target Implosions.
S. Padalino	SUNY Geneseo	Neutron Yield Measurements via Activation Diagnostics.
S. Padalino	SUNY Geneseo	Measurements of Core Conditions, Implosion Symmetry, and Anomalous Acceleration Mechanisms on OMEGA through Charged-Particle Spectrometry.
R. Petrasso	Massachusetts Institute of Technology	Measurements of Core Conditions, Implosion Symmetry, and Anomalous Acceleration Mechanisms on OMEGA through Charged-Particle Spectrometry.
J. MacFarlane	University of Wisconsin, Madison	Development of X-Ray Tracer Diagnostics for Radiatively Driven Ablator Experiments.
J. Seely	Naval Research Laboratory	Beam Uniformity Determination by Multi-Thermal Soft-X-Ray Imaging.
C. Hooper	University of Florida	Time-Resolved X-Ray Spectroscopy of Imploded Gas-Filled Microballoons: Systematic, Quantitative Comparison with Nuclear Diagnostics and Investigation of Line Shifts, Continuum Lowering, and Opacity.

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T. R. Boehly, R. L. McCrory, S. J. Loucks, J. M. Soures, C. P. Verdon, A. Babushkin, R. E. Bahr, R. Boni, D. K. Bradley, D. L. Brown, R. S. Craxton, J. A. Delettretz, W. R. Donaldson, R. Epstein, P. A. Jaanimagi, S. D. Jacobs, K. Kearney, R. L. Keck, J. H. Kelly, T. J. Kessler, R. L. Kremens, J. P. Knauer, S. A. Kumpan, S. A. Letzring, D. J. Lonobile, L. D. Lund, F. J. Marshall, P. W. McKenty, D. D. Meyerhofer, S. F. B. Morse, A. Okishev, S. Papernov, G. Pien, W. Seka, R. Short, M. J. Shoup, III, M. Skeldon, S. Skupsky, A. W. Schmid, D. J. Smith, S. Swales, M. Wittman, and B. Yaakobi, "The First Year of ICF Experiments on OMEGA — A 60-Beam, 60-TW Laser System," 16th IAEA Fusion Energy Conference, Montreal, Canada, 7-11 October 1996.

The following presentations were made at the XXVII Annual Symposium on Optical Materials for High Power Lasers, Boulder, CO, 7-9 October 1996:

J. F. Anzellotti, D. J. Smith, R. J. Sczupak, and Z. R. Chrzan, "Stress and Environmental Shift Characteristics of HfO_2/SiO_2 Multilayer Coatings."

S. Papernov and A. W. Schmid, "Heat Transfer from Localized Absorbing Defects to the Host Coating Material in HfO_2/SiO_2 Multilayer Systems."

A. L. Rigatti and D. J. Smith, "Status of Optics on the OMEGA Laser System after 18 Months of Operation."

D. J. Smith, J. F. Anzellotti, S. Papernov, and Z. R. Chrzan, "High Laser-Induced-Damage Threshold Polarizer Coatings for 1054 nm."

The following presentations were made at the OSA Annual Meeting/ILS-XII, Rochester, NY, 20–25 October 1996:

M. S. Adams and D. D. Meyerhofer, "Near Field and Spatial Coherence of the Third Harmonic Produced in a Noble Gas Target."

J. L. Chaloupka and D. D. Meyerhofer, "Second-Harmonic Generation from Oscillating Free Electrons in a Laser Focus."

S. D. Jacobs, "Producing Aspheres with Magnetorheological Finishing."

T. J. Kessler, L. S. Iwan, J. Barone, C. Kellogg, and W. P. Castle, "Optic Fabrication Using Photographic Lithography."

O. A. Konoplev and D. D. Meyerhofer, "Cancellation of the B -Integral for CPA Lasers."

E. M. Korenic, S. D. Jacobs, S. M. Faris, and L. Li, "Colorimetry of Cholesteric Liquid Crystals."

K. S. Lebedev, E. A. Magulariya, S. G. Lukishova, S. V. Belyaev, N. V. Malimonenko, and A. W. Schmid, "Reflective Nonlinearities of Nonabsorbing Chiral Liquid Crystals: Frustration of Selective Reflection by Powerful Laser Radiation."

Y. Lin, T. J. Kessler, and G. N. Lawrence, "Design of Continuous Surface-Relief Phase Plates by Simulated Annealing to Achieve Control of Focal Plane Irradiance."

D. D. Meyerhofer, J. P. Knauer, S. J. McNaught, and C. I. Moore, "Observation of Relativistic Mass Shift Effects during High-Intensity Laser-Electron Interactions."

L. Zheng and D. D. Meyerhofer, "Self- and Cross-Phase Modulation Coefficients in KDP Crystals Measured by a Z-Scan Technique."

The following presentations were made at the Second Annual International Conference on Solid State Lasers for Application to Inertial Confinement Fusion (ICF), Paris, France, 22–25 October 1996:

T. R. Boehly, R. L. Keck, C. Kellogg, J. H. Kelly, T. J. Kessler, J. P. Knauer, Y. Lin, D. D. Meyerhofer, W. Seka, S. Skupsky, V. A. Smalyuk, S. F. B. Morse, and J. M. Soures, "Demonstration of Three Enhancements to the Uniformity of OMEGA: 2-D SSD, New DPP's, and DPR's."

T. J. Kessler, Y. Lin, L. S. Iwan, W. P. Castle, C. Kellogg, J. Barone, E. Kowaluk, A. W. Schmid, K. L. Marshall, D. J. Smith, A. L. Rigatti, J. Warner, and A. R. Staley, "Laser Phase Conversion Using Continuous Distributed Phase Plates."

W. Seka, J. H. Kelly, S. F. B. Morse, J. M. Soures, M. D. Skeldon, A. Okishev, A. Babushkin, R. L. Keck, and R. G. Roides, "OMEGA Laser Performance with Pulse Shaping."

M. D. Skeldon, A. Babushkin, J. D. Zuegel, R. L. Keck, A. Okishev, and W. Seka, "Modeling of an Actively Stabilized Regenerative Amplifier for OMEGA Pulse-Shaping Applications."

The following presentations were made at the 22nd International Congress on High-Speed Photography and Photonics, Santa Fe, NM, 27 October–1 November 1996:

A. Babushkin, W. Seka, S. A. Letzring, W. Bittle, M. Labuzeta, M. Miller, and R. G. Roides, "Multicolor Fiducial Laser for Streak Cameras and Optical Diagnostics for the OMEGA Laser System."

D. K. Bradley and P. M. Bell, "Implementation of 30-ps Temporal Resolution Imaging on the OMEGA Laser System."

The following presentations were made at the Superabrasives Technology Meeting, 7–8 November 1996, Livermore, CA:

B. E. Gillman, Y. Zhou, and S. D. Jacobs, "Coolant/Tool Interactions in Deterministic Microgrinding of Glass."

J. C. Lambropoulos, Y. Zhou, P. D. Funkenbusch, B. Gillman, and D. Golini, "Brittleness and Grindability of Brittle Workpieces."

The following presentations were made at the 38th Annual Meeting, APS Division of Plasma Physics, 11–15 November 1996, Denver, CO:

R. Betti and J. P. Freidberg, "Shape, Pressure, and Kinetic Effects on the Resistive Wall Mode in Rotating Plasmas."

T. R. Boehly, D. D. Meyerhofer, J. P. Knauer, D. K. Bradley, R. L. Keck, J. A. Delettrez, V. A. Smalyuk, J. M. Soures, and C. P. Verdon, "Laser Imprinting Studies Using Multiple-UV-Beam Irradiation of Planar Targets."

D. K. Bradley, J. A. Delettrez, and P. A. Jaanimagi, "Initial Mix Experiments on the 60-Beam OMEGA Laser System."

R. S. Craxton, J. D. Schnittman, and S. M. Pollaine, "Uniformity in Tetrahedral Hohlräume."

J. A. Delettrez, D. K. Bradley, and C. P. Verdon, "Modeling of Mix due to the Rayleigh–Taylor Instability in Burnthrough Experiments Using the One-Dimensional Hydrodynamic Code *LILAC*."

R. Epstein, J. A. Delettrez, D. K. Bradley, C. P. Verdon, U. Alon, and D. Shvarts, "Simulations in One Dimension of Unstable Mix in the Ablation Region in Laser-Driven Plasmas."

Y. Fisher, T. R. Boehly, D. K. Bradley, J. A. Delettrez, D. Harding, and D. D. Meyerhofer, "'Shine-Through' Experiments Using 50-ps Laser Pulses."

A. C. Gaeris, Y. Fisher, J. A. Delettrez, and D. D. Meyerhofer, "Brillouin Scattering of Picosecond Laser Pulses in Preformed, Short-Scale-Length Plasmas."

V. Goncharov, R. Epstein, R. Betti, R. L. McCrory, and C. P. Verdon, "Feedthrough and Spatial-Temporal Evolution of the Ablative Rayleigh–Taylor Instability in ICF."

A. V. Kanaev, C. J. McKinstrie, and J. S. Li, "Spatiotemporal Interaction of Crossed Laser Beams."

J. P. Knauer, "OMEGA Experiments to Characterize the Rayleigh–Taylor Instability with Planar Foils."

J. P. Knauer, D. D. Meyerhofer, T. R. Boehly, D. Ofer, C. P. Verdon, D. K. Bradley, P. W. McKenty, and V. A. Smalyuk, "Initial Single-Mode Rayleigh–Taylor Growth Rates Measured with the OMEGA Laser System."

R. L. Kremens, K. Kearney, M. A. Russotto, B. Taylor, J. D. Zuegel, and M. D. Cable, "A Multichannel Neutron Time-of-Flight Spectrometer for Inertial Confinement Fusion Applications."

F. J. Marshall, D. K. Bradley, M. Cable, J. Delettrez, D. Harding, J. H. Kelly, J. P. Knauer, R. L. Kremens, S. A. Letzring, R. L. McCrory, S. F. B. Morse, J. M. Soures, C. P. Verdon, and B. Yaakobi, "Surrogate Cryogenic Target Implosion Experiments Performed with the OMEGA Laser System."

P. W. McKenty, P. A. Jaanimagi, R. L. Kremens, K. J. Kearney, C. P. Verdon, and M. D. Cable, "Convergence Studies of ICF Implosions Utilizing Doped-CH Ablators to Mitigate Instability Growth."

C. J. McKinstrie, E. J. Turano, and A. V. Kanaev, "Sideward Stimulated Raman Scattering of a Short Laser Pulse in a Plasma Channel."

D. D. Meyerhofer, J. P. Knauer, T. R. Boehly, D. Ofer, C. P. Verdon, P. W. McKenty, V. A. Smalyuk, O. Willi, and R. G. Watt, "Performance of Planar Foam-Buffered Targets on the OMEGA Laser System."

D. D. Meyerhofer, "Observation of Relativistic Ponderomotive Effects in Intense Laser-Electron Interactions."

W. Seka, A. V. Chirikikh, A. Babushkin, R. W. Short, and A. Simon, "Laser-Plasma Interaction Experiments on the 60-Beam OMEGA Laser System."

R. W. Short, "Diffractive Calculation of the Intensity Distribution in a Direct-Drive Laser-Fusion Target Corona."

A. Simon, "Two Types of Raman Scattering in Hohlräume and Gasbags."

V. A. Smalyuk, T. R. Boehly, J. A. Delettrez, L. S. Iwan, T. J. Kessler, J. P. Knauer, F. J. Marshall, and D. D. Meyerhofer, "Comparison of Laser and X-Ray Focal Images for a Single OMEGA Laser Beam."

J. M. Soures, J. P. Knauer, F. J. Marshall, D. K. Bradley, M. D. Cable, R. S. Craxton, R. L. Keck, J. H. Kelly, R. L. Kremens, R. L. McCrory, W. Seka, J. Schnittman, C. P. Verdon, T. J. Murphy, J. Wallace, J. A. Oertel, C. W. Barnes, N. D. Delamater, P. Gobby, A. A. Hauer, G. Magelssen, J. B. Moore, R. Watt, O. L. Landen, R. E. Truner, P. Amendt, C. Decker, L. J. Suter, B. A. Hammel, and R. J. Wallace, "Indirect-Drive Target Irradiation Experiments on OMEGA."

E. A. Startsev and C. J. McKinstrie, "Electron Acceleration by a Laser Pulse in a Plasma."

R. P. J. Town, R. W. Short, and C. P. Verdon, "Fokker-Planck Simulations of Foam-Buffered Targets."

E. J. Turano and C. J. McKinstrie, "Spatiotemporal Evolution of Stimulated Raman Scattering."

B. Yaakobi, F. J. Marshall, and R. Epstein, "High Temperature of Laser-Compressed Shells Measured with Kr^{34+} and Kr^{35+} X-Ray Lines."

The following presentations were made at the Twelfth Topical Meeting on Advanced Solid-State Lasers, Orlando, FL, 27-29 January 1997:

A. Babushkin, W. Bittle, S. A. Letzring, A. Okishev, M. D. Skeldon, and W. Seka, "Stable, Reproducible, and Externally Synchronizable Regenerative Amplifier for Shaped Optical Pulses for the OMEGA Laser System."

A. V. Okishev and W. Seka, "Diode-Pumped, Single-Frequency Pulsed Master Oscillator for the 60-Beam OMEGA Laser Facility."

A. Okishev and W. Seka, "Diode-Pumped, Single-Frequency Nd:YLF Laser for the 60-Beam OMEGA Laser Pulse-Shaping System," LASE '97, San Jose, CA, 8-14 February 1997.

The following presentations were made at the Second International Workshop on Laser Plasma Interaction Physics, Banff, Alberta, Canada, 19-22 February 1997:

A. V. Chirikikh, W. Seka, and J. A. Delettrez, "Scattered Light Spectra Taken on OMEGA Implosion Experiments."

C. J. McKinstrie, A. V. Kanaev, and E. J. Turano, "Sideward Simulated Raman Scattering of a Short Laser Pulse in a Plasma Channel."

R. W. Short and A. Simon, "Collisionless Damping of Plasma Waves Propagating in Filaments and Consequences for SRS in Filaments."

T. R. Boehly, R. L. McCrory, C. P. Verdon, W. Seka, S. J. Loucks, A. Babushkin, R. E. Bahr, R. Boni, D. K. Bradley, R. S. Craxton, J. A. Delettrez, W. R. Donaldson, R. Epstein, D. Harding, P. A. Jaanimagi, S. D. Jacobs, K. Kearney, R. L. Keck, J. H. Kelly, T. J. Kessler, R. L. Kremens, J. P. Knauer, D. J. Lonobile, L. D. Lund, F. J. Marshall, P. W. McKenty, D. D. Meyerhofer, S. F. B. Morse, A. Okishev, S. Papernov, G. Pien, T. Safford, J. D. Schnittman, R. Short, M. J. Shoup III, M. Skeldon, S. Skupsky, A. W. Schmid, V. A. Smalyuk, D. J. Smith, J. M. Soures, M. Wittman, and B. Yaakobi, "Inertial Confinement Fusion Experiments with OMEGA—A 30-kJ, 60-Beam UV Laser," 1997 IAEA Conference, Osaka, Japan, 10-14 March 1997.

R. L. McCrory and J. M. Soures, "Status of Direct-Drive Inertial Confinement Fusion Research at the Laboratory for Laser Energetics," Second Symposium on Current Trends in International Fusion Research Review and Assessment, Washington, DC, 10-14 March 1997 (invited).

The following presentations were made at Ultrafast Electronics and Optoelectronics, Incline Village, NV, 17-21 March 1997:

K. Green, M. Lindgren, T. Y. Hsiang, R. Sobolewski, and L. Fuller, "Picosecond Photoresponse in Polycrystalline Silicon."

D. Jacobs-Perkins, M. Currie, C.-C. Wang, C. Williams, W. R. Donaldson, R. Sobolewski and T. Y. Hsiang, "Analysis of an Interferometric Subpicosecond Imager Based on the Electro-Optic Effect."

M. Lindgren, M. Currie, C. Williams, T. Y. Hsiang, P. M. Fauchet, R. Sobolewski, S. H. Moffat, R. A. Hughes, J. S. Preston, and F. A. Hegmann, "An Ultrafast High- T_c Superconducting Y-Ba-Cu-O Photodetector."

The following presentations were made at the 1997 Applications of High Field and Short Wavelength Sources VII (OSA), Santa Fe, NM, 19-22 March 1997:

B. Buerke, J. P. Knauer, S. J. McNaught, and D. D. Meyerhofer, "Precision Tests of Laser-Tunneling Ionization Models."

J. L. Chaloupka, T. J. Kessler, and D. D. Meyerhofer, "A Three-Dimensional Ponderomotive Trap for High-Energy Electrons."

F. Yang, D. Golini, H. Edwards, D. H. Ragun, and S. D. Jacobs, "Planarization of Gratings Using Magnetorheological Finishing," MRS 1997 Spring Meeting, San Francisco, CA, 31 March-4 April 1997.

The following presentations were made at the Thirteenth International Conference on Laser Interactions and Related Plasma Phenomena (LIRPP), Monterey, CA, 13-18 April 1997:

R. Betti, V. N. Goncharov, R. L. McCrory, and C. P. Verdon, "Feedthrough and Dynamic Stabilization in Convergent Geometry."

T. R. Boehly, V. A. Smalyuk, D. D. Meyerhofer, J. P. Knauer, D. K. Bradley, C. P. Verdon, and D. Kalantar, "The Effect of Increased Irradiation Uniformity on Imprinting by 351-nm Laser Light."

R. L. Keck, A. V. Okishev, M. D. Skeldon, A. Babushkin, and W. Seka, "Pulse Shaping on the OMEGA Laser System."

J. P. Knauer, D. D. Meyerhofer, T. R. Boehly, D. Ofer, C. P. Verdon, D. K. Bradley, P. W. McKenty, V. A. Smalyuk, S. G. Glendinning, and R. G. Watt, "Single-Mode Rayleigh-Taylor Growth-Rate Measurements with the OMEGA Laser System."

D. D. Meyerhofer, C. Bamber, S. Boege, T. Koffas, T. Kotseroglou, A. C. Melissinos, D. Reis, C. Bula, K. T. McDonald, E. Prebys, D. L. Burke, C. Fields, G. Horton-Smith, A. C. Odian, J. C. Spencer, D. Walz, S. Berridge, W. Bugg, K. Shmakov, and A. Weidemann, "Observation of Non-linear Laser-Electron and Laser-Photon Scattering."

W. Seka, A. Babushkin, T. R. Boehly, D. K. Bradley, M. D. Cable, R. S. Craxton, J. A. Delettrez, W. R. Donaldson, D. R. Harding, P. A. Jaanimagi, R. L. Keck, J. H. Kelly, T. J. Kessler, J. P. Knauer, R. L. Kremens, F. J. Marshall, R. L. McCrory, P. W. McKenty, D. D. Meyerhofer, S. F. B. Morse, A. V. Okishev, G. Pien, M. D. Skeldon, J. M. Soures, C. P. Verdon, B. Yaakobi, and J. D. Zuegel, "OMEGA Experimental Program and Recent Results."

W. R. Donaldson, E. M. R. Michaels, K. Akowuah, and R. A. Falk, "Integrated Circuit Tester Using Interferometric Imaging," Diagnostic Techniques for Semiconductor Materials and Devices, Montreal, Canada, 4 May 1997.

The following presentations were made at CLEO/QELS '97, Baltimore, MD, 18-23 May 1997:

O. A. Konoplev, Y. Fisher, I. A. Walmsley, and D. D. Meyerhofer, "Determination of the Third-Order Nonlinearities of Materials Using Frequency-Domain Interferometry."

A. V. Okishev and W. Seka, "Diode-Pumped, Single-Frequency, Pulsed Master Oscillator for the 60-Beam OMEGA Laser Facility."

A. V. Okishev, W. Seka, S. F. B. Morse, J. M. Soures, M. D. Skeldon, A. Babushkin, R. L. Keck, and R. G. Roides, "Pulse-Shaping System Implementation on the 60-Beam OMEGA Laser."

M. D. Skeldon, R. Saager, A. Okishev, and W. Seka, "Thermal Distortions in Laser-Diode- and Flash-Lamp-Pumped Nd: YLF Laser Rods."

S. D. Jacobs and H. M. Pollicove, "Update on Magnetorheological Finishing," 9th International Precision Engineering Seminar, Braunschweig, Germany, 26 May 1997.

The following presentations were made at the 27th Annual Anomalous Absorption Conference, Vancouver, BC, Canada, 1–5 June 1997:

- P. Amendt, O. Landen, S. G. Glendinning, R. E. Turner, M. Cable, C. Decker, L. J. Suter, R. Wallace, B. A. Hammel, J. D. Kilkenny, D. K. Bradley, F. J. Marshall, R. Keck, R. Kremens, W. Seka, J. Schnittman, R. Craxton, C. P. Verdon, R. L. McCrory, T. Murphy, N. Delamater, J. Wallace, C. Barnes, G. Magelssen, P. Gobby, and A. A. Hauer, "Indirect-Drive Time-Dependent Symmetry Control on the OMEGA Laser."
- B. B. Afeyan, A. E. Chou, W. L. Kruer, A. J. Schmitt, R. P. J. Town, and J. P. Matte, "The Nonlinear Evolution of Stimulated Raman and Brillouin Scattering in Inhomogeneous and Nonstationary Plasmas."
- T. R. Boehly, D. D. Meyerhofer, J. P. Knauer, D. K. Bradley, R. L. Keck, J. A. Delettrez, V. A. Smalyuk, W. Seka, C. P. Verdon, R. G. Watt, O. Willi, R. J. Taylor, and D. Kalantar, "The Effect of Increased Irradiation Uniformity on the Stability of Planar Targets Accelerated by UV Light."
- D. K. Bradley, J. A. Delettrez, R. Epstein, F. J. Marshall, B. Yaakobi, and C. P. Verdon, "Measurements of Mix in Spherical-Imploding-Target Experiments Using Doped Layers."
- R. S. Craxton, W. Seka, and R. P. Drake, "Long-Scale-Length Plasma Designs for OMEGA."
- C. Decker, G. Glendinning, P. Amendt, R. E. Turner, O. L. Landen, L. J. Suter, B. A. Hammel, T. J. Murphy, J. Wallace, A. A. Hauer, J. P. Knauer, F. J. Marshall, D. K. Bradley, W. Seka, and J. M. Soures, "Measuring the Hohlraum Radiation Drive on the OMEGA Laser Using Dante and the Foam Ball Technique."
- J. A. Delettrez, D. K. Bradley, and C. P. Verdon, "Two-Dimensional Modeling of Imprint and Feedthrough in OMEGA Mix Spherical Experiments."
- R. Epstein, J. A. Delettrez, and C. P. Verdon, "Simulation of Mix in Laser-Driven Implosions—Compressible Flow in Multimode Modeling of the Rayleigh-Taylor Instability."
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D. C. Wilson, G. D. Pollak, C. J. Fontes, J. Abdallah, P. D. Goldstone, J. A. Cobble, S. R. Goldman, and P. A. Jaanimagi, "Non-LTE Effects on X-Ray Emission from Laser-Irradiated Gold Spheres."

The following presentations were made at the 6th International Superconductive Electronics Conference (ISEC '97), Berlin, Germany, 26-28 June 1997:

M. Currie, C.-C. Wang, R. Sobolewski, and T. Y. Hsiang, "Picosecond Nodal Testing of Centimeter-Size Superconducting Microstrip Interconnects."

M. Lindgren, W.-S. Zeng, M. Currie, R. Sobolewski, S. Cherednichenko, B. Voronov, and G. N. Gol'tsman, "Picosecond Response of a Superconducting Hot-Electron NbN Photodetector."

R. Sobolewski, R. Adam, W. Markowitsch, C. Stockinger, and W. Lang, "Optically Induced Effects in Y-Ba-Cu-O Josephson Junctions and Thin Films."

J. M. Larkin, W. R. Donaldson, T. H. Foster, and R. S. Knox, "Multiphoton Excitation of Rose Bengal by Near-Infrared Light," 25th Annual Meeting of the American Society for Photobiology, St. Louis, MO, 5-10 July 1997.

S. D. Jacobs, W. I. Kordonski, and H. M. Pollicove, "Precision Control of Aqueous Magnetorheological Fluids for Finishing of Optics," 6th International Conference on Electrorheological Fluids, Magnetorheological Suspensions, and Their Applications, Yonezawa, Japan, 22-25 July 1997.

The following presentations were made at the Fifth International Conference on Advances in the Fusion and Processing of Glass, Toronto, Canada, 27-31 July 1997:

S. D. Jacobs, "Deterministic Manufacturing of Precision Glass Optics Using Magnetorheological Finishing (MRF)."

J. C. Lambropoulos, S. D. Jacobs, B. Gillman, F. Yang, and J. Ruckman, "Subsurface Damage in Microgrinding Optical Glasses."

The following presentations were made at SPIE's 42nd Annual Meeting on Optical Science, Engineering, and Instrumentation, San Diego, CA, 27 July-1 August 1997:

B. E. Gillman, B. M. Reed, M. A. Atwood, J. L. Ruckman, D. J. Quesnel, T. T. Ochinerio, and S. D. Jacobs, "Application of Coolants in Deterministic Microgrinding of Glass."

D. Golini, S. D. Jacobs, W. Kordonski, and P. Dumas, "Precision Optics Fabrication Using Magnetorheological Finishing."

S. D. Jacobs, F. Yang, E. M. Fess, J. B. Feingold, B. E. Gillman, W. I. Kordonski, H. Edwards, and D. Golini, "Magnetorheological Finishing of IR Materials."

J. C. Lambropoulos, B. E. Gillman, Y. Zhou, S. D. Jacobs, and H. J. Stevens, "Glass-Ceramics: Deterministic Microgrinding, Lapping, and Polishing."

D. D. Meyerhofer, X. D. Cao, Y. Fisher, O. Konoplev, I. Walmsley, and L. Zheng, "Measurements of Material Properties Using Frequency Domain Interferometry," Ultrafast Optics 1997, Monterey, CA, 4-7 August 1997.