Experimental Investigation

of the Nonlinear Rayleigh-Taylor Instability

in CH Foils Irradiated by UV Light

by

Vladimir A. Smalyuk

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Professor David D. Meyerhofer

Department of Mechanical Engineering

The College

School of Engineering and Applied Sciences

University of Rochester

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Curriculum Vitae

The author was born in a village of Sudilkov, Khmelnitsky region, Ukraine, former USSR on March 6, 1967. He attended the Moscow Institute of Physics and Technology from 1984 to 1990, and received a Master of Science degree in physics in 1990. He worked as a Research Engineer at the Kurchatov Institute of Atomic Energy from 1990 to 1993. He came to the University of Rochester in the Fall of 1993 as a graduate student in the Mechanical Engineering department. He received a LLE Fellowship in 1993-1998. He pursued his research in the experimental investigations of Rayleigh-Taylor instability at ICF relevant conditions under the direction of Professor David Meyerhofer.
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Abstract

In direct-drive inertial confinement fusion (ICF), nonuniformities in the laser drive can create (or imprint) mass perturbations in the target and seed hydrodynamic instabilities. Understanding of this imprinting and the temporal evolution of resulting features is of utmost importance to high-gain target designs. Time-integrated optical measurements of a single OMEGA laser beam show that smoothing techniques such as distributed phase plates (DPP's) and smoothing by spectral dispersion (SSD) generate a broadband spectrum of laser nonuniformities. These nonuniformities are imprinted into targets, then, during a laser-driven acceleration, the Rayleigh-Taylor instability amplifies these features. The temporal evolution of these mass perturbations is studied using x-ray through-foil radiography of planar CH targets employing pinholes and framing cameras. To study these phenomena, it is important that the spatial response of the instrumentation be properly characterized. The noise in these experiments is limited by the photon statistics of the uranium backlighter x rays and the system resolution is limited by 8-μm pinhole. Using the measured system resolution, noise, and sensitivity, a Wiener filter has been constructed to filter the noise from the images and deconvolve the system MTF in order to recover the target's areal density modulations. The Fourier analysis of the radiographic images shows that the perturbations evolve to longer wavelengths and the amplitudes of the shorter wavelengths saturate. The saturation amplitudes and the rates of growth of these features are consistent with the predictions of Haan for broadband spectrum [1]. Using high-contrast, thin 5-μm teflon targets, the imprint has been measured at the time of the shock breakout using the same radiographic system.
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Chapter 1

Introduction

1.1 General principles of inertial confinement fusion (ICF)

Nuclear fusion reactions are a major source of energy in stars. In this process the nuclei of light elements are fused together at very high temperatures to produce heavier nuclei, releasing energy in the process. A reaction of the fusion of two isotopic ions of hydrogen, deuteron (D) and triton (T) is one of the most interesting fusion reactions leading to a release of a large amount of energy. Reaction products are a helium ion (He or $\alpha$-particle) and a neutron (n) which together carry released energy of 17.6 MeV,

$$D + T \rightarrow He + n + 17.6 MeV. \quad (1.1)$$

The tremendous potential of nuclear fusion was demonstrated by the development of the hydrogen bomb in the early 1950s. Today nuclear fusion is a very attractive potential source of energy for our civilization. It is clean (products of reactions are not radioactive) and the fuel is abundant in sea water and practically inexhaustible on earth.
The goal of nuclear fusion research is to create a high temperature and high density DT plasma (with parameters similar to and even more extreme than those in the stars) in the laboratory in order to produce more energy through fusion reactions than is required to heat and sustain such plasma. There are two major approaches in fusion research: magnetic confinement fusion (MCF) and inertial confinement fusion (ICF). MCF uses magnetic fields to confine a plasma with relatively low density \( (n \sim 10^{14} \text{ cm}^{-3}) \) for long times \( (\tau \sim 1 \text{ s}) \). ICF relies on the inertia of the fuel mass to provide confinement for short times \( (\tau \sim 10^{-10} \text{ s}) \) at high densities \( (n \sim 10^{25} \text{ cm}^{-3}) \). Today the prospects of a usable MCF reactor are dimmer than they were a few years ago, while ICF has become more promising after several decades of development.

### 1.2 Early research in ICF

With the invention of the laser in 1960, physicists throughout the world immediately recognized its utility for the inertial fusion. In 1968, thermonuclear neutrons were detected in experiments with laser irradiated LiD targets \([2]\), which proved that lasers were able to heat plasmas to thermonuclear temperatures. In 1972 Nuckolls proposed laser-driven implosions of DT shells as a way to obtain positive gain from thermonuclear reactions \([3]\). At that time laser technology and numerical modeling were becoming mature enough to evaluate the requirements for the success of ICF. Nuckolls has estimated that a 1 kJ laser would be able to achieve a 'breakeven' point (where the energy released in fusion reactions is equal to the energy required to create and sustain the plasma) and at the laser energy of 1 MJ the estimated gain was about 100. In his model, implosions with high laser intensities, approaching \( 10^{17} \text{ W/cm}^2 \), were required to obtain ignition at laser energies of 1 kJ.
1.3 Direct-drive ICF

The first implosion experiments used a direct-drive configuration, in which targets are driven directly by overlapped laser beams (see Fig. 1.1(a)). The first targets were hollow glass spheres containing DT gas. The sphere explodes when it is rapidly heated by a short, high-intensity laser pulse. The shell of the sphere expands outward, while the rest of the sphere is forced inward to conserve momentum. In these so-called 'exploding pusher' implosions high temperatures \( \sim 5 \) keV have been achieved but at low densities \( (< 1 \text{ g/cm}^3) \) [4]. At high laser intensities \( 10^{17} \text{ W/cm}^2 \), a large number of hot electrons and hard x-rays preheat the fuel rapidly, resulting in its low compression. The 'exploding pusher' implosions are unsuitable for achieving high gains.

In 'ablative compression' implosions, a fuel capsule is coated with a thick, low-Z ablator and irradiated by a long pulse at low laser intensities. The ablator serves to absorb the incident energy and to provide high-compression of the core by the rocket reaction of the ablated material [3]. This approach toward implosions is accepted by the ICF community as a way to achieve ignition and high gain.

In the early 1970s, the hydrodynamic Rayleigh-Taylor (RT) instability [5, 6]
was recognized as the major factor in degrading target performance and inhibiting
the achievement of the ignition during an implosion. A simple classical example
of the RT instability is the case when a layer of water is supported by a layer
of oil in a gravitational field. The density of water is higher than the density
of oil. The interface between two fluids will remain stationary in the absence of
any perturbations. However, the amplitude of any perturbation between these
two fluids will grow exponentially, and eventually the two fluids will exchange
positions.

Under ICF conditions, the RT instability occurs when the laser-heated low
density plasma accelerates high density fuel during implosion [7]. In the region
around the ablation surface, the density gradient is directed toward the target,
but the pressure gradient is directed toward the laser. In the case of two classical
fluids (see the example above) the directions of density and pressure gradients
are also opposite. Therefore, the interface between the cold, high density target
and hot, less dense plasma is unstable. Any nonuniformities present on the target
are amplified by the RT instability destroying the symmetry of the implosion and
compromising its performance.

In the mid-1970s direct-drive experiments around the world using 1 µm neodymium
glass lasers showed reduced laser light absorption and hot electron production by
parametric laser-plasma instabilities [8], which severely degraded implosions at
high laser intensities. In addition, the spatial quality of laser beam was also much
worse than could be tolerated for the implosion uniformity required. Laser beam
nonuniformities provide a source of perturbations (referred to as laser imprinting)
[9], which are further amplified by RT instability during implosion.

In the last decade, the direct-drive approach to ICF has succeeded in deval-
oping efficient conversion of 1 µm laser light to second and third harmonics [11],
which dramatically increased its absorption in plasma and reduced the deleterious
consequences of laser-plasma parametric instabilities such as laser light scattering and hot electron production. Several different laser smoothing techniques have been implemented on direct-drive laser systems. These have dramatically increased the uniformity of laser irradiation, reducing the initial seed to RT instability.

1.4 Indirect-drive ICF

In order to reduce the requirements on laser-beam uniformity and to decrease the sensitivity to hydrodynamic instabilities some laboratories have concentrated most of their efforts on indirect-drive approach to ICF [10] since the mid 1970’s. In indirect drive (see Fig. 1.1(b)), the target is located inside a high-Z enclosure, a 'hohlraum'. About 80% of the laser energy, absorbed by the walls of 'hohlraum', is converted to x-rays, which drive the capsule implosion. In the indirect-drive approach the drive uniformity has been increased dramatically and the growth of perturbations due to the RT instability has been reduced relative to direct-drive ICF. This occurs because x-ray driven implosions have much higher ablation rates and hence lower RT growth. A potential disadvantage of indirect-drive ICF is a high possibility that parametric instabilities, easily exited in long scale-length plasma inside the 'hohlraum', may scatter laser light (degrading drive uniformity) and produce hot electrons which can cause target preheat, reducing a target compression.

Both direct and indirect drive approaches promise achievement of the ignition and even moderate gain $\sim 20-30$ [12] on the 1.8 MJ national ignition facility (NIF) [12], which is currently is under construction in Livermore, California.
1.5 Requirements for the ignition in direct-drive ICF

After nearly three decades, theoretical and experimental research has defined the conditions for successful achievement of the ignition and high gain, of the order of 100, for direct-drive laser ICF [13]. Using DT fuel, a high compression approaching 1000 times liquid density is required, together with maximum temperatures of about 5 keV.

In order to achieve these conditions for high gain at reasonable driver energies (a few megajoules or less), incident intensities of the order of $10^{15}$ W/cm$^2$ are required with 0.351 $\mu$m laser light. At these intensities the incident irradiation is strongly absorbed by the plasma and parametric instabilities do not scatter much of the light out of plasma [14] or create hot electrons.

For one of several designs, the 'all DT' design, the target is a spherical shell filled with low-density gas ($\sim 1$ mg/cm$^3$). The shell (with radius of 2.8 mm) consists of frozen DT (with thickness of 340 $\mu$m), the main fuel. A laser pulse (about 15 ns) generally has a low intensity foot after which the intensity rises up to about $10^{15}$ W/cm$^2$. The laser pulse shape has to be designed such a way that it compresses the fuel optimally during the implosion and limits the effects of the target nonuniformity growth due to RT instability.

A direct-drive laser system must have a large number of beams in order to ensure symmetric drive, critical to the success of an implosion. In order to minimize the target perturbations coming from the imprinting of laser nonuniformities, direct-drive laser systems normally employ such laser smoothing techniques as random phase plates [15], smoothing by spectral dispersion [16] and polarization smoothing [17]. These smoothing techniques will be discussed in Chapter 3 along with quantitative measurements of their performance on the OMEGA laser.
system.

1.6 Summary

In inertial confinement fusion (ICF), the Rayleigh-Taylor (RT) instability can amplify target perturbations sufficiently to disrupt the implosion and degrade its performance. These perturbations can result from existing target imperfections and, in the case of direct-drive ICF, laser imprinting. To achieve high-gain ICF implosions, one must control laser imprinting and the hydrodynamic stability of the targets. The Rayleigh-Taylor evolution of laser-imprinted perturbations is the main subject of this thesis.

Chapter 2 of this work is devoted to the review of theoretical and experimental studies in the field of Rayleigh-Taylor instability and laser imprinting. The first part of this Chapter reviews the classical linear RT instability and its modification for conditions relevant to ICF. In the nonlinear regime, the RT instability is presented by the following nonlinear models: the mode coupling model, Haan’s saturation model for broadband spectrum, bubble competition and coalescence model. The second section of this Chapter reviews studies of imprinting physics. It includes a description of the impulsive Richtmyer-Meshkov instability, a model for nonuniform shock propagation, and the 'cloudy-day' model of laser imprinting.

The success of direct-drive ICF depends on the effective laser imprinting reduction by means of single-beam laser smoothing techniques. Chapter 3 describes such techniques as distributed phase plates (DPP’s), smoothing by spectral dispersion (SSD), and distributed polarization rotators (DPR’s). These techniques are effective in smoothing the laser light generated by solid-state laser systems. Quantitative measurements of single-beam uniformity improvements on OMEGA are presented in Chapter 3.
Face-on, through-foil x-ray radiography is the primary experimental technique to study RT evolution of imprinting in planar targets. Chapter 4 describes methods and principles for target perturbation measurements with this technique. It contains experimental characterization of system sensitivity, resolution, and noise. Using this information, a Wiener filter, which distinguishes signal from noise, is applied to filter the noise and deconvolve the finite system resolution from the signal, allowing observation of the evolution of target perturbations.

Using this experimental method, the saturation of RT growth was measured for broadband initial spectrum of target nonuniformities, generated by laser imprinting. The results of these measurements are in agreement with the Haan model for broadband spectrum saturation [1]. These experiments are discussed in Chapter 5.

Chapter 6 presents results of experiments, where high-contrast teflon targets are used to measure the imprinting spectrum early in time, prior to the acceleration phase (in contrast with the experiments from Chapter 5).

Chapter 7 summarizes all experiments and findings presented in this thesis.
Chapter 2

Linear Theory and Nonlinear Models of the Rayleigh-Taylor Instability

2.1 Introduction

In both direct and indirect drive ICF implosions a target is accelerated by the ablation of material heated by laser or x-ray light, respectively. Near the ablation surface the lighter ablating fluid pushes the heavier compressed fluid. This process is Rayleigh-Taylor (RT) unstable when the density gradient is directed towards the target but pressure gradient is directed at the laser. Any mass perturbation present at the ablation surface, due to target imperfections or imprinted from laser nonuniformities, will grow, distorting or possibly destroying the target [18]. That is why RT instability is of utmost concern to ICF [12]. It is therefore critical to the success of ICF that the RT linear growth and nonlinear effects be measured and understood.

This Chapter reviews the theoretical and experimental work performed in the area of the RT instability. First, an analytical derivation of classical RT instability growth rates is presented. It is followed by a discussion of stabilizing effects relevant to ICF conditions. Then, the nonlinear models of RT instability
such as mode-coupling, Haan's model for the saturation of a broadband spectrum growth, and bubble competition and coalescence are reviewed.

Target perturbations coming from laser nonuniformities (laser imprinting) are ones of the most dangerous for direct-drive ICF (the inner surface perturbations of DT ice have the same or even higher importance). The understanding the imprinting development is crucial to a success of direct-drive ICF. The second part of this chapter reviews models of imprinting development, imprint's evolution caused by impulsive Ritchmeyer-Meshkov (RM) instability and nonuniform shock propagation at the imprinting stage, i.e. before the shock reaches the rear side of the foil and the target begins to accelerate.

2.2 Rayleigh-Taylor instability

2.2.1 Classical Rayleigh-Taylor instability

This section is devoted to the derivation of growth-rates for the classical Rayleigh-Taylor (RT) instability. A heavier fluid could stay in equilibrium supported by a lighter fluid in a gravitational field. However, this equilibrium is unstable for any perturbation at the surface dividing these two fluids. This perturbation will grow as a result of the RT instability and eventually all heavier fluid will move underneath and support the lighter fluid. Figure 2.1 shows a schematic of a fluid with density $\rho$ underneath of a fluid with density $\rho_1$ in the field of gravity with the acceleration $g$. The depths of these two fluids are $h$ and $h_1$, respectively. For simplicity let's consider that the upper and lower fluid boundaries do not move in the problem (as shown in the Fig. 2.1). Coordinates $x$ and $y$ are along the surface dividing two fluids which has a sinusoidal perturbation at $z = \xi(x)$. Following an argument from Landau and Lifshitz's text book [19], let's investigate the
Figure 2.1: Classical RT instability. Lighter fluid with density $\rho$, velocity potential $\phi$ and height $h$ is underneath of the heavier fluid with density $\rho_1$, velocity potential $\phi_1$ and height $h_1$.

Stability of the surface dividing two fluids using the fluid equations

$$\frac{\partial \rho}{\partial t} + \nabla \cdot (\rho \mathbf{v}) = 0, \quad (2.1)$$

$$\frac{\partial \mathbf{v}}{\partial t} + (\mathbf{v} \cdot \nabla)\mathbf{v} = -\nabla\left(\frac{P}{\rho}\right) + \mathbf{g}, \quad (2.2)$$

where $\rho$, $\mathbf{v}$ and $P$ are the fluid's density, velocity and pressure, respectively. If fluid velocities are small in the problem, it is possible to neglect nonlinear term $(\mathbf{v} \cdot \nabla)\mathbf{v}$ compared to the term $\frac{\partial \mathbf{v}}{\partial t}$ in Eq. (2.2). The conditions at which this is possible can be derived from the following argument. The fluid velocity $\mathbf{v}$ can be approximated by the expression $\mathbf{v} = \xi / \tau$, where $\tau$ is the time at which the interface perturbation has grown up to the amplitude of $\xi$. The velocity's time derivative is the order of $\mathbf{v} / \tau$, while the spatial derivative is the order of $\mathbf{v} / \lambda$, where $\lambda$ is a wavelength of the perturbation. So the condition $(\mathbf{v} \cdot \nabla)\mathbf{v} \ll \frac{\partial \mathbf{v}}{\partial t}$ is equivalent to

$$\frac{1}{\lambda} \left(\frac{\xi}{\tau}\right)^2 \ll \frac{\xi}{\tau}, \quad (2.3)$$

or

$$\xi \ll \lambda, \quad (2.4)$$
which means that the perturbation amplitude has to be much lower than the perturbation wavelength. Assuming that the fluid is incompressible ($\rho = \text{const}$) and perturbation amplitude is much smaller then it’s wavelength ($\xi \ll \lambda$) the fluid equations are reduced to the following

$$\Delta \phi = 0$$  \hspace{1cm} (2.5)

$$P = -\rho gz - \rho \frac{\partial \phi}{\partial t},$$  \hspace{1cm} (2.6)

where $\phi$ is the velocity potential, defined as $\mathbf{v} = \nabla \phi$. Equations (2.5) and (2.6) are satisfied for both heavier and lighter fluids. On the surface dividing two fluids where $z = \xi(x)$, the pressure is continuous which leads to the following equation

$$\rho g \xi + \rho \frac{\partial \phi}{\partial t} = \rho_1 g \xi + \rho \frac{\partial \phi_1}{\partial t}.$$  \hspace{1cm} (2.7)

Using

$$v_z = \frac{\partial \phi}{\partial z} = \frac{\partial \xi}{\partial t},$$  \hspace{1cm} (2.8)

and taking partial time derivative in Eq.(2.7), it becomes (on the surface dividing two fluids at $z = 0$)

$$(\rho - \rho_1) g \frac{\partial \phi}{\partial z} = \rho_1 \frac{\partial^2 \phi_1}{\partial t^2} - \rho \frac{\partial^2 \phi}{\partial t^2}.$$  \hspace{1cm} (2.9)

The fluid velocities are continuous at the interface, which means that at $z = 0$

$$\frac{\partial \phi}{\partial z} = \frac{\partial \phi_1}{\partial z}.$$  \hspace{1cm} (2.10)

Equations(2.8) and (2.9) are the system which has to be solved. Solutions in both fluids are

$$\phi = C_1 \cosh(k(z + h)) \cos(kx - \omega t)$$

$$\phi_1 = C_2 \cosh(k(z - h_1)) \cos(kx - \omega t).$$  \hspace{1cm} (2.11)

Solutions (2.11) satisfy the condition that the velocity $v_z = 0$ on the top and the bottom boundaries. Substitution of Eq.(2.11) into Eq.(2.9) and Eq.(2.10) gives
two linear equations for $C_1$ and $C_2$, which can be solved giving [19]

$$\omega^2 = \frac{k g (\rho - \rho_1)}{\rho \coth(kh) + \rho_1 \coth(kh_1)}. \quad (2.12)$$

Let's consider the case when both fluids are very deep ($kh$ and $kh_1 \gg 1$). In case of the heavier fluid on top of the lighter fluid, one of the the solutions is unstable with the following growth rate,

$$\gamma = \sqrt{Akg}, \quad (2.13)$$

where $A$ is the Atwood number,

$$A = \frac{\rho_1 - \rho}{\rho + \rho_1}. \quad (2.14)$$

For conditions relevant to ICF, $\rho_1 \gg \rho$, so $A \simeq 1$. Surface tension and viscosity stabilize the RT instability at short wavelengths, however for ICF conditions this occurs at too short wavelengths to be of any significance. Fortunately, there are other physical processes stabilizing the RT instability during ICF implosions.

### 2.2.2 Rayleigh-Taylor instability in ICF

Under conditions relevant for ICF, the low-density plasma accelerates the high-density ablator (CH) or fuel (DT) with high pressure coming from the ablation of material from the unstable surface. Normally, the density profile during implosion is not sharp, as in the classical case discussed in previous section, but has a density gradient scale length, $L_m = \rho/\nabla \rho$. Typically, this length is in the range from $L_m \sim 1$ to $10 \mu m$. It is bigger for ablators with higher $Z$ [20]. The density gradient has a stabilizing effect which is explained via effective reduction of Atwood number as shown in Fig. 2.2. According to the Eq. (2.11), the distance with which the perturbation $\xi(z)$ extends into the fluid is $\sim \exp(-kz)$. If the density profile has a scale length $L_m$, the density change, which perturbation 'sees', is
Figure 2.2: The density gradient scale length stabilization of RT growth. The distance with which the perturbation $\xi(z)$ extends into the fluid is $\sim \exp(-kz)$. The density profile has a scale length $L_m$. The effective Atwood number is reduced by the factor $\sim \exp(-kz)$ because of the density gradient scale length $L_m = \rho/\nabla \rho$.

reduced from $(\rho_1 - \rho)$ to $(\rho_1 - \rho) \exp(-kL_m)$. Then, the effective Atwood number is reduced by the factor of $\exp(-kL_m)$. In the case of $kL_m \ll 1$ this factor can be approximated by $1/(1 + kL_m)$ giving the resulting growth rate

$$\gamma = \sqrt{kg/(1 + kL_m)}. \quad (2.15)$$

Ablative stabilization is another important effect which reduces the classical growth rates in conditions relevant to ICF. In the case of the classical RT instability the perturbation grows with its classical growth rate $\gamma_{cl}$ as $\xi(t, z) \sim \exp(\gamma_{cl} t) \exp(-k|z|)$. In ablative accelerated targets the perturbation $\xi(t, z)$ is being removed (or ablated) by material flow away from the unstable interface. Because of the ablation, the unstable interface moves a distance $z = v_a \Delta t$ during the time interval $\Delta t$. As a result, the amplitude of the perturbation is reduced by a factor of $\exp(-kv_a \Delta t)$ as shown in the Fig. 2.3. The net growth factor at
Figure 2.3: Ablative stabilization of RT growth. It is explained by a reduction of the perturbation amplitude $\xi(t)$ because of the material flow away from unstable interface.

the ablation surface is $\exp(\gamma_d \Delta t - k v_a \Delta t)$, resulting in lower growth rate

$$\gamma = \gamma_d - k v_a. \quad (2.16)$$

More rigorous analytic and numerical studies have examined the RT growth rates in conditions relevant to ICF. In his earlier work [21, 22], Takabe numerically found the steady-state equilibrium with material flow through the absorption region, and numerically calculated the eigenfunctions of the problem. As a result, the best fit to the resulting growth rates was given by the following dispersion relation, which is valid for direct drive with $k L_m \ll 1$,

$$\gamma = 0.9 \sqrt{k g} - 3 k v_a. \quad (2.17)$$

Recently, linearized and perturbed fluid equations were solved analytically by Betti and Goncharov as an eigen value problem [20, 23, 24, 25, 26] giving the growth rates of the RT instability in a self-consistent treatment including ablative and density gradient stabilization, thermal and radiative transport. The result is
given in a dispersion relation form,

\[ \gamma = \alpha \sqrt{\frac{k g}{1 + kL_m}} - \beta k V_a, \]

(2.18)

where \( \alpha \), \( \beta \) and \( L_m \) are functions of laser intensity and material parameters. For example, for typical direct-drive ICF laser intensity \(~ 10^{14} \) W/cm\(^2\) with flat top pulse shape and for 20-\(\mu\)m-thick CH targets: \( \alpha = 0.99 \), \( \beta = 1.7 \), and \( L_m = 1.1 \) \(\mu\)m. For 200 \(\mu\)m DT targets at the same intensity and same pulse shape: \( \alpha = 0.94 \), \( \beta = 2.6 \), and \( L_m = 0.49 \) \(\mu\)m [20].

Ablative stabilization and the density scale length are quite different for DT and CH targets resulting in quite different growth rates. For direct drive, with small radiation effects, the Takabe formula is found to be an excellent approximation. However, only Betti-Goncharov analytical theory is acceptable for more general cases.

The ablative RT instability has been experimentally studied for about two decades using both direct and indirect drive with pre-imposed initial perturbations on planar, cylindrical and spherical targets. X-ray, face-on, through-foil radiography [7, 27, 28, 29, 30, 31] is the most powerful technique to study the RT growth, allowing quantitative measurements of the areal density (the product of the target's thickness and density) of the target in planar geometry. In this technique the x-rays from a backlighter traverse the target and create images of target perturbations. Chapter 4 is devoted to the detailed characterization of such an experimental configuration. The other experimental techniques used to study the RT growth are the side-on radiography [32, 33, 34], and the burnthrough of buried tracer layer [35, 36, 37] are less quantitative (for studies of 3-D perturbations; side-on radiography, however, is acceptable for studies of 2-D perturbations), although they are very useful where face-on, through-foil radiography does not work, especially together with numerical predictions of the burnthrough times in
case of a buried tracer layer technique. For example, the side-on radiography has been used in cylindrical geometry [38, 39], while the burnthrough technique can be used in spherical geometry [37].

The first face-on radiographic measurements were made in the early 1980s using 2-D pre-imposed perturbations in planar CH foils. They were only qualitative because of small laser spots driving the foils [40] and low driving intensities [41, 42]. During the early 1990s, the series of RT experiments using indirect x-ray drive has been performed on NOVA laser system by Remington et al. [30, 43, 44]. These experiments were very well simulated by hydrocodes. Using smoothed laser beams at high laser intensities ($\sim 10^{13}$ W/cm$^2$) the RT growth has been measured in direct-drive by Desselberger [45] and Glendinning [46]. However, the large initial amplitudes (close to saturation amplitudes) used in these works and the nonuniform shock propagation [48, 49, 50] did not permit growth rate measurements of the RT instability directly.

Only recently the linear growth rates of RT instability have been measured in planar targets accelerated by a direct drive acceleration [47, 51, 50]. Numerical simulations are still necessary to explain all the experimental details. Glendinning found that simulated growth rates (based on classical Spitzer-Harm electron heat transport) were $\sim 18\%$ higher than measured values [47]. Shigemori found that measured growth rates in his experiment are well simulated with Fokker-Planck (FP) treatment of electron heat transport and are $\sim 30\%$ lower than predicted by simulations based on Spitzer-Harm (SH) electron heat transport [51]. The difference between the SH and FP electron heat transports is that the FP electron heat transport results in lower target density because of a preheat by energetic electrons originating in plasma corona which penetrate beyond the ablation surface. Lower target density results in higher ablation velocities and subsequently stronger ablative stabilization of the growth rates. Knauer concluded from his ex-
periments, that measured perturbation evolution is well simulated using SH electron heat transport and that the apparent growth rate reduction was explained by the early time nonuniform shock propagation [50]. Recently an experimental comparison of the growth rates in classical and ablative RT instability was performed by Budil using indirect drive [52].

2.2.3 Mode coupling in weakly nonlinear stage of RT instability

As the individual modes grow in the linear regime, the amplitudes of the fastest growing modes eventually reach the value where their root-mean square (rms) amplitude is an appreciable fraction of their wavelength (the \((\mathbf{v} \cdot \nabla)\mathbf{v}\) term in Eq. (2.2) can no longer be neglected). At this point the growth becomes nonlinear. The first nonlinear effect is a mode coupling, when the fastest growing modes couple to drive up harmonics and coupled modes. For the case of a single-mode initial spectrum, it is generally recognized that a mode with wavelength \(\lambda\) growth exponentially until it reaches an amplitude of about \(\xi_k \sim 0.1\lambda\) [1, 53]. Deviations of the actual amplitude from the exponential growth value on the order of a few percent are already present when \(\xi_k \simeq 0.02\lambda\) [54]. Third-order perturbation theory solution for unstable interface \(\xi(x)\) in case of the single-mode initial spectrum is [55],

\[
\xi(x) = \xi_k \cos(kx) - \frac{1}{2}(k\xi_k)\cos(2kx) + \frac{3}{8}(k\xi_k)^2\cos(3kx) - \frac{1}{4}(k\xi_k)^2\xi_k\cos(kx),
\]

(2.19)

where \(\xi_k = \xi_k(t = 0)\exp(\gamma(k)t)\), \(\xi_k(t = 0)\) is the initial amplitude, and \(\gamma(k)\) is the instability growth rate. As the ratio \(\xi_k/\lambda\) grows, harmonics of the primary mode due to a second-order term (at wavenumber 2\(k\)) and a third-order term (at wavenumber 3\(k\)) arise. A third-order term also causes a reduction in the primary mode growth rate. These effects become noticeable when \(\xi_k/\lambda \simeq 0.1\)
Figure 2.4: Harmonic generation for a single mode initial spectrum. The solid line is the foil modulation at the ablation surface in optical depth after $\sim 1.5$ ns of acceleration which has grown from the initial sinusoidal perturbation at wavelength $60 \, \mu m$ with the amplitude $a_0 = 0.5 \, \mu m$. The dashed line, which serves as a guideline, is the nonlinear sinusoidal fit to the data which includes up to six harmonics of the fundamental spatial frequency.

(or $k \xi_k \sim 1$), where the primary mode amplitude reduction is about 10%, and the second and third harmonic amplitudes are 30% and 15% of the amplitude of the fundamental mode, respectively. However, the exponential growth in an acceptable approximation up to about $\xi_k \approx 0.1$-0.2 $\lambda$ for many purposes [55].

Harmonics of the fundamental mode, generated by mode coupling during the exponential (linear) phase, lead to the formation of bubbles (penetration of lighter fluid into heavier) and spikes (penetration of heavier fluid into lighter). Figure 2.4 shows the effect of such nonlinear behavior (mode coupling) on the shot where the initial sinusoidal perturbation with 60-$\mu$m-wavelength and initial amplitude of $a_0 = 0.5 \, \mu m$ on 20-$\mu$m-thick CH foil has grown during the foil acceleration. It was driven by the 351-nm laser light at an intensity of $2 \times 10^{14} \, W/cm^2$ with a 3-ns flat-top pulse shape. The target modulation in optical depth, presented
by a solid line, after \( \sim 1.5 \) ns of acceleration is no longer purely sinusoidal. The modulation in optical depth is proportional to the amplitude of the ablation surface perturbation. The details of experimental techniques will be presented in Chapter 4. The dashed line, which serves as a guideline, is the nonlinear sinusoidal fit to the data which includes up to six harmonics of the fundamental spatial frequency. The sharp spikes (bottom of the Fig. 2.4) and broad bubbles (top of the same figure) are very pronounced indicating harmonic generation is occurring during an acceleration and a growth.

Using perturbation theory solution, Haan formulated mode coupling model for a broadband initial spectrum. Modes with wavevectors \( \mathbf{k} \) and \( \mathbf{k}_2 \) drive \( 2\mathbf{k}, 2\mathbf{k}_2, \mathbf{k} + \mathbf{k}_2 \) and \( \mathbf{k} - \mathbf{k}_2 \) in his second-order ideal fluid theory \cite{53}. Assuming that the flow is incompressible, Laplace's equation is \( \nabla^2 \phi = 0 \) and Bernoulli's equation is \( P = \rho (d\phi/dt - 1/2u^2 - gz) \), where \( \phi \) is the velocity potential (\( u = \nabla \phi \)) and \( u \) is the fluid velocity. This can be expanded up to second order in the Fourier mode amplitudes \( \xi_k \), resulting in the equation,

\[
\ddot{\xi}_k - \gamma^2(k)\xi_k + A_k \sum_{k_2} \left( \xi_{k_2}^2 \xi_{k_2'}^2 \left( 1 - k_2 \hat{k} \right) + \xi_{k_2} \xi_{k_2'} \left( 1/2 - k_2 \hat{k} - 1/2 k_2^2 \right) \right) = 0,
\]

(2.20)

where \( \hat{k} = k / k \) is a unit vector, \( k_2' = k_2 - k \), and \( A = (\rho_1 - \rho)/(\rho_1 + \rho) \) is the Atwood number. Equation (2.20) may be solved to second-order accuracy, neglecting terms of order \( \xi^3 \). The result is \cite{53},

\[
\xi_k = \xi_k^{\exp} + 1/2kA \left( \sum_{k'} \xi_{k'}^{\exp} \xi_{k+k'}^{\exp} - 1/2 \sum_{k' < k} \xi_{k'}^{\exp} \xi_{k-k'}^{\exp} \right),
\]

(2.21)

where \( \xi_k^{\exp} = \xi(t = 0) \exp(\gamma(k)t) \) is an exponential (or first-order) amplitude of mode \( \mathbf{k} \). Equation (2.20) and its solution (2.21) are valid only during very early weakly nonlinear stage of the perturbation evolution. In a more strongly nonlinear regime it quickly diverges from the exact solution. Haan found that, in case of classical RT instability, the mode coupling from high-frequency modes
(most unstable) has profound effect on the evolution of all other modes (much stronger than the exponential growth of these modes) even if they are in the linear regime. However, under the conditions relevant to ICF (with ablative stabilization of high-frequency modes), the effect of mode coupling on any mode is comparable to the exponential growth of this mode (in the linear regime). He concluded that under these conditions (relevant to ICF) his other model for evolution of broadband initial spectrum (which simply neglects mode coupling) is reasonable [1]. This model will be presented in the next section of this Chapter.

Most of the RT growth experiments with single mode initial spectra show nonlinear harmonic generation with formation of bubbles and spikes. In addition to single mode experiments, the mode coupling experiments with two or more modes of the initial spectrum (up to eight modes) have been performed using x-ray drive by Remington [30, 56, 57]. Budil [58] has demonstrated the coupling of two modes (each of which could not be resolved by the diagnostics) by the appearance of longer-wavelength coupled mode, once the growth has proceeded into the nonlinear regime. The mode coupling model [53] and numerical hydrocodes have been used to interpret the experimental data. Some nonlinear behavior has been measured in indirect-drive RT experiments with pre-imposed 3-D multimode initial perturbations in planar targets [59, 60].

2.2.4 Haan’s nonlinear saturation model for broadband spectrum

For a single mode initial perturbation spectrum, nonlinear effects cause the exponential growth of the mode to saturate at an amplitude of about $\xi_k = 0.1\lambda$ and to subsequently grow linearly in time with a constant velocity (see previous section). Haan pointed out that “for broadband initial spectrum of modes, modes of similar wavelengths can constructively interfere to create local structures with an
amplitude much larger then the individual mode’s amplitude” [1]. The modified saturation amplitude can be estimated by taking the rms of the amplitudes of modes in the vicinity of wavevector \( \mathbf{k} \) and comparing this to the "single mode saturation amplitude" \( 0.1\lambda \) [1],

\[
\left( \sum_{k'} \xi_{k'}^2 \right)^{1/2} \sim 0.1\lambda, \tag{2.22}
\]

where \( \frac{|k' - k|}{k} \leq \varepsilon \). Typically \( \varepsilon = 0.1 - 0.5 \), so \( \frac{1}{\varepsilon} \) is the number of spatial cycles before the modes fall out of phase and no longer construct the local structures responsible for the saturation. The model [1] assumes that a mode with wave number \( k \) grows exponentially until reaching the saturation amplitude (in 3-D) given by [61, 62]

\[
S(k) = 2/Lk^2, \tag{2.23}
\]

where \( L \) is a size of the analysis box. One can notice that the \( 1/k^2 \) dependence of the saturation amplitude \( S(k) \) occurs because the \( \lambda = 2\pi/k \), and area of the Fourier region under the summation sign in Eq.( 2.22) is proportional to \( k^2 \). It was also assumed that \( \xi_k^2 = S^2(k) \). The numerical coefficient 2 in Eq.( 2.23) was derived by normalizing the model to the experimental results of Read and Youngs [63, 64], who measured the rms bubble front amplitude \( b(t) = 0.035Agt^2 \) and the rms spike front amplitude \( s(t) = 0.035(1 + A)Agt^2 \) in the nonlinear regime. The \( L \) dependence occurs because the individual Fourier amplitudes of the broadband features depend on the size of analysis region; whereas the rms amplitude does not. The number of Fourier modes decreases as the box size is reduced. The nonuniformity’s sigma rms is a square root of the sum of all modes absolute values squared so the amplitudes of the modes must, concomitantly, increase to keep the nonuniformity’s sigma rms constant.

In Haan’s model, modes are assumed to grow linearly in time after reaching
the saturation amplitude $S(k)$, as

$$|\xi_k(t)| = S(k) \left( 1 + \ln \frac{|\xi_k^{exp}(t)|}{S(k)} \right). \quad (2.24)$$

This model has several desirable features. If $\xi_k^{exp}(t)$ grows exponentially at some growth rate, $\xi_k(t)$ grows secularly at a correct bubble velocity (it depends on the bubble scale-length). The transition from the linear to the nonlinear stage is continuous. The rms amplitude of the total spectrum $\sigma(t)$ is used to estimate the net bubble amplitude $b(t) = \sqrt{2}\sigma(t)$, and the spike amplitude $s(t) = (1 + A)\sqrt{2}\sigma(t)$. Haan’s model, however, does not provide a meaningful approximation for the full spectrum, because it does not account for the phases of individual modes making up the bubbles and spikes. For example, the bubble amplitude for a growing single mode is constructed from the single spatial frequency and all its harmonics as

$$b_k = \sum_n (-1)^n |\xi_{nk}|, \quad (2.25)$$

while the spike amplitude is the sum of all modes $\xi_{nk}$ without alternating signs. In a more complete modal description of $b(t)$, the tendency for $\xi_{2k}$ to cancel $\xi_k$ is required. Haan’s model neglects all these details.

In direct-drive ICF it is important to understand the evolution of broadband initial spectra (such as produced by the laser imprinting, see Chapter 5). Nonlinear effects are inherent and very important to the evolution of such broadband spectra. Signatures of nonlinear evolution of imprinted features (such as shift of the typical perturbation size toward longer wavelengths) have been shown in several experiments which measured the late-time evolution of the broadband initial spectra coming from laser imprint [17, 65, 31] or from 3-D surface perturbations using x-ray drive [30]. The results of experiments [66, 67] where three-dimensional broadband imprinted features exhibited growth that saturated at amplitudes consistent with the model of Haan are presented in Chapter 5.
2.2.5 Bubble competition and mixing phase of nonlinear RT instability

When bubble 'competition' and 'coalescence' \([55, 68, 69, 70, 71, 72, 73]\) start to dominate the spectral evolution, the RT instability evolves into a strongly nonlinear or mixing phase. The physics of the spike and bubble growth is different. The spikes are thin and massive parts of a heavier fluid which can be expected to fall inertially into light fluid with free fall rate, with the amplitude evolving as \(\xi_s \sim 0.5gt^2\). Bubbles on the other hand, have significant frontal area, so they experience a drag in the heavier fluid. In the absence of adjacent bubbles, the bubble will rise with constant velocity \(V_b = 0.23(g\lambda)^{1/2}\). However, for two adjacent bubbles with different size, the 'bubble competition' will result in the accelerated growth of a bigger bubble as the smaller bubble will gradually stop growing and will eventually be absorbed by a bigger bubble during 'bubble coalescence'. In the multimode perturbation spectrum 'bubble competition and coalescence' results in the rms amplitude of bubble growth \(b(t) = 0.035gt^2\) \([55, 72]\). As the nonlinear growth phase develops, the shear of the light fluid past the dense spikes initiates the Kelvin- Helmholtz (KH) instability, which gives rise to the vortex motion on the spikes, causing the tips to broaden into 'mushroom'-like structures. In this regime, the free fall growth of the spikes slows down. The bubbles continue to expand through this stage, competing and coalescing with each other \([55, 68, 69, 70, 73]\). The structure of the system no longer resembles the initial seed. The characteristics of the mixing phase are the KH cascade to shorter wavelength. In this stage, an isolated droplets of lighter fluid are mixed into the heavier and vice versa \([74]\).

In conditions relevant to ICF, the 'bubble competition' and 'coalescence' has not yet been clearly observed experimentally. In some experiments \([30, 65, 31, 66]\) round bubbles evolve from elongated structures with the scale length of these
features becoming larger later in time. These are nonlinear features of the RT instability. However in order to observe 'bubble competition' and 'coalescence', the bubble amplitudes have to be much further into the nonlinear regime (with their amplitudes much higher than Haan's saturation level) than in the above experiments.

2.3 Laser imprinting

In direct-drive ICF, the nonuniformities in the drive laser can create mass modulations in the target by a process called laser imprinting. As the target accelerates, these mass modulations can grow exponentially, creating large perturbations in the target shell. The understanding and control of laser imprinting is critical to successful design of a high-gain ICF target. Schematically, the imprinting process is shown in Fig. 2.5. During the first period of the laser drive (10's to 100's of picosecond), any nonuniformities in the laser can create mass modulations in the foil. Nonuniformities in the laser drive create perturbations in the ablation pressure, which leads to the perturbations in ablation velocity. Over a time-scale of \( \sim 10-100 \) ps this process creates modulations in the target thickness as schematically shown in Fig. 2.5(a). During this early-time interaction with the foil, a shock wave is launched into the target material, propagating from the front to the rear surface of the foil, compressing the target. If the front surface of the target is not uniform (because of initial or imprinted surface perturbations), the shock front will be nonuniform (or rippled). Approaching the rear side of the target, the ripples of the shock front will couple to the nonuniformities in the rear side of the target surface (this process is called a 'perturbation feed-in') which will further threaten a target stability during the deceleration phase.

As the laser light heats the target, material begins to be ablated creating a slab of hot plasma between the laser beams and the ablation surface, the point where
Figure 2.5: (a) During the first period of the laser drive (10's to 100's of picosecond), any nonuniformities in the laser can create mass modulations of the foil. (b) As soon as the laser light reaches a target, the target material starts to move away (or ablate) from it, creating a slab of hot plasma between the laser beams and the ablation surface which serves as the smoothing media for all nonuniformities in the laser drive. (c) When the shock reaches the rear surface of the foil, it starts to accelerate and the RT instability amplifies any ablation surface perturbations.
the steep temperature front meets the overdense material produced by the shock. As this plasma expands toward the laser, the critical surface (the surface with such electron density that the laser light could not penetrate it) expands with the plasma [8]. The absorption of laser energy occurs at densities lower than critical. Electron thermal conduction transports this energy toward the ablation surface. Since the temperature nonuniformities (coming from drive nonuniformities) must diffuse through the conduction region from the critical to ablation surfaces, lateral (perpendicular to the direction of the laser drive) as well as axial (along with the laser drive direction) energy transport is inevitable. This results in a reduction of the pressure nonuniformities at the ablation surface, smoothing the incident drive. In a sense, the laser nonuniformities become decoupled from the ablation surface as shown in Fig. 2.5(b). If the critical and ablation surfaces are separated by a distance $D_{ac}$, then the energy transport through electron thermal conduction reduces an imposed spatial variation in intensity $\Delta I/I$ by a factor $\alpha$ at the ablation surface:

$$\frac{\Delta P}{P}_{\text{abl}} = \alpha \frac{\Delta I}{I}_{\text{crit}},$$

(2.26)

where $P_{\text{abl}}$ is the ablation surface pressure. Classical diffusion gives [75]:

$$\alpha = \exp(-n\pi D_{ac}/\lambda_{\text{pert}}),$$

(2.27)

where $\lambda_{\text{pert}}$ is the perturbation wavelength and $n$ is a number determined by theory or simulations. Gardner [76] found that $n=2$, while according to Manheimer [77] $n=4$. Key has used numerical simulations in order to find an approximate expression for $D_{ac}(t)$ [78]:

$$D_{ac}(t) = 0.08\left(\frac{I}{10^{14} W cm^{-2}}\right)^{1/3}(\frac{\lambda}{\mu m})^{2/3}(\frac{t}{ps})\mu m,$$

(2.28)

where $I$ and $\lambda$ are the laser intensity and wavelength, respectively. The reduction of intensity nonuniformities via electron thermal conduction (see Eq.( 2.26)) is called a 'cloudy-day' effect.
As soon as the shock reaches the rear surface of the foil, it starts to accelerate as shown in Fig. 2.5(c). Any small perturbation at the unstable ablation surface will grow exponentially in the linear regime according to dispersion relation of the RT instability, threatening to destroy the target.

2.3.1 Richtmyer-Meshkov instability and ablation stabilization

As soon as the laser radiation reaches the foil, a compression shock is launched into the target, propagating towards the rear side of the foil. As the shock emerges from the rear side of the target, acceleration of the compressed foil begins. Richtmyer showed that the amplitude of any initial perturbation on the target’s surface, $a_0$, will grow with constant velocity as soon as the shock is launched into the target [79]:

$$a(t) = a_0 + \frac{1}{2} A k u_s a_0 t,$$  \hspace{1cm} (2.29)

where $A$ is the Atwood number, $k$ is the perturbation's wavenumber and $u_s$ is the shock's velocity. Such linear growth of initial perturbations due to shock propagation in the foil is known as the Richtmyer-Meshkov (RM) instability [79, 80]. According to Eq.(2.29), the amplitude of the perturbation with wavelength of 60-µm on 20-µm-thick foil will double at the time of a shock breakout (if the shock is launched at the laser intensity of $\sim 10^{13}$ W/cm$^2$, conditions typical to those described in Chapter 5), and shorter wavelength perturbations will grow even more.

In the conditions relevant to ICF, the material ablation from the unstable surface stabilizes RM instability (as in the case of RT instability). Using numerical simulations, Taylor [81] and Velikovich [82] found scaling laws relating saturation times and amplitudes to the laser drive intensity and the perturbation wavelength. According to their findings, the perturbation amplitudes of wavelengths
Figure 2.6: The ripple shock propagation of initial sinusoidal amplitude of the ripple front $a_0$ in the direction $z$. The amplitude of the shock front is $a$ after it has propagated a distance $z_s$ toward the rear surface of the target. The compressed density of the target is $\rho_s$, while the initial target density is $\rho_0$.

in the range interesting to most ICF planar experiments, 10-100 $\mu$m, do not grow significantly and even decrease their amplitudes toward the shock breakout time for 20-$\mu$m-thick CH foils at laser drive intensities $\sim 10^{14}$ W/cm$^2$.

2.3.2 Nonuniform shock propagation

The RT instability studied in this thesis exists primary at the ablation surface, the point where the steep temperature front meets the overdense material produced by the shock. Perturbations in the target result from both mass modulations (ripples at the ablation surface) and density modulations produced in the bulk of the target. The latter are created primarily by the propagation of nonuniform shocks as is shown schematically in Fig. 2.6.

Face-on, through-foil, x-ray radiography is the primary quantitative experimental technique to study evolution of 3-D target nonuniformities in planar targets (see section ‘Rayleigh-Taylor instability in ICF’ of this Chapter). For 2-D perturbations, side-on radiography results in quantitative measurements as well. Radiographic systems are sensitive to the density-thickness product (optical depth) of the target, and as such, cannot distinguish between mass and
density modulations. Even in the case of no growth of the mass modulations at the ablation surface, the measured areal density evolves in time because of the nonuniform shock propagation.

An analytical theory [83] of the ripple shock propagation in the \( z \) direction and the shock velocity \( u_s \) (see Fig. 2.6) predicts that the shock front amplitude, \( a \), satisfies following wave equation

\[
\frac{\partial^2 a}{\partial t^2} = \frac{u_s^2}{q} \frac{\partial^2 a}{\partial z^2},
\]

(2.30)

where \( q \) is a function of shock Mach number \( M \) and the specific-heat ratio \( \gamma \) [83, 48]

\[
q = \left[ \frac{M^2}{M^2 - 1} \right] \left[ 1 + \frac{2(1 - \mu^2)}{(\gamma + 1)\mu} \right] \left[ 1 + 2\mu + \frac{1}{M^2} \right],
\]

(2.31)

and \( \mu \) is given by the following expression,

\[
\mu = \sqrt{\frac{(\gamma - 1)M^2 + 2}{2\gamma M^2 - \gamma + 1}}.
\]

(2.32)

For an initial sinusoidal shock front with the wavenumber \( k \) and initial amplitude \( a_0 \cos(kx) \), the solution of Eq. (2.30) is

\[
a(t) = a_0 \cos(kx) \cos\left( k \frac{z_s(t)}{\sqrt{q}} \right),
\]

(2.33)

where \( z_s(t) = u_s t \) is the position of the shock as shown in Fig. 2.6. For a strong shock, \( M \gg 1 \) and \( \gamma = 5/3 \), the phase inversion is expected to occur at the shock propagation distance \( z_f/\lambda = 0.53 \) [48]. Equation (2.33) predicts an oscillation of the shock amplitude with the shock propagation distance.

Such behavior of the ripple shock amplitude was measured experimentally in planar CH foils by Endo et al. [48]. In this experiment, the ripple shock amplitude was measured together with the amplitude of areal-density perturbation (using face-on radiography) during shock propagation from the front to the rear surface in planar targets. At the same time, the amplitude of ablation surface
perturbation was measured using side-on radiography. The measured amplitude of the areal-density perturbation \( D \) was found to be

\[
D(t) = D_0 \left( 4 - 3a(t)/a_0 \right),
\]

where \( D_0 \) was the initial areal-density perturbation and the density ratio for strong shock compression was considered \( \rho_s/\rho_0 = 4 \). There was no observed growth of initial mass perturbation at the ablation surface from side-on radiography measurements. This experimental result was explained theoretically by Ishizaki and Nishihara [49] using numerical simulations and shock propagation models.

### 2.3.3 'Cloudy-day' model of laser imprinting

Using 1-D numerical simulations (such as the hydro-code LILAC [100]), the evolution of such foil parameters as ablation velocity, shock velocity, absorbed laser power in laser-heated planar CH foils can be found. Using these simulated parameters, Azechi et al. [84] formulated a model for laser imprint development which is called the 'cloudy-day' model. It is based on the effect of the laser nonuniformity smoothing by the electron thermal conduction ('cloudy-day' effect) and the model for the nonuniform shock propagation (both of these models were discussed earlier in this section). The 'cloudy-day' model, as presented by Azechi [84], does not include RM growth nor ablative stabilization of mass perturbations at the ablation surface. If the incident on the target laser intensity \( I_0 \) has a sinusoidal perturbation \( \delta I \), then the pressure perturbation \( \delta P \) will be proportional to the intensity perturbation, but reduced by the 'cloudy-day' effect [84]

\[
\delta P = \frac{2}{3} \frac{\delta I}{I_0} P_0 \frac{\int_{-\infty}^{+\infty} S_{abs}(z) \exp(-k|z - z_{a0}|)dz}{\int_{-\infty}^{+\infty} S_{abs}(z)dz},
\]

where \( P_0 \) is the unperturbed pressure at the ablation surface, \( S_{abs}(z) \) is the absorbed laser power per unit length, and the exponential factor accounts for the
thermal smoothing from a point of the laser absorption at \( z \) to the ablation surface at \( z_{a0} \) (the subscript 0 stands for the unperturbed quantities, see Fig. 2.7).

From the equation of motion, the time derivative of the momentum perturbation per unit surface area is equal to the pressure perturbation at the ablation surface [84]

\[
\frac{d\bar{\delta}(mv)}{dt} = \bar{\delta}P. \tag{2.36}
\]

The imposed momentum perturbation may be decomposed into one carried by the fluid moving along the axial direction (the \( z \)-direction) and the other which is lost due to the lateral fluid motion from the strongly to weakly pushed region

\[
\bar{\delta}(mv) = \bar{\delta}(mv)_{axial} + \bar{\delta}(mv)_{lateral}. \tag{2.37}
\]

Let's assume that the density is spatially uniform behind the shock \( \delta \rho_s = 0 \), and the velocity perturbation decreases exponentially with \( z \) as \( v_z(z) = v_{a0} + \delta v_a \exp(-k(z - z_a)) \), where \( \delta v_a \) is the velocity perturbation at the ablation surface. Then the term \( \bar{\delta}(mv)_{axial} \) can be expressed (neglecting higher order terms) as

\[
\delta(mv)_{axial} = \int_{z_a}^{z} \rho_s v_z dz - \int_{z_{a0}}^{z_{a0}} \rho_s v_{a0} dz = \rho_{a0} v_{a0} (\delta z_a - \delta z_a) + \rho_{a0} \delta v_a \left( \frac{1 - \exp(-k(z_{a0} - z_{a0}))}{k} \right), \tag{2.38}
\]
where $z_s$ is the position of the shock front, and $\rho_s$ and $v_a$ are the density and fluid velocity behind the shock, respectively. The time derivative of the lateral loss of the momentum perturbation is

$$\frac{d\delta(mv)_{\text{lateral}}}{dt} = \int_{z_a}^{z_s} \left( \frac{\partial \rho_s v_x v_z}{\partial x} \right) dz = -\rho_{s0} v_{a0} \int_{z_{a0}}^{z_{s0}} \left( \frac{\partial v_z}{\partial z} \right) dz,$$

(2.39)

where, using $div v = 0$ (assuming the fluid is incompressible), the mass loss term along the $x$-direction $\rho_{s0} \frac{\partial v_x}{\partial x}$ was replaced by $-\rho_{s0} \frac{\partial v_z}{\partial z}$ and high-order terms were ignored. The perturbed shock velocity $\frac{d(\delta z_s)}{dt} = \delta v_s$ divided by the unperturbed shock velocity $v_{a0}$ can be approximated

$$\frac{\delta v_s}{v_{a0}} = \frac{\delta v_a}{v_{a0}} \exp(-k(z_{s0} - z_{a0})).$$

(2.40)

Substituting Eq.( 2.38), Eq.( 2.39) into Eq.( 2.36), and using Eq.( 2.40) and Eq.( 2.37), the following differential equation for $\delta v_a(t)$ can be obtained

$$\rho_{s0} \left( \frac{1 - \exp(-k \Delta v t)}{k} \right) \frac{d(\delta v_a(t))}{dt} + 2\Delta v \exp(-k \Delta v t) \delta v_a(t) = \delta P(t),$$

(2.41)

where $\Delta v = v_{a0} - v_{a0}$. The quantities $\rho_{s0}$, $v_{a0}$, $v_{s0}$ can be numerically simulated by the 1-D hydro-codes, and the $\delta P(t)$ can be calculated from Eq.( 2.35) using simulated values. The imprint amplitude can be obtained from $\delta z_a(t) = \int \delta v_a(t') dt'$, where $\delta v_a(t')$ is the numerical solution of the differential Eq.( 2.41).

Comparing the calculated imprint from the 'cloudy-day' model and experimental measurements of the laser-induced target perturbations Azechi [84] found that the model gives reasonably good agreement with the experimental results. However, the perturbation of the target areal-density (or optical depth) was not directly measured prior to the shock breakout time (at the imprinting stage). A significant RT growth of this perturbation during the acceleration stage was necessary to increase the perturbation's amplitude to measurable magnitudes. This is a typical problem in most imprint experiments: the sensitivity of through-foil,
x-ray radiographic systems is low for high energy (more than 1 keV) x-rays and CH targets. The early-time imprinted target mass modulations are so small, that they are dominated by the system noise [17, 31, 66, 67, 84].

Using XUV radiography near 260 eV, Taylor [65] was able to measure the imprinted modulations at the time of shock breakout using a more sensitive radiographic system based on XUV multi-layer spherical mirrors. The measured imprint had a broadband spectrum down to wavelengths of order 10 μm.

If actual nonuniformity spectrum is limited to long wavelengths (> 20 μm), then the Fourier amplitudes of such spectrum must be high enough to be measured with standard high-energy x-ray radiography (as described in Chapter 4). Using spatial phase-plates, Glendinning [85, 86] was able to generate a broadband laser imprint with a minimum nonuniformity wavelength of ~ 15 μm. At the shock breakout time (prior to acceleration phase), this imprint was measured using the standard radiography technique with x-rays at ~ 1.5 keV. The advantage of these experiments was that the measurements did not require expensive and fragile XUV multi-layer optics, as in Taylor's experiments.

The laser imprint, produced by laser smoothing techniques on major ICF facilities such as NOVA, OMEGA, and NIKE, has a broadband spectrum in the spatial wavelength range down to about 2-3 μm. This scale is determined by the minimum speckle size which, in turn, is determined by the system's f-number. The σrms amplitudes of these imprinted nonuniformities are about the same as in Taylor's or Glendinning's experiments. Because the nonuniformity spectrum occupies much broader region in spatial frequency requires that the Fourier amplitudes of these modulations are much lower than for those measured in Taylor's or Glendinning's experiments. These amplitudes are so small, they cannot be detected by ~ 1 keV x-ray radiographic systems.

However, using targets with shorter x-ray attenuation length, the experimen-
tal sensitivity can be increased so much that the early-time imprint measurements become possible. For example, the use of teflon, or titanium targets, together with standard radiography at $\sim 1.3$ keV, provided enough sensitivity to measure the imprint prior to the acceleration stage even for very small imprint amplitudes under typical conditions for the OMEGA laser system. These measurements are described in Chapter 6.

Another method to measure small-amplitude broadband imprint was utilized in experiments done by Kalantar with yttrium (Y) [87, 88, 89] and germanium (Ge) [89] x-ray lasers. These measurements were done using thin (2-3 $\mu$m) Al and Si foils. Wolfrum [90] has continued similar experiments to measure the imprint efficiency with 2-D, pre-imposed perturbations in the laser drive.

2.4 Summary

This Chapter reviewed theoretical and experimental research performed in the areas of laser imprinting and its Rayleigh-Taylor evolution in ablatively accelerated targets driven by laser light. These areas are of utmost importance for direct-drive ICF. The derivation of growth-rates in linear classical RT instability has been followed by a discussion of stabilizing mechanisms under conditions relevant to ICF. The nonlinear features of the RT instability have been discussed when several nonlinear RT models were described. These models include mode-coupling model, Haan’s model for the saturation of broadband spectrum’s growth, and bubble competition and coalescence model.

The second section of this Chapter reviewed studies of imprinting physics. It has included a description of impulsive Richtmyer-Meshkov instability, model for nonuniform shock propagation, and 'cloudy-day' model of laser imprinting.
Chapter 3

Single-beam Laser Uniformity

3.1 Introduction

The experimental program at the Laboratory for Laser Energetics (LLE) supports the national inertial confinement fusion (ICF) effort by performing experiments on OMEGA [91] to investigate the requirements for attaining ignition using direct-drive targets [12]. One of the primary challenges in direct drive is the mitigation of the target perturbations created by the irradiation nonuniformities. These imprinted perturbations can be amplified by hydrodynamic instabilities to large enough amplitudes to destroy an implosion (see Chapter 5). Single-beam nonuniformities appear to be the most dangerous for direct-drive ICF and their effect can be reduced by various beam smoothing techniques. This Chapter is devoted to experimental characterization of single-beam laser uniformity improvements on the OMEGA laser system. The time integrated uniformity measurements of OMEGA single beam show beneficial effect of beam smoothing techniques such as distributed phase plates (DPP's) [16], smoothing by spectral dispersion (SSD) [15] and distributed polarization rotators (DPR's) [17].

The approach to laser smoothing with DPP's, SSD, and DPR's as smoothing techniques does not preclude efficient frequency tripling of the 1 μm light. It does not produce any significant high-intensity spikes within the laser chain as
Figure 3.1: The phase-plate intensity pattern in focus plane consists of a smooth envelope upon which is superimposed a rapidly varying structure caused by the interference between rays from different phase-plate elements.

required for the high-power glass lasers used in fusion experiments. Other beam smoothing techniques include induced spatial incoherence (ISI) [92], implemented on NIKE laser system at Naval Research Laboratory (NRL), and partially coherent light (PCL) [93], implemented on GECCO XII laser system at Institute of Laser Engineering (ILE), Japan.

3.2 Distributed phase plates (DPP’s)

Figure 3.1 schematically shows the effect of the DPP on the intensity distribution in the focus of a single beam. A DPP breaks the laser beam into beamlets whose diffraction-limited focal spot equals to the size of the irradiated target (∼1 mm). Phase plates act to distribute the phase of the incident laser beam across the beam aperture (with a scale-length of ∼1-2 mm) which produces a speckle pattern with high spatial-frequency modulations in the beam’s focal plane. Superimposed on the smooth envelope is a rapidly varying intensity structure arising from the interference between different beamlets. For example, the electric field from two rays $L_1$ and $L_2$, shown in Fig. 3.1, in the target plane is [16]

$$E = E_1 \exp \{i(kL_1 + \phi_1 - \omega t)\} + E_2 \exp \{i(kL_2 + \phi_2 - \omega t)\}, \quad (3.1)$$
where the amplitudes $E_1 = E_2$ are the diffraction limited far-field envelopes, $\phi_1$ and $\phi_2$ are the phases of the rays $L_1$ and $L_2$, respectively and $k$ is the laser wavevector. The intensity pattern in the target plane $I = |E|^2$ results in high-amplitude fluctuations

$$I = E_1^2 + E_2^2 + 2E_1E_2 \cos \left\{ k(L_1 - L_2) + (\phi_1 - \phi_2) \right\}. \quad (3.2)$$

These high-frequency modulations are smoothed by SSD when averaged over time.

### 3.3 Smoothing by spectral dispersion (SSD)

The DPP’s transform the long-wavelength beam nonuniformities (typical of coherent laser beams) to wavelengths sufficiently short that displacements of the speckle pattern (caused by SSD) smooth out most of the nonuniformities. SSD produces a time-varying frequency across the wavefront, producing multiple modes (or colors) that, as a result of imposed dispersion in the laser system, are displaced at the target plane. When averaged over the frequency-modulation time scale, these multiple patterns smooth the speckle produced by the DPP. The general principle of SSD is shown in Fig. 3.2. The two major elements of 1-D SSD are the pair of gratings and the electro-optic modulator. These elements are placed before the phase plate and frequency-tripling crystal in the OMEGA laser system. Grating 1 introduces a 'predelay' that compensates for the time delay in the beam produced by the grating 2. An electro-optic modulator introduces a bandwidth into the beam (the 'R' and 'B' labels represent the places where the 'red' and 'blue' ends of the spectrum are dominant). The grating 2 angularly disperses different frequencies of the beam, incident on the phase plate. Since the frequency across the beam varies in time, the beam intensity distribution in the focal plane varies in time. For example, with SSD the intensity distribution in the
Figure 3.2: General principle of SSD. Grating 1 introduces a 'predelay' that compensates for the time delay in the beam produced by grating 2. An electro-optic modulator introduces a bandwidth into the beam (the 'R' and 'B' labels represent the places where the 'red' and 'blue' ends of the spectrum are dominant). Grating 2 angularly disperses the different frequencies of the beam which are incident on the phase plate. Since the frequency across the beam varies in time, the beam intensity distribution in the focal plane will vary in time.
target plane in any moment of time is still can be described by the Eq. 3.2, but two rays \( L_1 \) and \( L_2 \) now have different frequencies \( \omega_1 \) and \( \omega_2 \), and wavenumbers \( k_1 \) and \( k_2 \) [16], respectively

\[
I = E_1^2 + E_2^2 + 2E_1E_2 \cos \{k_1L_1 - k_2L_2 + (\phi_1 - \phi_2) + (\omega_1 - \omega_2)t\}.
\]  \hspace{1cm} (3.3)

At any moment of time, the intensity distribution still has high-frequency modulations, but they fluctuate in time according to the frequency difference \( (\omega_1 - \omega_2) \). When averaged over time, the interference term approaches zero as \( 1/(\omega_1 - \omega_2)t \) [16], and the intensity approaches the smooth diffraction-limited envelope. 2-D SSD is implemented by using two electro-optic modulators which work at two orthogonal directions. Experimental measurements of DDP and SSD effects on single-beam uniformity are presented in the following sections.

### 3.4 Experimental configuration

For this study, a portion of the UV energy from one of OMEGA sixty laser beams was diverted to an analysis station where the beam was focused by a 10 m lens. The intensity distribution at the focal plane of that lens is a magnified image of (and equivalent to) that beam intensity distribution at the plane of the fusion target. Time-integrated images of 1-ns Gaussian pulses were captured on Kodak Aerographic Duplicating film 4421 film [94] which was processed and then digitized using a Perkin-Elmer PDS micro-densitometer [95] with a 5-\( \mu \)m aperture. The effective magnification of these equivalent target plane (ETP) images was 5.77 and the OMEGA system irradiates targets with f/6.67 focus lenses.

Figure 3.3 shows several ETP images taken with various beam smoothing techniques applied to the laser. Figure 3.3(a) shows the intensity distribution produced by a UV beam with no beam smoothing (coherent beam). The long-
Figure 3.3: Time-integrated photographs of the single OMEGA laser beam. (a) Laser beam with no smoothing (coherent beam). (b) Laser beam with DPP. (c) Laser beam with DPP and 1-D SSD, first modulator (modulation frequency 3.6 GHz). (d) Laser beam with DPP and 1-D SSD, second modulator (modulation frequency 1.2 GHz). (e) Laser beam with DPP and 2-D SSD.

Scale length (λ=30-100 μm) structure in the image is created by phase distortions produced primarily in the IR laser amplifiers. This image was acquired at a plane several millimeters away from the geometric focus of the beam in order to produce a beam about the size of an OMEGA spherical target. Figure 3.3(b) depicts the speckle pattern produced by a DPP that was inserted into the beam. Note the high-frequency speckle is significantly smaller in size than in the unsmoothed beam. The smallest speckle size is that produced by DPP elements that are separated by the original beam aperture, i.e., the diffraction limit of the focus lens, ~2-3 μm. The spatial envelope of this distribution is designed to be nearly Gaussian. A nonlinear fit to the intensity-converted DPP-only image (shown in Fig. 3.3(b)) in the super-Gaussian form of $C \exp \left\{ -(r/\sigma)^n \right\}$ resulted in $n=1.8$, whereas $n=2$ corresponds to the Gaussian envelope.

On the OMEGA laser, smoothing by spectral dispersion is applied in two-
dimensions (2-D SSD); two modulators move the beam in orthogonal directions at rates determined by their angular dispersion and modulation frequencies. Figures 3.3(c),(d) are ETP images showing the effect of each modulator separately (1-D SSD). For these experiments, modulator one (Fig. 3.3(c)) had the IR bandwidth $\Delta \lambda = 1.5$ Å, and modulator two (Fig. 3.3(d)) had $\Delta \lambda = 0.6$ Å. The modulation frequencies were 3.6 and 1.2 GHz respectively. The linear striations result from the beam movement along a particular direction, produced by SSD. This can be thought of as many speckle patterns (such as the speckle pattern in Fig. 3.3(b)) displaced along a single axis and superimposed. Figure 3.3(e) shows the effect of 2-D SSD, when both modulators were turned on. Note that little high-frequency speckle can be seen in this image.

3.5 Fourier images

The single-beam uniformity is analyzed in the spatial frequency domain. Two-dimensional Fourier power spectra of sub-regions of the ETP images in Fig. 3.3 with square area $L = 430$ μm in object plane, or 2.5 mm in image plane are shown in Fig. 3.4. The power spectrum for the beam with no smoothing (the coherent beam, see Fig. 3.4(a)) extends up to spatial frequency $30$ mm$^{-1}$ (or wavelength $\lambda = 30$ μm). The DPP transforms the long-wavelength beam nonuniformity of coherent beam to higher spatial frequencies up to $430$ mm$^{-1}$ ($\lambda = 2.3$ μm, see Fig. 3.4(b)). These wavelengths are short enough that displacements of the speckle, caused by SSD, smooth out most of the nonuniformities, when averaged over time. 1-D SSD moves the whole beam speckle pattern back and forth in a particular direction by some distance. As a result, the transforms of the 1-D SSD images (see Fig.3(c),(d)) have lines where Fourier amplitudes equal zero. From the distance between these lines the beam displacement can be calculated. The first ($\Delta \lambda = 1.5$ Å) and second ($\Delta \lambda = 0.6$ Å) modulators displace beam in
Figure 3.4: Fourier power spectra of time-integrated photographs of the single OMEGA laser beam. (a) Laser beam with no smoothing (coherent beam). (b) Laser beam with DPP. (c) Laser beam with DPP and 1-D SSD, first modulator (modulation frequency 3.6 GHz). (d) Laser beam with DPP and 1-D SSD, second modulator (modulation frequency 1.2 GHz). (e) Laser beam with DPP and 2-D SSD.
a far-field by 80 and 30 μm, respectively (see Figs. 3.4(c) and (d)). In the image with 2-D SSD (see Fig. 3.4(e)), there is no nonuniformity detected for the spatial frequencies higher then 100 mm\(^{-1}\), (λ = 10 μm) and the nonuniformity for longer wavelengths is greatly reduced.

### 3.6 Noise filtering

To quantify the laser uniformity in the images shown in Fig. 3.3, the signal has to be distinguished from noise. The noise has to be filtered out and the finite resolution of the imaging system has to be deconvolved. Such image processing is normally performed in Fourier space. The film noise and digitizing noise are two major noise sources in these images. To distinguish the signal from the noise in Fourier spectra, the noise spectra of the imaging system has been measured.

The digitization contribution to the total noise has been measured by digitizing uniform light exposures (with no film) using six different filters with transmission of 0.5, 1.1, 1.5, 1.9, 2.4 and 2.9 optical density. To measure the contribution from the film noise, the film has been exposed to uniform irradiation at six different exposure levels of 0.5, 1.1, 1.5, 1.9, 2.4 and 2.7 optical density and then digitized. The results are shown in Fig. 3.5. Digitization noise spectra are flat functions of spatial frequency (see Fig. 3.5(a)) as expected because they are not affected by the system resolution. Digitizing noise is added to the signal after the effect of the system resolution. The noise amplitudes increase when the light transmission through the filter decreases. However, the noise amplitudes (for intensity-converted images) do not increase as a square root of the number of transmitted photons, as it would be expected from statistical nature of light. This is because the densitometer is designed such that the digitization time for more dense areas (where fewer photons transmit through the film per unit time) is longer then for less dense areas of the film. Such design of the densitometer
Figure 3.5: (a) Average Fourier amplitudes as a function of spatial frequency for six different exposure levels at 0.5, 1.1, 1.5, 1.9, 2.4 and 2.9 optical density (box size $L = 430$ $\mu$m in the object plane, or 2.5 mm in the image plane) (b) The same for six different exposure levels of the film at 0.5, 1.1, 1.5, 1.9, 2.4 and 2.7 optical density.

allows efficient (fast) measurements with less noise overall.

The film noise contribution is about ten times higher then the digitization noise, as evident from the Fig. 3.5(b). This noise also depends on the exposure level. At high spatial frequencies ($f > 20$ mm$^{-1}$) it’s amplitude decreases following the densitometer’s modulation transfer function $MTF = \sin c(\pi \Delta tf_x) \sin c(\pi \Delta tf_y)$ [97], where the densitometer’s aperture width is $\Delta t = 5$ $\mu$m, and $f_x$ and $f_y$ are spatial frequencies in two perpendicular directions of the aperture). At low frequencies ($f < 20$ mm$^{-1}$), the main contribution to the noise comes from nonuniformities in a film development (a chemical distribution during film’s development is not uniform over whole area of the film). Since the MTF of the Kodak 4421 film stays constant for the range of spatial frequencies up to 100 mm$^{-1}$ [96], the film resolution did not have any effect on these measurements.

The average Fourier amplitudes of film’s optical density as a function of spatial frequency are plotted in Fig. 3.6 for three images (from Fig. 3.3): DPP only, DPP and 1-D SSD (modulator 2), and DPP and 2-D SSD. The analysis has been performed for square sub-images with box size $L = 430$ $\mu$m in the object plane,
Figure 3.6: (a) Average Fourier amplitudes as a function of spatial frequency for three images in optical density: with DPP only, with DPP and 1-D SSD (modulator 2), and with DPP and 2-D SSD, box size $L = 430 \, \mu m$ in the object plane, or 2.5 mm in the image plane. (b) The same with a vertical scale magnified.

or 2.5 mm in the image plane. Figure 3.6(a) shows spectra of signal and noise in full scale, while Fig. 3.6(b) shows low amplitude parts of the same spectra in order to display spectra at high spatial frequencies in more details.

Comparing Fig. 3.5 with Fig. 3.6 it is obvious that the main noise contribution in ETP images is the film noise. The noise spectrum line at exposure level of 1.1 (see Fig. 3.5(b)) is almost ideally fits noise level in Fig. 3.6(b). Note, that there is no signal detected above the noise for all three images at spatial frequencies $f > 70 \, \text{mm}^{-1}$, and there is no signal detected above the noise in case of DPP with 2-D SSD at spatial frequencies $f > 20 \, \text{mm}^{-1}$. Using this information, noise has been filtered out from all measured single-beam nonuniformity power per mode (see all definitions in data analysis shown in next section) by subtracting power per mode of the noise from the measured power per mode of the signal plus the noise. The resulting power per mode of a signal was corrected for system MTF, which was, in this case, the MTF of the 5-\mu m square aperture of the densitometer. The resulting nonuniformity’s power per mode for filtered and deconvolved single-beam images shown in Fig. 3.3 is shown in the next section.
3.7 Uniformity analysis

The laser speckle intensity distribution is predicted to be [98],

\[ I(\mathbf{r}) = I_0(\mathbf{r}) \{1 + S(\mathbf{r})\}, \quad (3.4) \]

where the \( I_0(\mathbf{r}) \) is the slowly-varying envelope of the intensity distribution and \( S(\mathbf{r}) \) is the speckle (or deviation from the envelope) term. The intensity images were obtained by intensity converting laser-beam photographs (Fig. 3.3) using measured film's sensitivity or D log(H) curve shown in Fig. 3.7. The speckle function \( S(\mathbf{r}) \) defined in Eq. (3.4) has been obtained by subtracting, then dividing by the envelope function \( I_0(\mathbf{r}) \) found by smoothing the image of the \( I(\mathbf{r}) \) by 70 \( \mu \text{m} \) (81 pixels). Figure 3.8 shows the lineouts of the intensity and envelope functions for the DPP-only image. The speckle function \( S(\mathbf{r}) \) does not depend on the intensity or the envelope of the beam distribution allowing different images taken at different conditions to be compared. Since deviations from the envelope function \( I_0(\mathbf{r}) \) are normalized to the envelope function, as shown from the definition of \( S(\mathbf{r}) \), its values are presented in \%.

The quantitative measure of the beam perturbations is the nonuniformity power per mode. At a given spatial frequency, \( f \), it is given by summing the
Figure 3.8: The lineouts of the intensity and envelope functions for the DPP only image.

The power spectrum of the function $S(r)$ over $2\pi$,

$$
P(f) = \sum_{\theta=0}^{2\pi} |S(|f|, \theta)|^2,
$$

(3.5)

where $P$ is the power per mode, $S(|f|, \theta)$ is complex Fourier transform of a speckle function, $|f|$ is a magnitude of the spatial frequency and $\theta$ is the angle in Fourier domain. During the summation in Eq. (3.5), the spatial frequencies with $|f| - 0.5\Delta f \leq |f| < |f| + 0.5\Delta f$ are taken to be a spatial frequency $f$, where $\Delta f = 1/L$ and $L = 430 \ \mu m$ is the analysis box size. The total sigma rms of the nonuniformity can be calculated by summing up the power per mode over spatial frequencies and taking the square root of the sum. The effect of the DPP on the beam uniformity is shown in Fig. 3.9(a), where the power per mode of speckle functions $S(f)$ (with noise filtered and MTF deconvolved) for coherent and DPP beams are presented. The effect of SSD smoothing is shown in Fig. 3.9(b), where the power per mode of the DPP, DPP with 1-D SSD (modulator 2), and DPP with 2-D SSD beams are presented. The calculated nonuniformity's sigma rms's are 45±3 % for the coherent beam and 97±12 % for the beam with DPP only. The nonuniformity sigma rms is reduced by the 1-D SSD and 2-D SSD down to 28±12 % and 8±4 %, respectively.
Figure 3.9: (a) The power per mode of speckle functions $S(f)$ versus spatial frequency for coherent (no smoothing) and with DPP laser beams. (b) The same for DPP-only, DPP with 1-D SSD (modulator 2), and DPP with 2-D SSD laser beams.

3.8 Distributed polarization rotators (DPR’s)

The beam-smoothing rate produced by SSD is proportional to the rate and amount of frequency modulation that is applied to the beam; typical smoothing times are of the order of 100-300 ps. Since significant imprinting may occur during these times, it is beneficial to produce time-independent smoothing, i.e., to instantaneously reduce the nonuniformity of the beam. Distributed polarization rotators (DPR’s) instantaneously decrease the nonuniformity of a single-beam with a DPP by $1/\sqrt{2}$ by splitting the beam into two non-interfering displaced beams with orthogonal polarizations.

An implementation of DPR’s on OMEGA is shown in Fig. 3.10. The wedge of a birefringent KDP crystal is inserted such that the beam polarization bisects the ordinary and extraordinary crystal axes. Since the element is wedged, the two orthogonally polarized beams emerge from the crystal at different angles, producing two speckle patterns that are shifted by 80 $\mu$m (see Fig. 3.10). Since the beams are orthogonally polarized the two speckle patterns do not interfere and, therefore, produce a reduction in nonuniformity by $\sqrt{2}$. The effect of the
Figure 3.10: An implementation of DPR's on OMEGA. The wedge of a birefringent KDP crystal is inserted such that the beam polarization bisects the ordinary and extraordinary crystal axes. Since the element is wedged, the two orthogonally polarized beams emerge from the crystal at different angles, producing two speckle patterns that are shifted by 80 $\mu$m.
Figure 3.11: The power per mode of speckle functions $S(f)$ versus spatial frequency for the DPP-only beam and the beam with DPP plus DPR.

DPR has been measured comparing the power per mode of the DPP-only beam and the DPP with DPR beam, (see Fig. 3.11). The nonuniformity sigma rms of 93±12 % for the DPP only beam has been reduced to 66±12 % (with is a factor of $1/\sqrt{2}$) by the DPR.

3.9 Summary

This Chapter describes measurements of uniformity improvements for a single laser beam on the OMEGA laser system. The far-field time integrated photographs of the laser beam were captured on film in an equivalent target plane (ETP) set-up. The pulse shape used in experiments was a 1-ns Gaussian. The effects of laser uniformity techniques such as distributed phase plates (DPP's), smoothing by spectral dispersion (SSD) and distributed polarization rotators (DPR's) were measured. The film noise, a dominant source of noise in these measurements, was characterized in uniformly exposed photographs. The system resolution was deconvolved and noise was filtered during the image processing. The long-scale length structure of laser nonuniformity (sigma rms~ 50 %) for a beam without any smoothing techniques has been spread over a broadband range of spatial frequencies by the DPP. As the result, the spatial features extend down
to 2-3 μm in size with nonuniformity ~100 %. The SSD smoothes out the high spatial frequencies of laser nonuniformities down to 8 % at the end of 1-ns pulse. The effect of the DPR is an instantaneous reduction of nonuniformity by $\frac{1}{\sqrt{2}}$ for spatial features with size less than ~ 80 μm. The effect of these laser uniformity improvements on a target stability will be discussed in Chapter 5.
Chapter 4

X-Ray Radiographic System

4.1 Introduction

This Chapter is devoted to an experimental characterization of the face-on, through-foil, x-ray radiographic system used to measure the evolution of perturbations in planar targets. As was mentioned in Chapter 2, face-on radiography is the most powerful tool for quantitative measurements of 3-D target perturbations. In order to reliably interpret images of target nonuniformities, the experimental system needs to be properly characterized.

In an ICF implosion, the target is hydrodynamically unstable and as a result, mass modulations in the target (either existing or created by the laser) can grow to be sufficiently large to disrupt the implosion, thereby reducing its thermonuclear yield [18]. In direct-drive ICF, nonuniformities in the drive laser can create mass modulations in the target by a process called laser imprinting. As the target accelerates, these mass modulations can grow exponentially, creating large perturbations in the target shell. The primary method of studying imprinting is through-foil x-ray radiography [27] of laser accelerated targets, where the growth of these mass modulations can be observed. Planar targets are used because they are easily diagnosed and are a reasonable approximation to the early portions of a spherical implosion.
The primary experiments use multiple laser beams to drive both the subject target and to produce x rays on another target. These x rays are filtered and imaged after they traverse the driven target. Modulations in these images are related to the optical depth (or density-thickness product) in the target. By properly interpreting these images, the character of the imprinted features and their temporal evolution are studied.

A direct measurement of the initial imprinted perturbations is difficult because of their low amplitudes (see discussion in Chapter 2). Additional complications result from the propagation effects of nonuniform shock waves [50, 48, 49]. Low-amplitude imprinting has been measured directly using an XUV laser to probe target nonuniformities produced by a laser on very thin (~2-3 \( \mu \)m) silicon and aluminum [89] or teflon targets as described in Chapter 6. The experiments described in Chapter 5 use 20-\( \mu \)m-thick CH targets, which more closely resemble the target shells normally used on OMEGA spherical implosions. These experiments are closely related to those that measure the growth of pre-imposed mass perturbations [50], which were well-simulated by hydrodynamic calculations. This provides confidence that both the energy coupling and the amount of unstable growth are well-modeled for these experiments. This provides a baseline calibration for various hydrodynamic effects that occur in the imprinting experiments. A caveat is that imprinting is not directly measured in these experiments, rather some unstable RT growth is needed to amplify the perturbations to detectable levels.

It should also be noted that the hydrodynamic instabilities studied here exist primary at the ablation surface, the point where the steep temperature front meets the overdense material produced by the shock. Perturbations in the target result from both mass modulations (ripples on the ablation surface) and density modulations produced in the bulk of the target. The latter are created primarily
by the propagation of nonuniform shocks. Radiographic systems are sensitive to the density-thickness product (optical depth) of the target, and as such, cannot distinguish between mass and density modulations. After about 1 ns of acceleration in these experiments, the variations in optical depth produced by nonuniform shocks become negligible, compared to those produced by the ablation front amplitude. After this point, it is reasonable to ascribe most of the measured optical depth to the amplitude of perturbation at the ablation surface [50].

The backlighting source typically has multiple spectral components. As a result, simulations of the resultant optical depth of the target are critical to interpreting the data. This worked extremely well for experiments using two-dimensional pre-imposed sinusoidal perturbations [44, 50]. In contrast, the features created by imprinting are three-dimensional and therefore significantly more difficult to simulate, therefore it is advantageous to obtain relationship between measured optical depth modulations and the amplitude of ablation surface modulations experimentally. This becomes possible by making several reasonable assumptions about the detection system.

In the following sections the radiographic imaging system and methods to recover the amplitude of the target perturbations from the radiographic imaged are discussed. The results of experiments are presented that characterize the sensitivity, resolution and noise of the system. Using this information, a Wiener filter has been formulated that was designed to enhance the radiographic images. In essence, the analysis provides a way to distinguish signal from the noise and deconvolve the system resolution.

4.2 Experimental configuration

The primary experiments use the following configuration. Unperturbed (smooth surface) CH ($\rho = 1.05$ g/cm$^3$) 20-μm-thick targets were irradiated at $2 \times 10^{14}$
Figure 4.1: Experimental configuration. Five overlapped beams drive a 20-μm-thick CH foil. An additional 12 beams produce x rays from uranium backlighter foil. X-rays traverse the target and are imaged by a pinhole array on a framing camera.

W/cm² in 3-ns square pulses by five overlapping UV beams (see Fig. 4.1). The targets were backlit with x rays produced by a uranium backlighter, located 9 mm away from the driven target and irradiated at $\sim 1 \times 10^{14}$ W/cm² (using 12 additional beams). X-rays transmitted through the target and a 3-μm thick Al blast shield (located at the center between the backlighter and drive foils) were imaged by 8-μm pinholes on a framing camera filtered with 6-μm of aluminum. This yielded the highest sensitivity for an average photon energy of $\sim$ 1.3 keV. The framing camera produced eight images of duration $\sim$ 80 ps, each occurring at different times. The distance between the target and the pinhole array was 2.5 cm, and the distance between the pinhole array and the framing camera was
35 cm, resulting in the magnification of \( \sim 14 \). The use of optical fiducial pulses coupled with an electronic monitor of the framing camera output produced a frame timing precision of about 70 ps. The output of the framing camera is captured on Kodak T-Max 3200 film [99] which is then digitized with Perkin-Elmer PDS microdensitometer [95] with a 20-\( \mu \)m-square scanning aperture.

Figure 4.2 shows a block diagram of the entire detection system which is comprised of four major parts: 8-\( \mu \)m pinhole, the framing camera with a microchannel plate (MCP) and phosphor plate, the film, and the digitization process. At each stage of the measurement, noise is added to the signal, and the signal plus noise are convolved with the point spread function (PSF) of each component of the system. In the frequency domain, the spectra of both the signal and the noise are multiplied by the modulation transfer function (MTF) of that subset of the imaging system.

In radiography, x-rays with a nominally wide spectrum are attenuated expo-
Figure 4.3: Uranium spectrum (solid line) and instrumental response (dashed line) as a function of x-ray energy.

Figure 4.4: X-ray spectrum propagated through 3-μm Al blast shield, 20 μm CH target, and 6-μm Al filter on MCP, then absorbed and converted into electrons by the MCP.

Eventually by the target being probed. In addition to the target, the filters and imaging devices affect the transmission of x rays to the detector. Figure 4.3 (solid line) shows a backlighter uranium spectrum used for imaging [27]. The spectral response function of the imaging system (Fig. 4.3 dashed line) includes the transmission of aluminum filters and mass absorption rate of a gold photocathode on the microchannel plate (MCP) in the framing camera. Figure 4.4 shows the spectrum used for imaging: absorbed and converted into electrons by the MCP. It is obtained by multiplying the two curves in Fig. 4.3 together and taking the attenuation of the 20-μm-thick CH target into account.
The output of the framing camera is proportional to the convolution of the x-ray spectral intensity incident on a target, its attenuation factor, and the PSF's of the pinhole $R_1(r, E, t)$ and the framing camera $R_2(r, E, t)$ including filters, where $E$ represents the x-ray energy. Assuming that no saturation occurs in these devices, the output intensity of the framing camera incident on the film is

$$I_2(r, t) \sim \int dE \int dr' R_{1,2}(r - r', E, t) f_{Al}(E, t) \mu_{Au}(E) S_{bkl}(r', E, t) \times \exp \left\{- \int_0^{z_0} dz' \mu_{CH}(E, t) \rho(r', z', t) \right\} \times \exp \left\{- \mu_{CH}(E, t) \rho_{abl}(t) \xi(r', t) \right\}. \quad (4.1)$$

In this equation, x-rays propagate along the target normal which is oriented along the $z$ axis. The coordinate $r$ is the position vector perpendicular to that axis. $I_2(r, t)$ is the output intensity of the framing camera. $R_{1,2}(r, E, t)$ is a point spread function of a pinhole and framing cameras, which is, in general, a function of the x-ray energy $E$. The aluminum filter transmission is $f_{Al}(E, t)$, the mass absorption rate of the gold photocathode (in the MCP) is $\mu_{Au}(E, t)$. $S_{bkl}(r, E, t)$ is a backlighter spectral intensity. The target density and thickness are $\rho(r, z, t)$ and $z_0(t)$, respectively. The target density and the amplitude of the target thickness modulation at the ablation surface are $\rho_{abl}(t)$ and $\xi(r, t)$. The mass absorption rate of $CH$ target is $\mu_{CH}(E, t)$.

The film converts the incident light intensity $I_2(r, t)$ into optical density $O_3(r, t)$ according to it's sensitivity or D log(H) curve W. Convolving that with PSF of the film $R_3(r)$ yields

$$O_3(r, t) = \int dr' R_3(r - r') W \left\{ \log_{10} \left( \int_{t - \tau/2}^{t + \tau/2} dt' I_2(r', t') \right) \right\} \quad (4.2)$$

where $\tau = 80$ ps is a time resolution of the framing camera. During film digitization, the optical density $O_3(r, t)$ is convolved with PSF $R_4(r)$ of the 20-μm-square aperture in the densitometer to give the digitized or measured optical density

$$O_4(r, t) = \int dr' R_4(r - r') O_3(r'). \quad (4.3)$$
The measured optical density of the film \(O_4(\mathbf{r}, t)\) is converted to intensity using the inverse film sensitivity \(W^{-1}\). The measured optical depth \(D_5(\mathbf{r}, t)\) of the target is obtained by taking the natural logarithm of that intensity-converted image,

\[
D_5(\mathbf{r}, t) = \ln(10^{W^{-1}(O_4(\mathbf{r}, t))}).
\]  

(4.4)

The primary objective of these experiments is to recover the amplitude of the perturbation at the ablation surface using the measured optical depth modulations. To do this rigorously requires significant effort. However, there are several aspects of the imaging system which enable assumptions that greatly simplify the analysis of the radiographic images. First, as a result of Al filters, a relatively narrow band \((\Delta E \simeq 200 \text{ eV})\) of x rays around 1.3 keV is used for radiography. The effect of the spectral component of uranium M-band emission around 3.5 keV (see Fig. 4.4) on the system sensitivity and resolution was measured and calculated to be insignificant. Second, the backlighter spectrum and filter transmission remain constant in time during the measurement. Third, the backlighter is produced by 12 beams that have phase plates, resulting in a very uniform and predictable spatial backlighter shape. Fourth, there is little heating of the solid part of the target (the mass absorption coefficient \(\mu\) is constant in time). Fifth, the amplitudes of the growing imprinted features are large enough that the propagation of a nonuniform shock has little contribution to the total optical depth of the target [50]. Given these assumptions, equation (1) becomes:

\[
I_2(\mathbf{r}, t) \sim I_{env}(\mathbf{r}, t) \int d\mathbf{r}' R_{1,2}(\mathbf{r} - \mathbf{r}') \exp \left\{ -D_{\xi_0}(\mathbf{r}', t) \right\}
\]  

(4.5)

where the modulation in a target optical depth \(D_{\xi_0}(\mathbf{r}, t)\) is simply:

\[
D_{\xi_0}(\mathbf{r}, t) = \frac{\xi(\mathbf{r}, t)}{\lambda_{CH}}
\]  

(4.6)
and the spectrally weighed attenuation length of the target $\lambda_{CH}$ is given as:

$$\lambda_{CH} \approx 1/(\mu_{CH}(1.3keV)\rho_{abt}).$$

(4.7)

$I_{ens}(r, t)$ is a slowly-varying envelope of the backlighter.

At this point, the target optical depth can be obtained from the measured optical depth by rigorously working backwards, compensating for noise and system response (PSF) at each stage. However, if the modulation in the target optical density $D_{\xi_0}(r, t)$ is small,

$$D_{\xi_0}(r, t) \ll 1$$

(4.8)

(which is the case in all our experiments), it is possible to consider the entire imaging system to be linear. This greatly simplifies the relation between the measured optical depth and the target optical depth. Let’s introduce a new variable: the optical depth modulation in the output of the framing camera $D_{\xi_2}(r, t)$ through the following equation

$$I_2(r, t) \sim I_{ens}(r, t) \exp(-D_{\xi_2}(r, t))$$

(4.9)

Then, assuming that $D_{\xi_0}(r, t)$ and $D_{\xi_2}(r, t)$ are small, let’s expand the exponential functions in Eq. (4.5) and Eq. (4.9) into Taylor series and, retaining only 0-th and 1-st orders in these expansions, it results in

$$D_{\xi_2}(r, t) \cong \int dr' R_{1,2}(r - r', t) D_{\xi_0}(r', t).$$

(4.10)

Here the fact that the point spread function $R_{1,2}(r, t)$ is normalized $\int dr R_{1,2}(r, t) = 1$ is used. The T-MAX 3200 film has a constant MTF up to a spatial frequency $\sim 50$ $\text{mm}^{-1}$, the highest spatial frequency we consider our in the measurements [99], so the PSF of the film is set to a $\delta(r)$-function. Since the 'linear' part of the $D \log(H)$ curve is used, the modulations in measured optical depth $D_{\xi_5}(r, t)$ are linearly related to the optical depth modulation in the target $D_{\xi_0}(r, t)$

$$D_{\xi_5}(r) = \int dr' R_{sys}(r - r') D_{\xi_0}(r', t),$$

(4.11)
where $R_{sys}(r)$ is the PSF of the entire system. It is the convolution of PSF's of the pinhole, the framing camera, and the digitizing aperture of the densitometer. In frequency space, the system MTF is the product of the MTF's of each of these components.

In summary, several approximations to the system performance have been used that enable a straightforward relationship between measured optical depth and the modulation of the ablation surface to be developed. As opposed to requiring detailed computer simulations to interpret experimental results, it is possible to find, for a class of experiments, a direct relationship between the measurement and target perturbations. Equation (4.11) has been derived assuming that modulations of the target optical depth are small compared to unity. Since Eq. (4.11) is a linear approximation, it does not treat the generation of harmonics and coupling of modes produced by nonlinearities in the system. However, by simulating these nonlinearities for modulation amplitudes even greater than those measured routinely in primary experiments, it was found that nonlinear effects were negligible compared to the noise in the system.

### 4.3 System sensitivity

Once the modulation in target optical depth is obtained, the perturbation amplitude in the target can be found, provided various criteria are met. Variations in optical depth are produced either by changes in the target density or target thickness. Apparent changes in optical depth can also result from changes in the x-ray spectrum or the attenuation coefficient of the target material. Several experiments have been performed to characterize the system performance.

The sensitivity of the system is defined by the spectrally weighted x-ray attenuation length $\lambda_{CH}$. This length is inversely proportional to the mass absorption coefficient and the target density (Eq. (4.7)). In practice, $\lambda_{CH}$ can be constructed
using target compression $C_p$, calculated by the 1-D hydrocode LILAC [100], and the attenuation length of undriven target $\lambda_x$

$$\lambda_{CH} = \frac{\lambda_x}{C_p}. \quad (4.12)$$

This relation can be used as long as the driven target maintains a cold value of its mass absorption coefficient. Typically, during these experiments the target temperature is far below the values which are expected to change the mass absorption coefficient.

The attenuation length $\lambda_x$ was measured in undriven 25-μm-thick $CH_2$ ($\rho = 0.92$ g/cm$^3$) targets using backlighter beams only with the same experimental configuration as for driven-target experiments shown in Fig. 4.5(a). A thin strip of $CH_2$ was mounted at the usual position of the experimental target so that the radiographic system could view x-rays that both miss and traverse the target, as shown on the image taken at some time during the 3-ns backlighter pulse (see Figure 4.5(b)). Using the backlighter spectrum, the system spectral response shown in Fig. 4.3 and taking into account the attenuation of the $CH_2$ strip, the attenuation length $\lambda_x$ was calculated using the following expression

$$\lambda_x = \frac{t}{OD_B - OD_A}, \quad (4.13)$$

where $t = 25 \mu m$ is the thickness of the 'strip' target, $OD_B - OD_A$ is the measured difference of optical depths in two regions out of $(B)$ and in $(A)$ the 'strip' regions. The calculated attenuation length for this material was 11.5 μm, and the measured 10±1 μm using same Eq.(4.13) from the measured difference of optical depths in these two regions (lines $A$ and $B$ in Fig. 4.5(c)). This value was constant for the ~1-mm backlighter spot and did not vary over the duration of the 3-ns backlighter pulse. Also, the radiographs were taken of undriven 20-μm-thick CH ($\rho = 1.05$ g/cm$^3$) targets that had pre-imposed, low-amplitude
Figure 4.5: (a) Experimental configuration for the 'strip'-target experiment is the same as for driven-target experiments. (b) Image of the undriven $CH_2$ 'strip' target and (c) lineouts of the measured optical depth are shown by lines $A$ and $B$. Images through both the $CH_2$ and open areas allow the optical depth of the target to be measured.
Figure 4.6: Images through both the $CH_2$ and open areas allow the linearity of the framing camera response to be measured. (a) The lineouts of the measured optical depth in and out of the 'strip'. (b) The difference between these two lines in optical depth as a function of distance.

$(0.5 \mu m)$ sinusoidal modulations with wavelengths of 60 and 30 $\mu m$. Using these modulations as control references, $\lambda_x$ was measured to be $10\pm2 \mu m$. These experiments showed that both backlighter spectrum and filter transmission remained constant in time during the measurements.

4.4 Linearity of the framing camera response

One of the approximations in these experiments is that the response of the microchannel plate (MCP) in the framing camera is linear. For high enough x-ray intensities incident on the MCP, the number of secondary electron leaving the photocathode is limited. This causes saturation of the MCP’s response. Such saturation is very difficult to compensate for during image processing because it’s effect is similar to the blurring by finite system resolution which does not depend on the input x-ray intensity. In general, the problem of signal recovery under such circumstances becomes mathematically incorrect. It is better to avoid complications and keep x-ray intensities at low enough levels that the framing camera response is linear.
The undriven, 'strip' target shots (see previous section) were also designed to check the linearity of framing camera response. Figure 4.6(a) shows two lineouts (in and out of the 'strip') of x-ray intensity in optical depth at the output of the framing camera, similar to the lines A and B in Fig. 4.5(c). In this case the lineout was taken over the region where the backlighter intensity was significantly varying. If the ratio remains constant at the highest optical depth then the MCP is not saturating. These two lines were chosen because they represent whole range of x-ray intensities possible in the driven target experiments. The lower (in the 'strip') line begins at the lowest x-ray intensity (in fact it is dominated by a film fog level from 0 to about 100 μm), while the upper (out of the 'strip') line ends at the highest possible intensities at optical depth ~7 (it corresponds to a film exposure of ~2.75 while the highest exposure for driven target experiments was ~2.5, which corresponds to an optical depth of ~6.3). Figure 4.6(b) shows the optical depth difference between these two lines as a function of the distance obtained by subtracting two lines in Fig. 4.6(a). If the framing camera was saturated, then the difference between two lines would fall at higher x-ray intensities (or at high optical depths). As was pointed out earlier, at low x-ray intensities the signal is dominated by the film fog level, so one should ignore all data for the low optical depth (from 0 to ~100 μm). At high x-ray intensities in the range of 100 to 350 μm (this corresponds to about two order of magnitudes in intensity) the difference between two line is about constant well within error bars (see Fig. 4.6(b)). This measurement confirms that there is no saturation in the framing camera response up to optical depth of 7.

An undriven, 'step' target was also designed to check system performance. Figure 4.7 shows a radiographic image of the 'step' target. It consists of three steps corresponding to the changes in CH thicknesses of 5 to 10 μm, 10 to 15 μm, and 15 to 20 μm, respectively. As with the 'strip' target, the 'step' target
Figure 4.7: Image of the 'step' target. It consists of three steps corresponding to the changes in CH thicknesses of 5 to 10 µm, 10 to 15 µm, and 15 to 20 µm, respectively. (a) Two upper lines represent lineouts along the direction of the step in the 5 and 10 µm CH regions. The lower line shows the difference between these two lines. (b) The same in the step 10 and 15 µm CH regions. (c) The same in the step 15 and 20 µm CH regions.
was illuminated by back lighter x rays only. The two upper lines in Fig. 4.7(a) represent lineouts in the step at the 5 and 10 μm regions along the direction of the step. The lower line on the same Fig. 4.7(a) shows the difference between these two lines. The fact that it stays at constant level confirms that there is no saturation of the framing camera response, as in the case of the 'strip' target. The same result is evident in Figs. 4.7(b) and (c), where the difference between lineouts from the other steps is also constant.

4.5 System resolution

The resolution of the system was characterized by measuring its response from a backlit image of a sharp, opaque edge (machined platinum) which was placed in place of a target as shown in Fig. 4.8(a). The image of the edge is shown in Fig. 4.8(b). The dashed line in Fig. 4.8(c) represents the shadow of the light intensity incident on the edge, the thin solid line is the measured light intensity propagated through the system (and averaged in the direction parallel to the edge). The thick solid line in Fig. 4.8(c) is the fit to experimental data assuming the system MTF as a two-Gaussian [101]

\[ M_{sys}(f) = \alpha_1 \exp \left\{ -\sigma_1 f^2 \right\} + 0.045 \exp \left\{ -\sigma_2 f^2 \right\}, \]  

(4.14)

where \( \alpha_1 = 0.955 \pm 0.002 \), \( \alpha_2 = 0.045 \pm 0.002 \), \( \sigma_1 = 14.2 \pm 1.8 \mu m \), \( \sigma_2 = 248.3 \pm 0.9 \mu m \), and \( f \) is the spatial frequency.

The MTF is essentially the product of the responses of three system components: the pinhole camera, the 20-μm digitizing aperture, and the framing camera. The former two are straightforward calculations based on geometry and spectral energy. The MTF of the framing camera was determined by measuring the camera response to a 150-μm-wide slit placed 1 mm in front of the camera backlit by x rays from the gold back lighter and filtered with 20 μm of beryllium
Figure 4.8: (a) Experimental configuration for the 'edge'-target experiment is the same as for driven-target experiments. (b) Image of the 'edge' target (Pt strip). (c) The dashed line represents the shadow of the light intensity incident on the edge. The thin solid line is the measured light intensity propagated through the system (and averaged in the direction parallel to the edge), and the thick solid line is the fit to experimental data, assuming the system MTF as a two-Gaussian.
(see Fig. 4.9(a)). This image of the slit shown in Fig. 4.9(b) was digitized with a 5-μm scanning aperture. The slit width and its proximity to a camera were sufficient to neglect any diffraction effects. That made a switch from \(~1.3\) keV x-rays of uranium backlighter to \(~2.5\) keV x-rays of gold backlighter insignificant. However, an increased photon flux (in case of the gold backlighter) incident on the framing camera, made this measurement more valuable because it enabled us to test a sensitivity of framing camera’s MTF from the photon scattering between microchannel and phosphor plates which was expected to be important at high photon fluxes.

The dashed line in Fig. 4.9(c) represents the light intensity incident on the slit, the thin solid line is the measured light intensity propagated through the system (and averaged in the direction parallel to the slit). The thick solid line is the fit to experimental data assuming the framing camera MTF is a two-Gaussian function as shown in Eq. 4.14 with \(\alpha_1 = 1.05 \pm 0.01\), \(\alpha_2 = -0.05 \pm 0.01\), \(\sigma_1 = 103.8 \pm 0.4\) μm, and \(\sigma_2 = 95.8 \pm 0.4\) μm. The measured MTF of the framing camera is shown in Fig. 4.10. This MTF is similar to that measured in other experiments performed at LLNL [102, 103] and NRL [104] with aluminized phosphor plates, which have reduced the long scale length scattering of photons and electrons between the phosphor and microchannel plates. This scattering resulted in reduction of the MCP resolution up to about 20 per cent at low spatial frequencies \(< 5\) mm\(^{-1}\). In our experiments phosphor plate was not aluminized. However, no significant reduction of the MTF at low spatial frequencies has been detected due to such scattering because of much lower levels of irradiation used in present experiments.

Figure 4.11 summarizes the various MTF’s discussed above. The thin solid line is the MTF of the entire system as determined its edge-response. The dashed line is the system MTF calculated as the product of the MTF’s of: the 8 μm pin-hole (dotted line), the framing camera (dot-dashed line), and the 20 μm digitizing
Figure 4.9: (a) Experimental configuration for the 'slit'-target experiment. (b) Image of the 'slit' target, installed in front of the MCP. (c) The dashed line represents the light intensity incident on the slit, the thin solid line is the measured light intensity propagated through the system (and averaged in the direction parallel to the slit), and the thick solid line is the fit to experimental data assuming the framing camera MTF as a two-Gaussian function.
Figure 4.10: Resolution of the framing camera.

Figure 4.11: Resolution of the system. The thin solid line is the measured MTF of the entire system. The dashed line is the system MTF calculated as the product of the MTF’s of the 8 μm pinhole (dotted line), the framing camera (dot-dashed line), and the 20 μm digitizing aperture (thick solid line).
aperture (thick solid line). These MTF's assumed a system magnification of \( \sim 14 \). It can be seen that for spatial frequencies below 70 mm\(^{-1}\) the measured system MTF (thin solid line) is in good agreement with the aggregate response of the individual components (dashed line). So in the analysis of target nonuniformity evolution (discussed in the section on Wiener filtering), the analyzed signal was considered only at frequencies less than 70 mm\(^{-1}\).

4.6 System noise

Using radiographs of strip targets (see Fig. 4.5(a)), the noise of the system was characterized. Since the strip targets were uniform with very smooth surfaces, all of the nonuniformity measured in the radiographs of these targets is noise. The primary noise sources in this system are: photon statistical noise of backlighter x-rays, noise in the microchannel (MCP), and phosphor plates, film noise, and noise produced during digitization. It is possible to determine the origin of noise based on its spectrum, since in frequency space, the signal and noise at each stage are multiplied by the MTF of that portion of the system.

Figure 4.12 depicts the azimuthally-averaged Fourier amplitudes of the optical depth for two 150 \( \mu \)m square regions through and around the strip. At high-frequencies (>200 mm\(^{-1}\)) the averaged noise is nearly constant, indicative of the noise from film and digitization. At lower spatial frequencies the noise amplitude depends on the MTF's of the pinhole camera and MCP. This suggests the dominant noise source is the photon statistics of backlighter x rays. In optical depth space, the noise amplitude is inversely proportional to the square root of the number of photons. There is more noise in the region of the strip where there are fewer x ray photons than in the region out of the strip.

This relationship between noise levels and the photon flux can be explained by the following consideration. If \( I_1 \) and \( I_2 \) are the average x-ray intensities in
Figure 4.12: System noise. The measured noise level for two portions of a radiograph through and around the 25 μm CH₂ strip target.

and out of the 'strip' regions, respectively, then the noise rms amplitudes in these regions are \( \sim \sqrt{I_1} \) and \( \sqrt{I_2} \), assuming a monochromatic x-ray spectrum. Since the signal's optical depth is the natural logarithm of its intensity, the variation of a signal from its averaged value in terms of the optical depth will be proportional to \( \frac{1}{\sqrt{I_1}} \) and \( \frac{1}{\sqrt{I_2}} \) after a series expansion of the logarithm, and retaining only the first term. It is assumed that the number of x-ray photons per pixel is greater than 1 (in this experiment it was \( \sim 200 \) in the region out of 'strip', and \( \sim 30 \) in the 'strip' region), which is necessary to justify this analysis. The fact that there is more noise in the optical depth in the 'strip' attenuated region with fewer x-ray photons (rms amplitude \( \sim \frac{1}{\sqrt{I_1}} \)) than in the region out of the 'strip' (rms amplitude \( \sim \frac{1}{\sqrt{I_2}} \)) supports the suggestion of the photon statistical nature of the noise.


4.7 Wiener filtering

Using the measured system sensitivity, resolution, and noise, the imprinted perturbations were recovered from the radiographic images. A broad spectrum of imprinted features has been generated by laser nonuniformities with spatial frequencies up to $425 \text{ mm}^{-1}$ determined by the laser system f-number. These initial imprinted nonuniformities come from the nonuniformities in drive laser beams which have distributed phase plates (DPP's) [15]. The RT instability has growth rates and saturation effects that depend upon spatial frequency. In addition, the resolution of the radiographic system begins to cutoff at spatial frequencies above $\sim 70 \text{ mm}^{-1}$. As the result, the detected signal resides in a narrow range of spatial frequencies $\sim 10$-70 mm$^{-1}$.

400-$\mu$m-square sections of radiographic images of the target have been analyzed by converting them to measured optical depth and compensating for the backlighter envelope using a fourth order, two-dimensional polynomial fit. The signal nonuniformity is expressed as the power per mode in optical depth by Fourier transforming the resulting optical depth.

A Wiener filter was developed to recover the true signal from the resulting images [105]. If $C(f)$ is the signal and noise measured by the system $C(f) = S(f) + N(f)$, then the restored signal $P(f)$ is [105]

$$P(f) = \frac{C(f)}{M_{sys}(f)} \frac{|S(f)|^2}{|S(f)|^2 + |N_{avg}(f)|^2}, \quad (4.15)$$

where $M_{sys}$ is total system MTF, $|N_{avg}(f)|^2$ is the average or Wiener noise spectrum, and $|S(f)|^2$ is the measured signal power spectrum. The average noise spectrum $|N_{avg}(f)|^2$ and system MTF have been measured as described above; the only unknown is $|S(f)|^2$, the signal power spectrum. In this technique, the signal is compared to the measured noise spectrum, and only points which are
Figure 4.13: (a) Power per mode of the noise $N_{ppm}$ (thin solid line) and power per mode of the signal plus noise $C_{ppm}$ of the driven foil image at 2 ns. (b) Power per mode of target modulations versus spatial frequency. $S_w$ is calculated using Wiener filter assuming MTF=1 (triangles), $S_{ppm}$ is calculated by subtracting the noise power per mode $N_{ppm}$ from the power per mode of the signal plus noise $C_{ppm}$ (lower solid line which agrees with triangles), $R_{ppm}$ is calculated using Wiener filter, assuming measured MTF (squares).

greater than twice the amplitude of that noise are at first considered, i.e.:

$$|S(f)|^2 = |C(f)|^2 - |N_{avg}(f)|^2, \text{ for } |C_{re}(f)| \text{ or } |C_{im}(f)| > 2|N_{avg}(f)|,$$  \hspace{2cm} (4.16)

where $|C_{re}(f)|$ and $|C_{im}(f)|$ are real and imaginary parts of the measured signal with noise. Due to the statistical nature of the noise spectrum, the signal that is less than twice the noise amplitude can be treated in three primary manners: a) rejected (i.e. set to zero), b) considered to be uniformly distributed between zero and twice the noise level, or 3) set equal to twice the noise level. These options are used to provide the uncertainties of the measured signal. At higher spatial frequencies (>70 mm$^{-1}$) the detector response is falling rapidly, so the signal-to-noise level is greatly reduced and the error bars are larger.

The thin solid line in Fig. 4.13(a) shows the power per mode of the noise $N_{ppm}$. The thick solid line in this figure represents the power per mode of signal and noise $C_{ppm}$ in a driven target image at $\sim$2 ns. These two lines are almost the same at high spatial frequencies > 80 mm$^{-1}$, suggesting that the noise dominates at these
spatial frequencies in the driven target image. However, there is a significant level of signal at lower spatial frequencies, which has to be separated from noise.

The result of the Wiener filter Eq. (4.15) is shown in Fig. 4.13(b). To demonstrate the effect of noise reduction, the system MTF was set to one (i.e., MTF=1, no resolution compensation) in Eq. (4.15); the calculated signal's power per mode $S_{ppm}^w$ is shown as the triangles in Fig. 4.13(b). The lower solid line shows the signal's power per mode obtained by simply subtracting the noise power per mode $N_{ppm}$ (the thin solid line in Fig. 4.13(a)) from the measured power per mode of signal plus noise $C_{ppm}$ (the thick solid line in Fig. 4.13(a)) $S_{ppm} = C_{ppm} - N_{ppm}$.

The agreement between these curves indicates that the noise compensation portion of the Wiener filter behaves reasonably. The upper curve in Fig. 4.13(b) depicts the resulting filtered and deconvolved power per mode $R_{ppm}$ using the proper MTF and represents the fully processed experimental results.

### 4.8 Summary

In summary, this Chapter described the experimental characterization of a through-foil, x-ray imaging system. By properly characterizing our detection system, a simplified relationship between radiographic images and the optical depth in the target has been found. Using the measured aspects of the system, the system response was considered to be linear, which was proven to be a good approximation for the experimental conditions. Experiments have been performed to measure the system sensitivity and resolution, which were constant for the duration of the experiment. Using the measured noise spectra, a Wiener filter was constructed that enabled the signal to be distinguish from noise and allowed the signal to be reconstructed to reconstructed by deconvolving system MTF. This technique is applied to the analysis of the primary experiments in the next Chapter.
Chapter 5

Measurements of the Saturation of Rayleigh-Taylor Growth of Broadband Nonuniformities

This Chapter is devoted to experimental studies of the Rayleigh-Taylor (RT) evolution of imprinting spectra. Using the OMEGA laser system, planar, 20- and 40-μm-thick CH targets have been accelerated by 351-nm laser beams utilizing different beam-smoothing techniques including distributed phase plates (DPP’s), smoothing by spectral dispersion (SSD), and distributed polarization rotators (DPR’s). These smoothing techniques were described in detail in Chapter 3. The RT evolution of 3-D broadband planar target perturbations seeded by laser nonuniformities was measured using x-ray radiography at ~ 1.3-keV. The characterization of x-ray radiography system was presented in Chapter 4. The experiments presented in this Chapter were designed to test nonlinear models of the RT instability under conditions relevant to ICF.

5.1 Introduction

In an ICF implosion, the target is hydrodynamically unstable and as a result, mass modulations in the target (either existing or created by laser imprinting) can grow sufficiently large to disrupt the implosion, reducing its thermonuclear
yield [18]. In direct-drive ICF, the nonuniformities in the drive laser can create mass modulations in the target by process called laser imprinting. The understanding and control of laser imprinting is critical to the successful design of a high-gain ICF target.

The experiments described in this Chapter use planar 20- and 40-μm-thick CH targets, which closely resemble the target shells normally used on OMEGA spherical implosions. The imprinting is not directly measured, rather some unstable RT growth is needed to amplify the perturbations to detectable levels. Therefore an understanding of the RT instability is necessary to quantify the early-time imprinting.

The linear growth of RT instability has been extensively studied in planar targets accelerated by both indirect drive (in which the laser energy is directed into a hohlraum and converted into x rays) [47] and direct drive (laser irradiation) [50, 51]. These experiments were generally done with pre-imposed two-dimensional (2-D) sinusoidal perturbations (see Chapter 2 for more review of these experiments). They were well simulated by hydrocodes, providing confidence that both the energy coupling and amount of unstable growth are well modeled. The experiments discussed in this Chapter are closely related to those that measure the growth of pre-imposed mass perturbations [50], which provide a baseline calibration for various hydrodynamic effects that occur in the imprinting experiments.

In direct-drive ICF it is important to understand the evolution of broadband initial spectra produced by laser imprinting. Nonlinear effects are inherent and very important to the evolution of such broadband spectra. Signatures of the nonlinear evolution of imprinted features have been shown in several works [17, 30, 65, 31] which use laser imprint as the initial perturbation for RT growth. However, because of complexity of nonlinear physics, only qualitative analysis of broadband
saturation was shown in those works. In recent years several models has been developed to explain nonlinear RT evolution in Fourier [1, 53, 61, 62, 55, 72] and in real space [55] (see discussions of these models in Chapter 2). Some multi-mode nonlinear behavior has been measured in indirect-drive RT experiments with pre-imposed 3-D multi-mode initial perturbations in planar targets [59, 60].

The goal of the experiments presented in this Chapter is to use x-ray radiography to measure the mass modulations that are created in planar targets as a result of imprinting and subsequent growth [17]. This Chapter presents the results of experiments where three-dimensional broadband imprinted features exhibit growth that saturates at amplitudes consistent with Haan’s model [1]. Target images taken at different times show the formation of bubbles and spikes from initially elongated structures. The long scale length effect of the target bowing, because of decreased drive intensity toward the edges of the target, has been quantified.

5.2 Experimental configuration

In these experiments, initially smooth, 20-μm and 40-μm-thick CH targets (ρ = 1.05 g/cm²) were irradiated at 2 × 10^{14} W/cm² in 3-ns square pulses by five overlapping UV beams. To enhance the on-target uniformity all five beams had distributed phase plates (DPP’s) [15]. For some shots smoothing by spectral dispersion (SSD) [16] and distributed polarization rotators (DPR’s) [17] were used. The targets were backlit with x rays produced by a uranium backlighter at ~ 1 × 10^{14} W/cm² (using 12 additional beams). X rays transmitted through the target and a 3-μm-thick aluminum blast shield were imaged by 8-μm pinholes on a framing camera filtered with 6-μm of aluminum (see Fig. 4.1). This yielded the highest sensitivity for an average photon energy of ~ 1.3 keV. The framing camera produced eight temporally displaced images of duration ~ 80 ps and magnification of 14. The use of optical fiducial pulses coupled with an electronic
monitor of the framing camera output produced a frame timing precision of \( \sim 70 \) ps. The framing camera images are captured on Kodak T-Max 3200 film which was digitized with Perkin-Elmer PDS microdensitometer with a 20-\( \mu \)m-square scanning aperture. The measured target optical depth (OD) is obtained from the natural logarithm of the intensity-converted images.

The experiments involve fourteen shots where radiographs were obtained at different times. For each shot, up to 6 images of the same area of the target (found by the method described in section 5.4) were processed with a 400 \( \mu \)m analysis box. These images were acquired during the time interval from 1.0 to 2.8 ns after the beginning of the drive. The backlighter shape removed by subtracting was a fourth-order two-dimensional envelope fit to data. The resulting images were the measured modulations of optical depth \( D_m(f) \). Using the measured system resolution, noise, and sensitivity we applied a Wiener filter to filter the noise from the images and deconvolved the system MTF in order to recover the target's areal density modulations \( D_t(f) \) [28] (see Chapter 4). The noise in these measurements is limited by photon statistics of the backlighter x rays and the system resolution is limited by a 8-\( \mu \)m pinhole [28].

At early times, the contribution of the propagation of nonuniform shocks to the areal density modulations in the bulk of a target is comparable to those from amplitude modulation at the ablation surface [50]. However, after about \( \sim 1 \) ns of drive, modulations with spatial frequencies in the region 10- 100 mm\(^{-1}\) (where measurements are performed) experience sufficient growth to dominate any density modulation produced by nonuniform shocks. Thus, at times > 1 ns, the measured modulations well represent the amplitude at the ablation surface.

Once the modulation in target optical depth is obtained as described in Chapter 4, the perturbation amplitude in the target can be found using spectrally weighted attenuation length \( \lambda_{CH} \) [28], which is inversely proportional to the mass
Figure 5.1: Fully processed image of the target optical depth (OD) perturbations captured at 2.4 ns for one of the six shots with all smoothing techniques including DPP's, SSD, and DPR's.

absorption coefficient and the target density. $\lambda_{CH}$ can be constructed using calculated by 1-D hydrocode LILAC [100] target compression $C_p$ and the measured attenuation length of undriven target $\lambda_x$

$$\lambda_{CH} = \frac{\lambda_x}{C_p}.$$  \hfill (5.1)

This relation can be used as long as the driven target maintains a cold value of its mass absorption coefficient. Typically, during our experiments the target temperature is far below the values which could change the mass absorption coefficient.

As an example, one of the fully processed images of the target optical depth at 2.4 ns is shown in Fig. 5.1 for a shot with all of the laser smoothing techniques including DPP's, SSD, and DPR's employed. The range of 3.2 OD corresponds to a target areal-density ($\rho r$) modulation of about $3.2 \times 10^{-3} \text{ g/cm}^2$. 
5.3 Haan’s model for broadband spectra

In the linear regime of the RT instability, individual modes do not interact and therefore grow exponentially at rates determined by the dispersion relation, given by [20, 21, 22]

\[
\gamma = \alpha \sqrt{\frac{k g}{1 + L_m k}} - \beta k V_a,
\]  

(5.2)

where \( \gamma \) is the instability growth rate, \( k \) is the wavenumber of the perturbed mode, \( g \) is the target acceleration, \( L_m \) is the density gradient scale length, and \( V_a \) is the ablation velocity. For CH targets \( L_m \sim 1 \mu m \) and the constants have values \( \alpha \sim 1 \), and \( \beta \sim 1.7 \). Equation (5.2) determines how the actual growth rate differs from the classical rate \( \gamma = \sqrt{kg} \) as a result of density scale length and laser ablation. For a single-mode initial perturbation, nonlinear effects cause the exponential growth of the mode to saturate at an amplitude \( \xi_k \sim 0.1 \lambda \) and to subsequently grow linearly in time [1]. Harmonics of the fundamental mode are generated by mode coupling [53] during the exponential growth (in linear phase) leading to the formation of bubbles (penetration of lighter fluid into heavier) and spikes (penetration of heavier fluid into lighter).

The evolution of 3-D broadband perturbations is more complicated. The fastest growing modes rapidly drive harmonics and coupled modes. The contribution of the mode coupling becomes comparable to the exponential growth for some of the modes even at small (in the linear regime) amplitudes. As a result, some modes grow faster than others, while others shrink and change their phase [72]. However, the average amplitude of all modes at given spatial frequency grows exponentially at a rate given by Eq. (5.2). Saturation occurs due to collective behavior of modes because adjacent modes can constructively interfere to create local structures with amplitudes much larger than those of individual modes. As these features experience saturation, the individual modes saturate
at amplitudes much less than $0.1\lambda$. After reaching saturation level, the modes grow, on average, linearly in time with a continuous transition from the linear to nonlinear stages. Haan formulated a model for the saturation of 3-D broadband spectrum and found a saturation level of the azimuthally average amplitude given by [61, 62]:

$$S(k) = \frac{2}{Lk^2},$$

where $L$ is a size of the analysis box. The $L$ dependence occurs because the individual Fourier amplitudes of the broadband features depend on the size of analysis region; whereas the rms amplitude $\sigma_{rms}$, a measure of the deviation of the function $\xi(i, j)$ from its average value $\bar{\xi}$, does not. Using the Fourier transform $\xi(k_i, k_j)$ of function $\xi(i, j)$, the rms amplitude $\sigma_{rms}$ is defined as [106]

$$\sigma_{rms} = \sqrt{\sum_{k_i, k_j = 0} |\xi(k_i, k_j)|^2 - |\xi(k_i = 0, k_j = 0)|^2}.$$

The rms amplitude is the physically measurable quantity and hence must have the same value independent of how it is derived. The number of Fourier modes decreases as the box size is reduced. The nonuniformity's sigma rms is the square root of the sum of all modes absolute values squared, as shown by Eq.(5.4), so the amplitudes of the modes must, concomitantly, increase to keep the nonuniformity's sigmarms constant.

The evolution of the average amplitude $\xi_k(t)$ after it reaches the saturation level $S(k)$ is given by [1]

$$\xi_k(t) = S(k) \left(1 + \ln \left\{ \frac{\xi_{k \exp}(t)}{S(k)} \right\} \right),$$

where $\xi_{k \exp}(t) = \xi_k(t = 0) \exp(\gamma t)$ is the exponential growth in the linear stage of instability. This is equivalent to the growth with a constant velocity $V(k)$ in the saturation regime,

$$V(k) = S(k)\gamma(k).$$
Figure 5.2: (a) Predicted Fourier amplitudes of optical depth using Haan model for an assumed spectrum of perturbations with initial flat power per mode ($\sigma_{rms} = 0.03 \, \mu m$) spectrum and $L = 400 \, \mu m$ at times 1.3, 1.4, 1.6, 1.8, 1.9, 2.0, and 2.2 ns. (b) The measured azimuthally averaged Fourier amplitudes of the optical depth modulations as a function of spatial frequency for 20-\(\mu m\)-thick foil at times 1.4, 1.8, 1.9, and 2.2 ns, and (c) the same for another shot with 20-\(\mu m\)-thick foil at times 1.6, 1.9, 2.2, and 2.4 ns. Haan’s saturation amplitude $S$ is shown by the dashed line.

The behavior predicted by this model is shown in Fig. 5.2(a) for an initial perturbation spectrum that has constant power per mode as a function of spatial frequency. This is representative of the features imprinted by irradiation nonuniformities that arise primarily from the speckle pattern produced by DPP’s and SSD [107]. The evolution of this spectrum (plotted as the average amplitudes versus spatial frequency for seven different times between $t = 1.3$ to 2.2 ns) is simply modeled by applying the growth-rate dispersion relation (Eq. (5.2)), the saturation level (Eq. (5.3)), and the evolution in the saturation regime (Eq. (5.5)) where the target acceleration was $g = 50 \, \mu m/ns^2$, ablation velocity $V_a = 2.5 \, \mu m/ns$ and the nonuniformity’s initial $\sigma_{rms} = 0.03 \, \mu m$. The amplitudes are converted to target optical depth by dividing by the measured, spectrally weighed, attenuation length $\lambda_x = 10 \pm 2 \, \mu m$ and multiplying by the simulated compression of the target (about 2 for times from 1.6 to 2.6 ns). The high-frequency modes grow most rapidly and saturate at the level given by Eq. (5.3) which is shown by the
dashed line, while the low-frequency modes grow more slowly. As a result, the mid-frequency modes experience the largest growth factors, producing a peak in the spectrum. As the evolution progresses, the mid-frequency modes begin to saturate, and that peak moves to longer wavelengths. This behavior is relatively insensitive to the initial spectrum or drive conditions; therefore, most broadband initial spectra will evolve similarly given sufficient time. Variations in growth rates of up to 50%, or the $\sigma_{\text{rms}}$ of initial spectrum up to two orders of magnitude have little effect on the predicted spectral evolution; the only requirement is that the spectrum be broadband. Experimental results which confirm this model, including Fig. 5.2(b), (c), are described in the following sections.

5.4 Image cross-correlations

The experiment involves multiple shots with 20-μm-thick CH targets with different smoothing techniques applied. For each shot, up to 6 unfiltered images of the same area of the target, found with a cross-correlation technique, were processed with a 400-μm analysis box. Figure 5.3 shows two images of the target acquired at 2.4 ns and 2.5 ns for one of the shot with all smoothing techniques employed. The same area of the target was found by aligning these images when the cross-correlation between two 400-μm-square regions is maximized. The cross-correlation function for two images with target optical depths $D_{t1}(r)$ and $D_{t2}(r)$ is given by

$$C(r) = \frac{\int dr' D_{t1}(r' + r) D_{t2}(r')}{\sqrt{\int dr D_{t1}(r)^2 \int dr D_{t2}(r)^2}}.$$  \hspace{1cm} (5.7)

If two images of the target are shifted by some distance $a$, then the maximum of the cross-correlation function $C(r)$ will be shifted by the same distance from the center of coordinates $(r = 0)$, as it can be seen from the Eq. (5.7). For example, if two images at 2.4 ns and 2.5 ns are misaligned by $a_x = 133$ μm in
Figure 5.3: Two images of the target acquired at 2.4 ns and 2.5 ns for one of the shots with all smoothing techniques employed. Two 400-μm-square regions inside square boxes are taken for the cross-correlation analysis.
Figure 5.4: The cross-correlation function between two images. (a) Two images are shifted by 133 μm in horizontal and 67 μm is vertical direction, respectively. The cross-correlation coefficient between two images is 17 %. (b) When two images are aligned, then the maximum of the cross-correlation function is located at the center of coordinates at \( r = 0 \). The cross-correlation coefficient between two images increases to 34 %.

horizontal and \( a_y = 67 \mu m \) in vertical directions, the peak of the cross-correlation function between these two images is shifted from the center of coordinates by same distances \( a_x \) and \( a_y \) as shown in Fig. 5.4(a). When one of the images is moved by the distances \( a_x \) and \( a_y \), aligning these two images, the peak of the cross-correlation function moves toward the center of coordinates as shown in Fig. 5.4(b). At the same time, the cross-correlation coefficient between these two images, which is defined as the maximum of \( C(r) \), increases from 17 % for misaligned images to 34 % for aligned images because larger areas of the target overlap for two aligned images.
Figure 5.5: (a) The cross-correlation function between two images taken at different shots. (b) The lineout through the center of coordinates of the cross-correlation function for two images taken at different shots (dotted line). The lineout through the center of coordinates of the cross-correlation function shown in Fig. 5.4(b) for 2.4 ns and 2.5 ns images taken at the same shot (solid line).

The details of target nonuniformity structure are unique and specific only to images taken at the same shot. Therefore, the cross-correlation technique should not find any cross-correlation between two images taken from different shots. The typical cross-correlation function of two images taken on two different shots is shown in Fig. 5.5(a). This function does not have any pronounced peaks, it fluctuates around zero, as shown by the lineout of this function through the center of coordinates in Fig. 5.5(b) by the dotted line. This indicates that there is no correlation between features in these two images. Different images for the same shot (such as captured at 2.4 and 2.5 ns and shown in Fig. 5.4) aligned by the cross-correlation technique are well correlated as shown by the lineout through the center of coordinates of their cross-correlation function shown in Fig. 5.5(b) by the solid line. The cross-correlation between different images has been greatly increased after they have been Wiener filtered. For example, the cross-correlation coefficient between unfiltered images of 34% has been increased up to 70% after they have been Wiener filtered. This indicates that the image processing efficiently reduced noise in these images.
The accuracy of the image alignment using such cross-correlation technique has been defined in the following way. For some particular shot, five images A, B, C, D, and E have been aligned with sixth image F by moving the peaks of all five cross-correlation functions toward their centers of coordinates. Then the cross-correlation functions of these five images A, B, C, D, and E have been calculated between each other. It was found that the peaks of all these cross-correlation functions were located not further than 1 pixel (1.67 μm) from the centers of coordinates (where the images are considered fully aligned) for all shots in the experiment. This indicates that the accuracy of the alignment can be considered not worse than 1.7 μm.

Figure 5.6 shows six fully processed sub-images (L =100 μm) of the target optical depth for one of the shots with full smoothing (DPP's, SSD, and DPR's applied) captured at 1.6, 1.9, 2.0, 2.2, 2.4, and 2.5 ns and aligned by the cross-correlation technique. All six raw (unfiltered) images are well correlated indicating that the evolution of the same features is observed. The correlation between different images has been greatly increased after they have been Wiener filtered. Figure 5.7 shows the cross-correlation coefficients C(r=0) between different Wiener filtered images. The solid line shows that the cross-correlation C(r=0) of the image at 2.5 ns with itself equals to 1 and the cross-correlation of the same image (2.5 ns) with other images decreases as a time separation between 2.5 ns image and other images increases. The dashed line in Fig. 5.7 shows the same behavior for the cross-correlation of the image at 2.2 ns with all other images. In fact, the same behavior has been observed for each time frame in all six shots, i.e. the cross-correlation between neighboring images is always higher than with more distant images. That confirms that the image processing allows an observation of the evolution of the same features of target perturbations.
Figure 5.6: Fully processed sub-images (with a box size of 100 μm) of the target optical depth captured at (a) 1.6, (b) 1.9, (c) 2.0, (d) 2.2, (e) 2.4, and (f) 2.5 ns for one of the six shots taken with laser conditions that include DPP’s, SSD, and DPR’s. The evolution in time to longer-scale structures is evident.

Figure 5.7: Cross-correlation of different time images for one of the six shots. The cross-correlations of images captured at 2.5 ns and 2.2 ns with all other images are shown by the solid line and the dashed line, respectively.
5.5 Nonlinear saturation of RT growth

The evolution of the averaged amplitudes of measured target optical depth modulation as a function of spatial frequency for two shots with planar 20-μm-thick CH targets is shown in Fig. 5.2(b) for times of 1.3 to 2.2 ns and in Fig. 5.2(c) for somewhat later times of 1.6 to 2.4 ns with full smoothing techniques employed including DPP’s, SSD, and DPR’s. To obtain the average amplitudes, the Fourier amplitudes are azimuthally averaged at each spatial frequency. One can readily see that the measured spectra are peaked and that the peak shifts toward longer wavelengths as time progresses, similar to the predicted behavior shown in Fig. 5.2(a). Moreover, the dashed line which shows the Haan saturation level is in reasonable agreement with the position of the spectral peak. Similar behavior was observed for all six shots taken under identical drive conditions with full smoothing techniques employed.

In Figs. 5.2(b) and (c) one can see that the measured growth of the amplitudes at 20-μm-wavelength is much less pronounced than that of 30-μm for two shots. To prove that observation, Figs. 5.8(a) and (b) show the evolution of the perturbation average Fourier amplitudes at 20-μm and 30-μm-wavelength, respectively for all six shots with full smoothing techniques employed. The growth of the 30-μm-wavelength perturbations is more pronounced than the 20-μm-wavelength perturbations for the time interval of 1.5 to 2.5 ns for all six shots. This is because the amplitudes at 20 μm are predicted to be already above Haan’s saturation level at 0.01 OD, while amplitudes at 30-μm experience a transition from exponential growth to the saturated growth with constant velocity at 0.022 OD.

The amplitude at 60 μm is predicted not to be saturated and therefore should grow exponentially during this time interval. To support the above observations, the growth of low-amplitude pre-imposed, 2-D, 60-μm-wavelength and 30-μm-
Figure 5.8: Evolution of the average Fourier amplitudes of the optical depth at (a) 20-μm-wavelength and (b) 30-μm-wavelength for six shots distinguished by the different symbols. Haan saturation levels are at 0.01 OD and 0.022 OD for 20-μm and 30-μm perturbations, respectively.

wavelength single-mode, sinusoidal perturbations on 20-μm-thick CH foils driven with the same irradiation conditions was measured. Targets with initial perturbation amplitudes of 0.05 and 0.125 μm at 60-μm-wavelength and 0.025 μm at 30-μm-wavelength were used. These initial amplitudes are sufficiently low that they are expected to be in the linear regime for at least 2.5 ns, yet have high enough amplitudes that mode coupling from the broadband spectrum has no effect on their evolution. Figure 5.9(a) shows that the broadband features at 60 μm (the combined data from six shots) grow at a similar rate as the pre-imposed 60-μm modulations (upper data points for two shots). Exponential fits to these data (three solid lines) indicate growth rates of $0.96 \pm 0.02$ ns$^{-1}$ and $1.02 \pm 0.02$ ns$^{-1}$ for the pre-imposed modulations and $0.91 \pm 0.05$ ns$^{-1}$ for the broadband modulations. Figure 5.9(b) shows that the broadband features at 30 μm (the combined data from six shots) experience a transition from linear to nonlinear phases at an amplitude of about 0.02 OD, which is 30 times lower than the single-mode saturation value of 0.6 OD (0.1 λ). At the same time, the two pre-imposed 30-μm modulations (upper data points) grow exponentially with growth rates of $1.45 \pm 0.02$ ns$^{-1}$ and $1.54 \pm 0.02$ ns$^{-1}$, respectively. These data clearly show the
Figure 5.9: (a) Average Fourier amplitudes of optical depth of imprinted features versus time at 60-μm wavelength for six shots (distinguished by different symbols) and the amplitude of pre-imposed 60-μm perturbations of corrugated targets with initial amplitudes of 0.05 and 0.125 μm (upper data). (b) The same for the 30-μm wavelength with the amplitudes of initial pre-imposed 30-μm perturbations of 0.025 μm (upper data). Exponential fits (solid lines) were used for the pre-imposed corrugation data and the 60-μm imprinted data, and a third-order polynomial fit was used to the imprinted data at 30 μm.

wavelength dependent saturation level.

In summary, perturbations were observed to saturate at levels in agreement with those predicted by Haan's model [1] and are much lower than the single-mode saturation levels 0.1 λ as has been shown in combined data from six shots with identical drive conditions using DPP’s, SSD, and DPR’s. This behavior was noted in both the shape of the spatial Fourier spectra and in the temporal behavior of modes at various wavelengths. In addition, it was found that the growth of perturbations from broadband spectrum in the linear regime is the same as that for the linear growth of pre-imposed 2-D perturbations, also in agreement with the Haan's model.

The relationship between the evolution of the pre-imposed modes and their coupling to the broadband spectrum needs to be clarified. The results inferred from Fig. 5.9 require that the growth of the pre-imposed modes remain in the linear regime (exponential growth) and are not affected by mode coupling. It
was experimentally observed that the absolute values of Fourier amplitudes of the broadband spectra are randomly distributed from zero to about twice the average level $2\tilde{\xi}_k(t)$ for any azimuthal lineout with wave-vector $k$ and at any time $t$. Therefore, the evolution of any mode $\xi_k(t)$ from the azimuthal lineout at wavenumber $k$ of the broadband spectrum is confined in the boundaries between zero and $2\tilde{\xi}_k(t)$. This means that all effects of the nonlinear mode coupling on any particular mode $\xi_k(t)$ from all other modes is of the order of $\tilde{\xi}_k(t)$ or less. This observation is in agreement with the theoretical work performed by Haan [53] for ablatively accelerated targets. If the amplitude of some particular mode $\xi_k(t)$, growing in the linear regime, is much higher than the average level of broadband spectrum $\xi_k(t) \gg \tilde{\xi}_k(t)$ as for pre-imposed modes from Fig. 5.9, then nonlinear effects will be only a small fraction of its amplitude $\xi_k(t)$. For example; it was shown in Chapter 2 section 2.2.4 'Mode coupling in weakly nonlinear stage of RT instability' that, to second-order accuracy and neglecting high-order terms, the evolution of the mode $\xi_k(t)$ can approximated by the following equation [53],

$$
\xi_k(t) = \xi_k^{\text{exp}}(t) + 1/2kA \left( \sum_{k'} \xi_k^{\text{exp}}(t)\xi_{k+k'}^{\text{exp}}(t) - 1/2 \sum_{k' < k} \xi_k^{\text{exp}}(t)\xi_{k-k'}^{\text{exp}}(t) \right),
$$

(5.8)

where $\xi_k^{\text{exp}}(t) = \xi(t = 0) \exp(\gamma(k)t)$ is an exponential (or first-order) amplitude of mode $k$. If the amplitude of some particular mode $\xi_k(t)$, growing in the linear regime, is much higher than the average level of broadband spectrum $\xi_k(t) \gg \tilde{\xi}_k(t)$ as for pre-imposed modes from Fig. 5.9, then nonlinear effects will be only a small fraction of its amplitude $\xi_k(t)$ (the first linear term in Eq.( 5.8) is much larger than the second nonlinear term). This means that the effect of mode coupling on the amplitudes of 60-μm and 30-μm pre-imposed modes (which are 10-20 times above the average level of broadband amplitudes, see Fig. 5.9) is small (of the order of 5-10 % of their amplitudes) and can be neglected compared with their exponential growth in the linear regime.
Figure 5.10: In the small regions (containing only one or two bubble), the material flow in the horizontal direction can be dominated by the flow from individual 3-D bubbles (closed solid lines). However, in the whole region with box size $L$, the material flow in the horizontal direction is dominated by the flow from the tips of the 2-D bubble (dotted lines).

The same concept can be expected in a simple physical picture. Figure 5.10 schematically shows an image of the driven target which has a single mode perturbation (dotted lines) and broadband features with individual 3-D bubbles shown by closed solid lines. In some small regions (containing only one or two bubbles), the material flow in the horizontal direction can be dominated by the flow from individual 3-D bubbles. The material from 3-D bubble flows in all directions, but it flows only in one horizontal direction from 2-D bubbles. However, in the whole region with box size $L$, the material flow in the horizontal direction is dominated by the flow from the tips of the 2-D bubble (dotted lines) because the overall contribution to the flow in this direction from the 3-D bubbles becomes much
smaller in whole region with box size $L$. Therefore overall effect of the broad-
band features on the evolution of a single mode is insignificant providing that the 
amplitude of the single mode is high enough. This consideration in real space is 
complimentary to that in Fourier space described in the previous paragraph.

The measured growth rates for 60-μm and 30-μm pre-imposed modulations 
are smaller by a factor of 2.5 than those measured for earlier times (from 0 to 1.2 
ns) by Knauer [50] in similar experiments but for larger amplitudes of initial pre-
imposed modulations ($a_0 = 0.5 \, \mu m$). This effect can be attributed to the finite 
foil thickness since the estimated late-time $\sigma_{rms}$ amplitude ($\sim 3 \, \mu m$), based on 
the measured spectra, is comparable to the thickness of a target ($\sim 6 \, \mu m$) and 
may reduce the growth of these perturbations. If small-scale bubbles penetrate 
the target, the mass available to feed the growth of a spike is limited, thereby 
limiting the perturbation growth [61, 108].

This effect has been studied numerically by Srebro et al. [109] using 2-D 
hydrodynamic simulations of the perturbation evolution with initial broadband 
spectrum coming from laser nonuniformities with and without 2-D pre-imposed 
60-μm-wavelength single-mode perturbations. These simulations were conducted 
for conditions corresponding to the experiment with 20-μm-thick targets irradiated 
with 5 overlapped laser beams having different smoothing techniques such as 
DPP’s and SSD, and with various pulse shapes. The evolution of 2-D pre-
imposed 60-μm-wavelength perturbation with initial amplitude of 0.05 μm has 
been studied for three cases: with smooth drive (the drive beams having no 
nonuniformities), with drive beams having nonuniformities generated by DPP’s, 
and with nonuniformities generated by DPP’s and SSD. The 2-D spectrum for 
a single-beam laser nonuniformity generated by DPP was taken with the same 
power per mode as shown in Fig. 3.9(b), and the rate of smoothing by SSD was 
taken to be 3.6 GHz.
Figure 5.11: The evolution of the amplitude of the target areal density (normalized to initial target density) of pre-imposed 60-μm-wavelength mode at these three conditions with smooth laser beams, with laser beams having DPP’s and with laser beams having DPP’s and SSD calculated using 2-D hydrodynamic simulations.

Even though the drive and smoothing conditions used in simulations were similar to those used in experiments described above in this section, the comparison between results of 2-D simulations and 3-D experiments has to be considered only qualitative because the RT evolution is similar but not identical in 2-D and in 3-D. Figure 5.11 shows the evolution of the amplitude of the target areal density (normalized to initial target density) of pre-imposed 60-μm-wavelength mode at these three conditions with smooth laser beams, with laser beams having DPP’s and with laser beams having DPP’s and SSD. [109] The pulse shape for these simulations was 2.5-ns ramp with peak intensity of $4 \times 10^{14}$ W/cm$^2$. The amplitude of pre-imposed modulations grows similarly for all three cases until $\sim 1.7$ ns. After 1.7 ns, the amplitude of the pre-imposed mode continues to grow in the case of the smooth drive, while in two other case with DPP-only drive and DPP plus SSD drive the growth is less pronounced and the amplitudes are decreasing
later in time after 2 ns of the drive. This evolution was more pronounced for less uniform drive condition, with DPP's only. This behavior was attributed the finite foil thickness since the late-time $\sigma_{rms}$ amplitudes, based on the simulated spectra, were comparable to the thickness of a target reducing the growth of all perturbations [109]. Similar results were obtained for two other pulse shapes: 3-ns flat top with intensity of $2 \times 10^{14}$ W/cm$^2$, and 1-ns flat then 2-ns ramp with peak intensity of $4 \times 10^{14}$ W/cm$^2$.

5.6 Experiments with DPP-only laser smoothing

To confirm that the saturation level and the late-time spectra depend neither on the target thickness nor the initial spectra of nonuniformity, Fig. 5.12 shows the results of two additional experiments using a laser drive with DPP's only for 20- and 40-$\mu$m foil thicknesses. Figure 5.12(a) shows the measured spectra at 1.0, 1.6, and 1.8 ns for one of the three shots performed with DPP's only on a 20-$\mu$m-thick foil. The saturation level was calculated using the simulated compression of the target (about 2.5 for 1.0 to 1.8 ns). Figure 5.12(b) shows the measured spectra for 40-$\mu$m CH foil at 2.7 and 2.8 ns for one of two shots. The saturation level was again calculated using the spectrally weighed attenuation length $\lambda_x$, (which does not change significantly going from 20- to 40-$\mu$m foils) and the predicted compression (about 2 at $\sim$2-3 ns). It can be seen that: (1) the higher amplitudes at early times (compare Fig. 5.12(a) to Fig. 5.2(b) and 5.2(c)) demonstrate that without SSD and DPR's the laser imprint is higher; (2) the finite target thickness does not affect the saturation levels since the shapes of measured spectra are similar for both thicknesses and in agreement with Haan's model; and (3) the measured spectra for DPP-only laser drive (Fig. 5.12) have slightly higher amplitudes at spatial frequencies $> 40$ mm$^{-1}$ than for the laser
Figure 5.12: Experimentally measured perturbation spectra using DPP-only drive beams. Azimuthally averaged Fourier amplitude of the optical depth modulation as a function of spatial frequency for (a) one of the three shots with 20-μm-thick foil at 1.0, 1.6, and 1.8 ns and (b) one of the two shots with 40-μm-thick foil at 2.7 and 2.8 ns. The saturation amplitude $S$ is shown by the dashed line.

drive with all smoothing (Fig. 5.2(b) and 5.2(c)), suggesting that SSD affects the shape of the initial imprint spectrum by lowering the amplitudes of high spatial frequencies [107]. Note that the resulting acceleration of the 40-μm-thick foils was about a factor of 2 lower and the shock breakout time a factor of 2 later compared to 20-μm foils. The measured noise level is also about twice as high since approximately four times fewer photons are transmitted through the 40-μm foil. As a result, the imprint was detected only near the end of the drive and with DPP-only drive. When additional smoothing by SSD and DPR's were added, the initial imprinted amplitudes of target nonuniformity were small enough, that even at at the end of the drive no signal was measured above the noise.
5.7 Variations of laser smoothing techniques on late-time spectra evolution

As it was pointed out earlier in this Chapter (section 5.3 'Haan's model for broadband spectra'), predictions of the Haan model for the shape of the late time spectrum is not very sensitive to the initial perturbation spectrum; therefore, most broad-bandwidth initial spectra evolve similarly. Variations in the initial amplitudes of the spectrum up to two orders of magnitude have little effect on the predicted spectral evolution. This prediction has been tested by varying the amplitudes of initial perturbation using different laser smoothing techniques. Figure 5.13 summarizes data from eleven shots performed with (1) DPP's only (two shots, Figs. 5.13(a), (d) and (g) for the average broadband amplitudes at wavelengths of 60 $\mu$m, 30 $\mu$m, and 20 $\mu$m, respectively), with (2) DPP's and SSD (three shots, Figs. 5.13(b), (e), and (h)), and with (3) DPP's, SSD, and DPR's (six shots, Figs. 5.13(c), (f) and, (i)) It is expected that SSD smoothes out perturbations at high spatial frequencies much more efficiently than at low spatial frequencies [107], therefore the expected initial imprinted spectra with and without SSD have not only different nonuniformity $\sigma_{rma}$'s, but also different shapes. However, according to Haan's model, later in time and after considerable RT growth, the perturbation spectra are expected to be similar but shifted in time with the evolution happening later in the case of DPP's plus SSD comparing to DPP's only.

The measured time evolution for all three cases is similar, as evident from the evolution of the amplitudes of at broadband perturbations 60-$\mu$m, 30-$\mu$m, and 20-$\mu$m wavelength. 30-$\mu$m and 20-$\mu$m wavelength perturbations saturate at different times for different smoothing conditions but at the same levels for all smoothing techniques. 60-$\mu$m perturbations grow in linear regime (exponential
Figure 5.13: (a) Average Fourier amplitudes of optical depth of imprinted features versus time at 60-μm wavelength for two shots (distinguished by different symbols) with DPP’s only. Exponential fit (solid line) indicates the growth rate of 0.70 ± 0.05 ns⁻¹. (b) The same for the three shots with DPP’s and SSD. Exponential fit (solid line) indicates the growth rate of 0.93 ± 0.05 ns⁻¹. (c) The same as in Fig. 5.9(a) with DPP’s, SSD, and DPR’s. (d) Average Fourier amplitudes of optical depth of imprinted features versus time at 30-μm wavelength for two shots with DPP’s only. The solid line is a third-order polynomial fit. (e) The same for the three shots with DPP’s and SSD. (f) The same as in Fig. 5.9(b) with DPP’s, SSD, and DPR’s. (g) Average Fourier amplitudes of optical depth of imprinted features versus time at 20-μm wavelength for two shots with DPP’s only. The solid line is a third-order polynomial fit. (h) The same for the three shots with DPP’s and SSD. (i) The same for the six shots with DPP’s, SSD, and DPR’s.
growth) with similar growth rates of $0.70 \pm 0.05\ \text{ns}^{-1}$, $0.93 \pm 0.05\ \text{ns}^{-1}$, and $0.91 \pm 0.05\ \text{ns}^{-1}$ for the cases of smoothing conditions with DPP's only, DPP's and SSD, and DPP's, SSD and DPR's, respectively. As it has been shown earlier in this Chapter, these growth rates are in very good agreement with the measured growth rates of pre-imposed 2-D sinusoidal perturbations at wavelength of 60 $\mu$m (Fig. 5.13(c), upper data points).

For the data with less laser uniformity with DPP's only (case (1)), the initial imprinted amplitudes are higher, and as a result the RT evolution is observed earlier in time (by $\sim 250$ ps earlier than in case (2), and $\sim 500$ ps earlier than in case (3)). This observation is evident for all three wavelengths of broadband perturbations at 60-$\mu$m, 30-$\mu$m, and 20-$\mu$m. Note also that the 30-$\mu$m-wavelength perturbations saturate at higher amplitude than the 20-$\mu$m perturbations as expected. In summary, the measured spectral evolution at $\sim 1$-3 ns including the saturation level is insensitive to the details of initial perturbation spectrum, which is in agreement with predictions of Haan's model.

5.8 Late-time perturbation evolution

In both Fig. 5.8 and Fig. 5.13, the amplitudes of both the 20 $\mu$m and 30 $\mu$m wavelength broadband perturbations decrease late in time (at 2.5- 2.8 ns for the case with all smoothing techniques employed). In the nonlinear regime, Haan's model predicts the growth of these perturbations to be linear in time (constant velocity growth) after they reach their saturation levels. This is in contradiction with experimental data. Figure 5.14 shows nonuniformity spectra at 2.5 ns (dotted line) and 2.8 ns (solid line) for one of the shot with all laser smoothing techniques employed including DPP's, SSD, and DPR's. It clearly shows that the amplitudes of all spatial frequencies higher than 30 mm$^{-1}$ decrease. Figure 5.15 shows images of the target at 2.5 ns (Fig. 5.15(a)) and 2.8 ns (Fig. 5.15(b)).
Figure 5.14: Experimentally measured late-time perturbation spectra. Azimuthally averaged Fourier amplitude of the optical depth modulation as a function of spatial frequency for a 20-μm-thick foil at 2.5, 2.8 ns. The saturation amplitude $S$ is shown by the dashed line.

Features in the 2.8 ns image become much more round and some close bubbles coalesce with a typical scale length of all features becoming larger. The bubble competition model [55] predicts (however in 2-D) such bubble coalescence occurs at higher bubble amplitudes much farther into the nonlinear regime. At the end of the drive (~3 ns) 1-D LILAC simulations predict an increase of the ablation velocity (by factor of two) with the target decompressing quickly. This indicates that, at this time, target falls apart and it’s opacity to backlighter x rays probably decreases. In this case, target perturbations in optical depth could decrease even though the perturbation amplitudes at the ablation surface continue to grow.

Another possible explanation for observed spectral behavior is that even earlier in the nonlinear regime (before the bubble coalescence), some short-scale length bubbles move underneath larger long-scale length bubbles. The areal density (or optical depth) measurement becomes insensitive to short-scale length bubbles because in the direction of x ray’s propagation, the short-scale length bubble becomes to be a part of the long-scale length bubble, even though physically the two bubbles are still separated.

This effect should be explored experimentally. To separate the influence of the
target's finite thickness and decompression, thicker targets (~60-100 μm) could be used together with longer laser drive pulses (6-10 ns) and higher x-ray energy backlighters (3-5 keV). In this case the evolution of perturbation amplitudes will not be affected by the finite target thickness effects even in the highly nonlinear stage. Under such conditions, both bubble competition and Haan models can be studied in yet unexplored highly nonlinear regimes.

5.9 'Bowling' of the target

The 'bowing' of the target is another effect that has been observed and quantified in these experiments. Figure 5.16 schematically explains this effect. Before the shot, CH target is attached to the massive Mylar washer. Without this washer a part of the laser energy from the laser beams, which irradiate the target at the angle to its normal, would miss a target and propagate toward opposing ports of OMEGA target chamber, threatening to damage the laser system. During the drive, a central ~1 mm part of a target accelerates, while the part attached to the washer stays undriven. Toward the end of the drive the target becomes 'stretched' (see Fig. 5.16). The experiments performed by Knauer [50] with same
Figure 5.16: Schematic of the foil setup. The CH foil is attached to the massive Mylar washer. After the drive begins, it bows (dashed contour) because the drive pressure is applied only to its central part.

drive condition on similar targets with pre-imposed 2-D corrugation showed that the wavelength of the corrugations increased up to 7-10 % toward the end of the drive. This effect was explained by the target 'bowing' and it resulted in stretching of all features on the target, as though they were magnified.

Figure 5.17 describes this effect quantitatively for one of the shots using six target images of different time frames at 1.6 ns, 1.9 ns, 2.0 ns, 2.2 ns, 2.4 ns and 2.5 ns. It was assumed (for simplicity) that the main effect of the bowing resulted in magnifying of late-time images with respect to early-time images. If so, then the correlation of magnified early-time image with late-time image has to be higher than the correlation of original images. The 'bowing' analysis has been performed using raw, un-filtered target images because Wiener filtered images were all processed with fixed box size, not allowing this box size to vary which is necessary to perform 'bowing' analysis. The dotted line in Fig. 5.17 shows the cross-correlation coefficient of the image at 2.2 ns with the image at 2.5 ns as the function of the magnification of the 2.2 ns image. The cross-correlation coefficient is maximized when the 2.2 ns image is magnified by 1.04 times. In order to maximize the same 2.2 ns image with the early time image at 1.6 ns, it
Figure 5.17: 'Bowing' of the target. The cross-correlation of the image at 2.2 ns with: the images at 2.5 ns (dotted line), at 2.4 ns (thin solid line), itself (thick solid line), at 2.0 ns (dashed-dotted line), at 1.9 ns (bid dashed line) and at 1.6 ns (dashed line) as a function of the magnification of 2.2 ns image. The smoothing techniques for this shot included DPP’s, SSD, and DPR’s.

has to be de-magnified by 0.97 times as shown by the dashed line in Fig. 5.17. The thick solid line shows cross-correlation coefficient of the image at 2.2 ns with itself. The shorter the time interval between images, the less magnification (or demagnification) is required for one of the images to maximize correlation between them as shown by all other lines in Fig. 5.17. This analysis shows that as time progresses, target images become larger with stretched distances between image features up to about 7-10 % between the 1.6 ns and 2.5 ns images. Still, it is small enough and have an insignificant effect on the perturbation evolution. For example, Fig. 5.17 shows that all features at 2.4 ns image are magnified by $\sim 2.5 \%$ with respect to earlier image at 2.2 ns. Figure 5.2(c) shows the spectral evolution for the same shot. The spectrum is peaked at $\sim 34 \text{ mm}^{-1}$ for the image at 2.2 ns, and the peak is shifted to $\sim 20 \text{ mm}^{-1}$ for the image at 2.4 ns. Such a change to longer wavelengths of the spectral peak, $\sim 40 \%$, could not be solely ascribed to 'bowing' effect which is responsible for only $\sim 2.5 \%$ of the spectral shift to longer wavelengths in these two images.
5.10 Summary

This Chapter has shown the measured evolution of 3-D broadband perturbations produced by laser imprinting in planar CH foils, accelerated by UV light. Using through-foil radiography these perturbations were observed to saturate at levels in agreement with those predicted by Haan’s model [1]. This behavior was noted in both the shape of the spatial Fourier spectra and in the temporal behavior of modes at various wavelengths. In addition, it was found that the growth of perturbations from broadband spectrum in the linear regime is the same as that for the linear growth of pre-imposed 2-D perturbations, also in agreement with the Haan’s model.
Chapter 6

High-contrast Measurements Using Teflon Foils

6.1 Introduction

This Chapter describes direct measurements of laser imprinting in planar geometry prior to the acceleration phase (before target perturbations are amplified by RT instability). In the experiments described in the previous Chapter, imprinting spectra were not directly measured, rather some unstable RT growth was needed to amplify the perturbations to detectable levels. Even in the case of the highest initial perturbation level (using DPP-only drive to create these imprinted nonuniformities), these perturbations became comparable or higher than noise level only after they had been amplified by RT instability by $\sim 10$ times.

Since the dominant source of noise in these experiments was photon (or statistical) noise of backlighter x rays, one of the directions toward measurements of much lower imprinted amplitudes would be to increase a backlighter flux by factor of $\sim 100$ (in order to decrease noise by factor of $\sim 10$ by increasing the backlighter irradiation intensity). However, in this case the noise level will not decrease by factor of $\sim 10$ because the dominant source of noise in the system will then be the film and digitizing noise (as it is evident from Fig. 4.12 in the Chapter 4).
Another approach to increase signal is to increase the target's opacity to x-rays either by (1) using a different spectral range of backlighter x-rays, or (2) by using different materials for the target, or (3) by using both of these methods. Method (1) was demonstrated by Taylor [65] by using an XUV radiographic system near 260 eV with same CH and CH$_2$ targets. The experimental foil attenuates the x rays more strongly in this spectral range (at 260 eV), which makes the experimental system more sensitive (the x-ray intensity perturbations are higher for more sensitive system at fixed amplitudes of target surface perturbations).

Method (3) was demonstrated in experiments done by Kalantar with yttrium (Y) [87, 88, 89] and germanium (Ge) [89] x-ray lasers. These measurements were done using thin (2-3 µm) Al and Si foils. These foils were chosen because they give a right opacity for x-rays generated by yttrium and germanium x-ray lasers.

Method (2) is described in this Chapter. Keeping the backlighter the same and using materials which are more attenuating than CH, such as teflon or titanium, it is possible to increase the system sensitivity to the levels at which direct measurements of imprinting are possible without using RT amplification.

6.2 Sensitivity of the teflon target and system noise

Figure 6.1 shows the attenuation length $\lambda(E)$ of CH (dotted line) as a function of x-ray's energy $E$. The attenuation length $\lambda(E)$ is the length at which the spectral intensity of incident x-rays $I_0(E)$ is attenuated by $e$ times

$$I(x, y, E) = I_0(x, y, E) \exp \left\{ -\frac{z(x, y)}{\lambda(E)} \right\}, \quad (6.1)$$

where $z(x, y)$ is the thickness of a target (it is a function of coordinates $x$ and $y$ which are perpendicular to $z$). The solid line in the same Fig. 6.1 represents the attenuation length of teflon as a function of x-ray's energy. At the energy of $\sim 1.3$
Figure 6.1: The attenuation length $\lambda(E)$ of teflon as a function of x-ray’s energy $E$ (solid line), and the attenuation length $\lambda(E)$ of CH as a function of x-ray’s energy $E$ (dotted line).

At 1.3 keV, the attenuation length of CH is $\lambda_{CH}(1.3 \text{ keV}) \simeq 10 \mu m$, while the attenuation length of teflon $\lambda_{tf}(1.3 \text{ keV}) \simeq 2 \mu m$. Because the ratio between the above attenuation lengths is about 5, the perturbation amplitude in experimentally measured optical depth (which is a natural logarithm of $I(x, y, E)$) is 5 times higher for the teflon target than for CH target when both these targets have same amplitude of the surface perturbation.

Teflon targets can be used in experiments with similar configuration to those discussed in Chapter 4, if the foil thickness has to chosen such a way that the reasonable x-ray flux will be transmitted through the foil and directed toward the detection system. Since the teflon is more attenuating than CH, this requires that the teflon target thickness be smaller than that of CH target. Figure 6.2 shows the spectra used for imaging: absorbed and converted into electrons by the MCP (similar to those in Fig. 4.4, see Chapter 4 for details). The upper solid line is obtained by assuming no target during the shot. One of the dotted
Figure 6.2: X-ray spectrum propagated through 3-μm Al blast shield, and 6-μm Al filter on MCP, then absorbed and converted into electrons by the MCP (solid line). The same but taking the attenuation of 20-μm CH (one of two similar dotted lines) and of 4-μm teflon into account (the other of two similar dotted lines).

lines is obtained by taking the attenuation of 20-μm CH into account, while the other dotted line (very similar to the previous one) is obtained by taking the attenuation of 4-μm teflon into account. Since the transmitted spectra with 20-μm CH and 4-μm teflon targets are very similar, the photon statistical noise level in the system will be similar (see Chapter 4 for details).

6.3 Experimental results

A test of utility of teflon targets was performed with two laser shots. One shot was with 5-μm-thick teflon target, the other one was with 20-μm-thick CH target, taken at the same experimental configuration as described in Chapters 4 and 5. An unperturbed (smooth surface) teflon (ρ = 2.22 g/cm³) 5-μm-thick target was irradiated by PS-26 pulse shape (see Fig. 6.3) using one OMEGA UV beam with DPP only. An unperturbed 20-μm-thick CH target was irradiated by the same
pulse shape using six OMEGA UV beams with DPP's only and with intensity six time higher than for the teflon target shot. The drive intensity for teflon target shot was chosen six time smaller than for CH target shot in order to keep a shock breakout time (~0.4 ns) in teflon foil approximately the same as in CH foil. All other details of experimental configuration were the same as described in Chapter 4.

Figure 6.4 shows two images, one with the 20-μm-thick CH target taken at 1.85 ns after the beginning of the drive (see Fig. 6.4(a)), and the other one with the 5-μm-thick teflon target taken at 0.5 ns after the beginning of the drive. In case of the CH target shot, significant RT growth by the factor of ~10 (which started at ~0.5 ns after the beginning of the drive) was needed to amplify the initial laser-seeded perturbations to levels above the noise. However, in case of the teflon target, the signal was detected above the noise (with was about the same as for CH target shot) near the shock breakout time (~0.4 ns). Imprinted perturbations as the initial seed for RT instability were not identical in the two shots because the drive intensities, number of drive beams, target materials were different in these two shots, but the differences were not expected to be large (at least not by an order of magnitude). The intent of this experiment was
Figure 6.4: (a) Image of the 20-μm-thick CH target at 1.85 ns. (b) Image of 5-μm-thick teflon target at 0.5 ns.

not to compare quantitatively laser imprinted perturbations in teflon and CH targets, but to show a possibility of the imprinting measurements with much higher contrast.

Figure 6.5 shows the measured nonuniformity powers per mode for two images shown in Fig. 6.4 with CH target at 1.85 ns (thick solid line), and with teflon target at 0.5 ns (thin solid line). One can easily see that the signal in the teflon target image has nonuniformities which are spread over a broad range of spatial frequencies (up to about ~80 mm\(^{-1}\), which is a detection limit of the experimental system), while the signal in the CH target image is peaked near ~30 mm\(^{-1}\). This observation is in agreement with experimental results presented in the previous Chapter, showing the shift of the spectrum’s peak toward longer wavelength (or shorter spatial frequencies) during nonlinear RT of the broadband spectrum evolution.
Figure 6.5: (a) The power per mode of the target nonuniformities of the 20-μm-thick CH target image at 1.85 ns. (b) the same for the teflon target image at 0.5 ns.

6.4 Summary

This Chapter presented experimental results which show that using more attenuating targets the experimental sensitivity can be increased so much that the early-time imprint measurements become possible. For example, the use of teflon targets together with a standard radiography at \( \sim 1.3 \) keV, provided enough sensitivity to measure the imprint at the beginning of an acceleration stage even for very small imprint amplitudes in the conditions typical for OMEGA laser system.
Chapter 7

Conclusions

This thesis presented results of experimental work performed in the area of direct-drive inertial confinement fusion (ICF). One of the biggest challenges in achieving ignition in the direct-drive ICF is mitigating the initial perturbations caused by laser irradiation nonuniformities (laser imprint). These can subsequently be amplified by the Rayleigh-Taylor (RT) instability to sufficiently large levels to disrupt the implosion, thereby reducing the thermonuclear yield. The primary method to reduce imprinting is to reduce irradiation nonuniformities, especially the single beam nonuniformities, the most dangerous for direct-drive ICF. For this purpose, several laser-beam smoothing techniques have been developed at the Laboratory for Laser Energetics (LLE) at the University of Rochester. The goal of this thesis was to quantify the effect of these smoothing techniques in laser uniformity measurements, then to apply these smoothing techniques in hydrodynamic stability experiments and to quantify their effect measuring the evolution including saturation of the imprinted target perturbations due to RT instability.

The effect of laser beam smoothing techniques on the single-beam laser uniformity has been quantified in time-integrated optical measurements of a single OMEGA laser beam. These techniques included distributed phase plates (DPP's), smoothing by spectral dispersion (SSD), and distributed polarization
rotators (DPR's). It was found that the long-scale length structure of laser nonuniformity (sigma rms $\sim 50\%$) for a beam without any smoothing techniques has been spread over a broadband range of spatial frequencies by the DPP. As the result, the spatial features extended down to 2-3 $\mu$m in size with nonuniformity $\sim 100\%$. The SSD smoothed out the high spatial frequencies of laser nonuniformities down to 8\% at the end of 1-ns pulse. The effect of the DPR was the instantaneous reduction of nonuniformity by $1/\sqrt{2}$ for spatial features with size less than $\sim 80\,\mu$m.

The effect of beam smoothing techniques on hydrodynamic stability of the CH targets was studied using an experimental system based on x-ray through-foil radiography with a time resolution determined by a framing camera used as a detector. The experimental system sensitivity, resolution, and noise have been characterized. Using this information, a Wiener filter has been formulated to distinguish signal from noise, to filter the noise and to deconvolve the finite system resolution from the signal, allowing the observation of the evolution of target perturbations. The noise in the system was dominated by the photon statistics of backlighter x rays, and the system resolution was limited by the pinhole camera resolution.

Using the well-characterized experimental system, the saturation of RT growth for broadband initial spectrum of target nonuniformities, generated by laser imprinting has been measured. The results of these measurements were in agreement with the Haan model [1] for broadband spectrum saturation. These measurements were the first clear demonstration of Haan's model in conditions relevant to ICF.

It was experimentally demonstrated that high-contrast teflon targets enabled measurements of the laser imprint early in time, prior to the acceleration phase.
7.1 Future work

Even though the RT saturation level of the broadband spectrum has been found in agreement with the Haan model predictions [1], it is still necessary to test Haan model further in highly nonlinear regime. It is also expected that 'bubble competition and coalescence' take place at amplitudes much higher than Haan saturation level, so both Haan model and 'bubble competition' model can be tested simultaneously. Experimentally it can be done with two different approaches. First one requires the use of high-resolution (with spatial resolution of at least 5 μm) x-ray optics such as KB-microscope or spherical crystals operating at ~1 keV of x-ray energy and face-on, through-foil radiography technique. The thickness of CH targets can be 10-20 μm and the backlighter can be uranium irradiated at ~10^{14} W/cm^2. All other conditions can be the same as for saturation experiments presented in Chapter 5. From the data presented in section 5.5 'Nonlinear saturation of RT growth' it is expected that amplitudes of short wavelengths (which are not resolved by the system based on the pinhole camera) are much higher than saturation level. The measurements of their evolution will test both Haan model and 'bubble competition' model in highly nonlinear stages.

The second approach requires the use of much thicker CH targets (60-80-μm-thick) together with backlighter operating at ~2 keV of x-ray energy (possibly M-band of the gold backlighter). The nonlinear evolution in highly nonlinear stage can be observed also with face-on, through-foil radiography using pinhole cameras in the range of 20-60 μm-wavelength of broadband perturbations. It is necessary to accelerate targets with ~6 ns laser pulses at laser intensities not lower than ~10^{14} W/cm^2. The amplitudes of initial imprinted perturbations have to be at least a factor of 10 higher than those used in experiments presented in section 5.5 'Nonlinear saturation of RT growth'. This can be achieved when
the targets are irradiated with coherent beams or beams with specially designed DPP's.

To further test the 'bubble competition' model, the evolution of 2-D perturbations in highly nonlinear regime can be studied with side-on and face-on, through-foil radiographies simultaneously. This will allow to compare two radiography methods which will help to interpret face-on radiography results with 3-D perturbations.
Bibliography


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IMAGE EVALUATION
TEST TARGET (QA-3)