Fast-Ion Spectrometry of ICF Implosions and Laser-Foil Experiments at the Omega and MTW Laser Facilities

by

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Submitted to the Department of Nuclear Science and Engineering in partial fulfillment of the requirements for the degree of

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Abstract

Fast ions generated from laser-plasma interactions (LPI) have been used to study inertial confinement fusion (ICF) implosions and laser-foil interactions. LPI, which vary in nature depending on the wavelength and intensity of the driver, generate hot electrons with temperatures ranging from tens to thousands of kilo-electron-volts. These electrons, which accelerate the ions measured in this work, can be either detrimental or essential to implosion performance depending on the ICF scheme employed.

In direct-drive hot-spot ignition, hot electrons can preheat the fuel and raise the adiabat, potentially degrading compression in the implosion. The amount of preheat depends on the hot-electron source characteristics and the time duration over which electrons can deposit energy into the fuel. This time duration is prescribed by the evolution of a sheath that surrounds the implosion and traps electrons. Fast-ion measurements have been used to develop a circuit model that describes the time decay of the sheath voltage for typical OMEGA implosions.

In the context of electron fast ignition, the produced fast ions are considered a loss channel that has been characterized for the first time. These ions have also been used as a diagnostic tool to infer the temperature of the hot electrons in fast-ignition experiments. It has also been shown that the hot-electron temperature scales with laser intensity as expected, but is enhanced by a factor of 2-3. This enhancement is possibly due to relativistic effects and leads to poor implosion performance.

Finally, fast-ion generation by ultra-intense lasers has also been studied using planar targets. The mean and maximum energies of protons and heavy ions has been measured, and it has been shown that a two-temperature hot-electron distribution affects the energies of heavy ions and protons. This work is important for advanced fusion concepts that utilize ion beams and also has applications in medicine.

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This thesis is composed of several studies across a range of ICF subfields and would not have been possible without the support of several scientists from the Laboratory for Laser Energetics (LLE) at the University of Rochester (UR). I thank Dr. Wolfgang Theobald and Dr. Gennady Fiksel for allowing me to collaborate with them and for their helpful suggestions and feedback on several manuscripts. I am grateful for many discussions with Dr. Christian Stoeckl, whose extensive knowledge of x-ray diagnostics was helpful in data analysis and interpretation. I thank Dr. Craig Sangster for insightful conversations and for ride-along opportunities. I also thank Prof. David Meyerhofer for taking time to review manuscripts several manuscripts. The development and fielding of diagnostics at LLE would not have been possible without the engineering staff at LLE. I thank Michelle Burke, Sam Roberts, Dino Mastrosimone and Chad Mileham for their assistance in experiment and diagnostic setup.

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Introduction

1.1 Nuclear Fusion

1.1.1 Historical Note

The foundations of controlled nuclear fusion were first laid by Einstein in his 1905 paper on mass-energy equivalence. His famous relation \( E = mc^2 \) suggests that in any nuclear reactions where the final mass of the products is less than the mass of the reactants, the excess mass must be released as energy. A number of important discoveries were made following the establishment of this theoretical framework. In 1929, Atkinson and Houtermans postulated that the fusion of light nuclei could release large amounts of energy. The hypothesis remained untested until 1932, when Oliphant demonstrated the fusion of light hydrogen isotopes in the laboratory and hence discovered tritium. It was also around that time that the neutron was discovered (Chadwick). Following this discovery, neutron bombardment experiments (Hahn and Strassmann, 1938) using uranium measured the first nuclear fission reactions (interpreted as fission by Frisch and Meitner in 1939), and the focus shifted over to nuclear fission. As word of these discoveries spread to the United States, Fermi and colleagues began a series of fission experiments first at Columbia University and later at the University of Chicago (1939). Fermi measured neutron production in \(^{235}\)U fission reactions and realized that chain reactions were possible and that they could release immense energy. This was the driving force behind a subsequent letter to President Roosevelt (the Einstein - Szilárd letter) urging the US to devote resources for the study of nuclear fission reactions, and warned the President of the implications of their results in the context of nuclear weapons development.

In 1941, as World War II had just begin, the US began an official program to apply research results for national defense. Then in 1942, Fermi demonstrated the first fission chain reactions in what is now known as the first nuclear reactor, the “Chicago Pile-1.” At this point, the US began an effort (the Manhattan Project) to develop these findings into a nuclear weapon which resulted in the first atomic bomb in 1945. Several nuclear reactors were built to enrich fissile material for bomb development, but nuclear reactors for civilian use were not commissioned until 1951. In the same year, Edward Teller who was deeply involved in the Manhattan project, revisited the idea of nuclear fusion in the context of weapons. Teller and Ulam worked on boosting the yield of conventional atomic bombs using nuclear fusion in addition to fission, which resulted in thermonuclear weapons.

After the development of thermonuclear weapons, nuclear fusion by inertial confinement has become a major fusion research program in the United States, with the ultimate goal to compress a small amount of deuterium-tritium (DT) fuel to the high densities and temperatures required for nuclear ignition and burn. Inertial Confinement Fusion (ICF) necessarily requires the creation of
Figure 1.1. Binding energy per nucleon (MeV/A) as a function of atomic mass (A) [Source: BNL Atomic Mass Evaluation]. For elements lighter than iron, indicated by the dotted vertical line, fusion of light nuclei can produce net energy. For elements heavier than iron, fission can result in net energy release.

high pressure plasmas ($\sim 300$ Gbar), or high-energy-density (HED) states of matter. The study of various HED physics and technology development on the path to thermonuclear ignition have significant applications in basic science and medicine as well.

1.1.2 Fusion Reactions

Nuclear fusion is the process by which light nuclei fuse to form heavier elements and simultaneously release significant amounts of energy. Nuclei are bound together by the strong nuclear force, which can be described as a potential well that contains the nucleons. The binding energy is a function of atomic mass as shown in Fig. 1.1. The maximum in this curve occurs for $^{56}_{26}$Fe, the most stable element on earth. Nuclei heavier than iron undergo fission to form lighter nuclei and release energy. Nuclei lighter than iron can undergo fusion to form a heavier nucleus and to release energy. In both cases, the produced nuclei are more tightly bound, and this larger binding energy is attained by the release of excess energy.

In order for light nuclei to fuse, the reactants must overcome the repulsive coulomb force ($F = Z_1 Z_2 q_1 q_2 / 4\pi\varepsilon_0 r_0$) associated with the positive charges of reacting nuclei. In this classical picture, the reacting nuclei require energies approaching $\sim 1$ MeV to overcome the coulomb barrier and fuse. Fortunately, quantum mechanics and the wave-particle duality of particles allows for significant numbers of reactions to occur at lower energies ($\sim 10$ keV). This is because low energy reactants have a low but finite probability of tunneling across the coulomb barrier and fusing with one another.

A number of fusion reactions are suitable for fusion energy. Ideal reacting nuclei are light ions with a low charge number ($Z$) that allows them to more readily overcome the coulomb barrier. Elemental abundance of the reactants is also desirable if fusion energy is to be a viable source. One reaction with considerable energy release is the fusion of deuterium (D) with tritium (T):

$$D + T \rightarrow n + \alpha + 17.6 \text{ MeV}.$$ 

Deuterium is attractive as a reactant because it is found in seawater with an atomic abundance of $\sim 100$ ppm relative to hydrogen, while tritium can be bred from abundance $^6 Li$ and $^7 Li$ using a closed fusion fuel cycle.² Reactions involving only deuterium are also possible, but release less energy:
The final reaction under consideration, deuterium-helium-3 (D$^3$He), is unique among those considered here as it produces only charged fusion products:

\[
D + ^3\text{He} \rightarrow p + \alpha + 18.3 \text{ MeV}
\]

The energy released in these reactions is in the form of kinetic energy of the fusion products, and conservation of momentum determines the amount of energy carried by each product. The cross sections of these reactions, which are essentially the probability for a fusion reaction to occur, are shown in Fig. 1.2. These cross sections, which are less than 1 barn (1 barn $\equiv 10^{-28}$ m$^2$), should be compared to the cross section for coulomb interactions between the ions. For a typical DT plasma, the coulomb cross section is 3800 barns, which is several orders of magnitude larger than the fusion cross section. It is for this reason that accelerating nuclei at a stationary target (i.e., beam fusion) is not viable as a net energy source. Most nuclei will scatter off one another instead of fusing and significant energy is lost to accelerating the target ions. To overcome this problem, all reactants are heated to temperatures of 10-20 keV and confined sufficiently long for fusion reactions to occur. At these temperatures, the electrons and ions of the reactants have dissociated, leaving them in a plasma state. The bulk nuclei in this plasma scatter off one another, resulting in a thermal velocity distribution of the nuclei. Nuclei in the tail of this thermal distribution have energies that are several times higher than that of the bulk (tens of keV), and hence have a high probability of quantum tunneling through the coulomb barrier and fusing. Since the plasma is confined, scattering of the bulk nuclei results in a re-distribution of energy and acts to maintain the thermal distribution. The problem of producing efficient fusion reactions is thus reduced to confinement of the plasma with minimal radiation and heat conduction losses.

A major consideration for controlled thermonuclear fusion energy is thus the confinement of the hot plasma for sufficiently long periods of time. The timescale for confinement must be larger than the timescale for fusion reactions to occur. For a thermal plasma, where the velocity distri-
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Figure 1.3. Maxwellian-averaged volumetric reactivities for the three prominent fusion reactions as a function of reactant temperature. The DD curve represents the D(d,p)T branch [Source: ENDFVII.1 library].

...bution of the reacting ions are approximately characterized with a single temperature (Maxwellian distribution of velocities), the timescales for fusion reactions are given by:

\[ \tau_f = \frac{1}{\langle \sigma v \rangle n}, \]

(1.1)

where \( n \) is the ion density of the reacting nuclei and \( \langle \sigma v \rangle \) is the velocity-averaged reactivity (see Fig. 1.3) that is defined by the integral:

\[ \langle \sigma v \rangle \equiv \int \sigma(|v_2 - v_1|)f(v_1)f(v_2)(|v_2 - v_1|)dv_1dv_2, \]

(1.2)

where \( v_1 \) and \( v_2 \) are the velocity vectors and \( f(v_1) \) and \( f(v_2) \) the velocity distribution functions of the reacting nuclei. For a collisional plasma, the velocity distribution is a Maxwellian that is characterized by an ion temperature \( T_i \) and given by:

\[ f(v) = \left( \frac{m}{2\pi k_B T_i} \right)^{3/2} \exp\left[ -\frac{mv^2}{2k_B T_i} \right]. \]

(1.3)

In this formulation the reaction rate, or the number of reactions occurring per second per unit volume is given by:

\[ R = \frac{f_1f_2}{1 + \delta_{12}} n^2 \langle \sigma v \rangle, \]

(1.4)

where \( n \) is the total ion density, \( f_1 \) and \( f_2 \) are the fuel density fractions for the reacting nuclei (e.g., \( n_1 = f_1 n \)), and the Kronecker delta is unity if the two species are equal (for like-particle reactions such as DD) and zero otherwise.

The velocity-averaged reaction rate was computed using a Maxwellian velocity distribution of the reacting nuclei. For a DT plasma at a temperature of \( \sim 15 \) keV and a number density of \( 10^{25}-10^{26} \) cm\(^{-3} \) (typical ICF conditions), the timescales for fusion reactions are of the order 0.1 - 1 ns, which must be shorter than the confinement time of reactants for an appreciable number of reactions to occur.
1.2 Inertial Confinement Fusion (ICF) Using Laser Drivers

1.2.1 Overview

Compression, heating and inertial confinement of a fuel pellet using lasers as drivers was first openly proposed in 1972. In this scheme, the spherical fuel pellet has to be compressed to large densities and high pressures for the fusion reaction timescales to be smaller than the confinement time set by the inertia. An analytical calculation illustrates the effectiveness of ICF. First, the mass confinement time ($\tau_c$) of the compressed fuel can be expressed as:

$$\tau_c = \frac{\rho R_f}{4 m_f c_s n_0},$$  (1.5)

where $\rho R_f$ is the line-integrated density (areal density) of the compressed fuel, $n_0$ the density of the fuel, $m_f$ the mass of the fuel, and $c_s$ is the ion sound speed in the fuel ($c_s \equiv [2T/m_f]^{1/2}$). The ratio of the confinement time to $\tau_f$ determines the amount of fuel that is burned during inertial confinement. In particular, the burn fraction is defined as $\phi = (n_i - n_f)/n_i$, where $n_f$ and $n_i$ are the number densities of final and initial number of DT fuel ions, respectively. Since the fuel ions are depleted as they burn, the number of ions at any given time can be written as the differential equation:

$$\frac{dn_f}{dt} = -\langle \sigma v \rangle \frac{n_i^2}{2},$$  (1.6)

which can be integrated from $t = 0$ to $t = \tau_c$ and subtracted from the number of initial ions to determine the number of fusion reactions that occur during confinement. After normalizing this quantity to the initial number of fuel ions, and using the relation $\rho = 2n_f m_f$, algebraic manipulation yields an expression for the fuel burn fraction:

$$\Phi = \frac{\rho R_f}{H_B + \rho R_f},$$  (1.7)

where $H_B (\equiv 8c_s m_f/\langle \sigma v \rangle)$ is the temperature-dependent “burn parameter” that originates from the Maxwellian velocity-averaged reactivity. For temperatures in the range of 20-100 keV, the burn parameter is relatively constant with a value of $H_B \sim 7$ g-cm$^{-2}$. Thus, for a $\sim 20$ keV thermal DT plasma an appreciable burn-fraction of $\sim 30\%$ is possible for $\rho R \sim 3$ g-cm$^{-2}$.

The burn fraction calculation assumes that the fuel is uniformly heated to tens of keV. This is impractical, since energies of $\sim 10$ MJ are required to externally heat a typical ICF plasma ($\rho R$ of 3 g-cm$^{-2}$, total mass of 1 mg) to a temperature of $\sim 20$ keV. The conventional approach is therefore to heat a small central region of the fuel (the “hot spot”) with a $\rho R$ of $\sim 0.3$ g-cm$^{-2}$ to modest temperatures $\sim 5$ keV. This will generate DT-α’s that will stop within the hot spot, increase the hot-spot temperature and set up a burn wave that will propagate outwards, heating the main part of the fuel that forms a shell surrounding the hot spot. The criteria for self-heating of the hot spot are discussed in Sec. 1.2.2, while schemes for fuel compression and hot-spot generation are discussed in Sec. 1.2.3.

Eq. 1.7 shows that for a given temperature, larger $\rho R$ values result in greater burn fraction, and hence more fusion power output. In addition, high compression of the fuel (large $\rho$) is also desirable, as demonstrated next. For a spherical fuel assembly, the fuel mass can be written as:

$$M = \frac{4}{3\pi}(\rho R)^3/\rho^2.$$  (1.8)
To ensure the integrity and longevity of the experiment or the fusion reactor, the amount of energy that can be released at once must be limited. This places an upper bound on the total mass of fusion fuel that can be burned at once. Since a high-burn fraction is desirable for maximum fusion power output, it is evident from Eq. 1.8 that a large \( \rho \) and \( \rho R \) are required for a small total fuel mass. Compression of the target is therefore a fundamental requirement to ICF.

High compression of the fuel requires that the laser compress it along a low adiabat. The adiabat in ICF is a non-dimensional parameter defined as the ratio of the plasma pressure to that of a fermi-degenerate gas \(( \alpha \equiv p/p_F )\). In relation to classical thermodynamics, it can be shown that for isentropic (constant adiabat) compression the equation of state takes on the form:

\[
p / \rho^{\gamma} = \alpha A_{\text{deg}},
\]

where \( p, \rho \) and \( \gamma \) are the pressure, density and adiabatic index of the fuel, respectively, and \( A_{\text{deg}} \) is a constant that relates the fermi-degeneracy pressure to the density.\(^3\) The amount of work required to compress the fuel to a given density is given by:

\[
W = \int p \, dV \propto \alpha \int \rho^{\gamma} dV. \tag{1.10}
\]

The energy expended by the laser to compress the target to a fixed density increases with the adiabat, and detailed simulations have shown that the maximum \( \rho R \) scales with adiabat and laser energy:\(^6\)

\[
(\rho R)_{\text{max}} = \frac{2.6}{\alpha^{0.54}} E_{\text{MJ}}^{1/3}.
\]

As shown in Eq. 1.11, the adiabat must be kept low to achieve high fuel \( \rho R \) for a given laser energy. The adiabat is prescribed by the pressure and the density of the fuel (Eq. 1.9), which is set at the start of the implosion using strong pressure waves or shocks. These shocks are launched into the fuel using shaped laser pulses, as discussed further in Sec. 1.2.3.

For these reasons, ICF targets consist of a hollow spherical shell with an inner layer of cryogenic fuel (e.g., DT-ice) and an outer ablator (of CH, CD or Be) to maximize the initial density (and minimize the adiabat). The ablator absorbs energy from the driver and vaporizes. The outward expansion of this layer drives the remaining shell mass inwards, which compresses the fuel. Compression and subsequent hot-spot formation of ICF targets is described in greater detail in Sec. 1.2.3.

### 1.2.2 Criteria for Self-heating and Ignition

The central hot spot is formed after the fuel is compressed to its peak \( \rho R \). A self-sustaining burn-wave is initiated in the hot spot that propagates outward, consuming the dense fuel shell and generating high burn fractions. The hot spot maintains its temperature by self-heating from charged-fusion products (e.g., DT-\( \alpha \)'s). For net self-heating to occur, the power per unit volume deposited by the DT-\( \alpha \)'s into the hot spot \( (W_\alpha) \) must be greater than the power density loss from the hot spot due to heat conduction \( (W_c) \), radiation \( (W_r) \) and mechanical work \( (W_m) \). When the hot spot is first formed, it must satisfy the following condition:

\[
W_\alpha > W_r + W_c + W_m. \tag{1.12}
\]

Several authors have evaluated each of these terms to develop the criterion for hot-spot self-heating.\(^3,7\) We summarize the salient results here. The energy deposited by \( \alpha \)-particles (for DT) per
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second per unit volume is given by:

\[ W_\alpha = \frac{f_1 f_2}{1 + \delta_{12}} \left( \frac{\rho_h}{m_i} \right)^2 f_\alpha E_\alpha \langle \sigma v \rangle, \]  

(1.13)

where \( E_\alpha \) is the energy of an \( \alpha \)-particle, \( m_i \) is the effective mass (\( \equiv \sqrt{m_1 m_2} \)) of the fuel ions, and \( f_\alpha \) is the fraction of the \( \alpha \)-particle energy that is deposited into the hot spot. The energy deposited is a function of temperature, since both \( f_\alpha \) and \( \langle \sigma v \rangle \) have a temperature dependence. Next, since the hot spot is surrounded by cold, dense fuel, heat will be lost through electron conduction. Ion conduction is slower by a factor of \( (m_e/m_i)^{1/2} \sim 50 \) and hence negligible. Thermal energy lost by the hot spot (with surface area \( S \equiv 4\pi R_h^2 \) and volume \( V \equiv 4/3\pi R_h^3 \)) to the surrounding cold fuel per second per unit volume is given by:

\[ W_c = -\chi_e \nabla T_e S/V \approx A_c T_h^{7/2}/R_h^2, \]  

(1.14)

where \( T_h \) is the hot-spot temperature, \( \chi_e \) is the Spitzer conductivity (\( \propto T_h^{5/2} \)), and \( A_c \) is a constant with a weak dependence on density (it incorporates the coulomb logarithm \( \ln \Lambda \)). Radiative losses are due to electrons stopping on high-Z ions and other electrons, and can be expressed as bremsstrahlung power losses per unit volume:

\[ W_r = A_r \rho_h^2 T_h^{1/2}, \]  

(1.15)

where \( A_r \) is a numerical constant. The last loss mechanism concerns mechanical work done by the hot spot on the surrounding cold fuel, and this term looks different depending on different hot-spot thermodynamic conditions. In this work, we consider isobaric and isochoric hot-spot formation (Fig. 1.4), as discussed in Sec. 1.2.3. The isobaric condition requires that the pressure of the hot spot is equal and opposite to the pressure exerted by the compressed cold fuel onto the hot spot. Under these conditions, the hot spot will have a lower density (and larger temperature) than the surrounding fuel. Furthermore, there is no change in the hot-spot volume since it is in thermodynamic equilibrium with the surrounding fuel, and the mechanical work done by the hot spot \( \int pdV \) is zero. For the isochoric case, the central volume of the fuel is heated by external means when the fuel is compressed to its peak density. Heating occurs quickly compared to fuel assembly time-scales (by design), resulting in a hot spot with a density comparable to the surrounding fuel. In this case, the hot spot is not in thermodynamic equilibrium with the surrounding cold fuel, and does work on the fuel that is given by:

\[ W_m = A_m \rho_h T_h^{3/2}/R_h, \]  

(1.16)

where \( A_m \) is a constant. The loss terms (Eqs. 1.14, 1.15 and 1.16) can be combined with the self-heating condition (Eq. 1.12) and the fusion power deposited (Eq. 1.13) to obtain the criterion for self-heating that depends only on the hot-spot temperature:

\[ \rho_h R_h > A_m^{3/2} T_h^{3/2} + \sqrt{A_m^{3/2} T_h^{3/2} - 4A_c A_r T_h^{1/2} + 4\langle \sigma v \rangle A_c f_1 f_2 f_\alpha T_h^{7/2} W_\alpha [m_i (1 + \delta_{12})]^{-1}} \]  

(1.17)

This condition is indicated by the solid lines in Fig. 1.5 for isobaric and isochoric hot-spot configurations (equimolar DT). Shown alongside these curves are regions where self-heating does not occur due to either radiation losses (at large areal densities) or conduction losses (at high temper-
Figure 1.4. Hot-spot density and temperature profiles for (a) isobaric and (b) isochoric configurations. Isobaric configurations are characterized by a central hot spot surrounded by a cold, dense fuel. Isochoric hot-spot configurations are inherently asymmetric since the hot spot is formed by externally heating a compressed shell.

The simple model presented here is useful as it provides physical insight and gives a rough idea of the temperatures and areal densities required for self-sustained fusion. A number of important effects, however, are not taken into account in this model. For some values of $\rho_hR_h$ and $T_h$ outside of the self-heating region, ignition can still occur. In particular, the heat flux leaving the hot spot, a loss mechanism in this model, will ablate and heat an inner portion of the surrounding cold dense fuel, which then becomes part of the hot spot. The lost heat flux effectively recycles the lost energy-density into the hot spot, and can result in a trajectory in the $\rho_hR_h-T_h$ space that leads to ignition. In addition, this model does not incorporate expansion losses of the hot spot associated with decompression of the fuel (even in the isobaric case, the hot spot and fuel disassemble after stagnation). Since the hot spot expands against the surrounding fuel shell, the expansion losses depend on the shell $\rho R$. Furthermore, in the isobaric case, the energy density and hence the pressure of the hot spot comes from the kinetic energy of the converging shell. For this reason, we expect the hot-spot parameters at stagnation to be linked to the fuel shell $\rho R$ as well. These dynamic effects have been studied using simulations for marginally igniting targets. The result is the following ignition condition:

$$\langle \rho R \rangle > \left( \frac{4.4}{\langle T_h \rangle} \right)^{2.2},$$

where the $\rho R$ (in g-cm$^{-2}$) and hot-spot temperature (in keV) are burn-averaged quantities over the implosion. This form is convenient since these burn-averaged quantities are routinely measured. Eq. 1.18 suggests that a burn-averaged temperature of $\sim 4$ keV and $\rho R$ of $\sim 1$ g-cm$^{-2}$ is required.

Figure 1.5. Hot-spot heating and cooling regimes as a function of the hot-spot areal density and temperature. The self-heating regimes for isobaric and isochoric configurations are indicated by the solid lines. For high temperatures and low densities, conduction losses dominate and cool the hot spot. For low temperatures and high densities, bremsstrahlung radiation losses dominate and cool the hot spot.
1.2 Inertial Confinement Fusion (ICF) Using Laser Drivers

for ignition. An important parameter for fusion energy production is the gain. The gain is defined as the ratio of the fusion energy output to the laser energy input:

\[ G = \frac{E_{\text{fusion}}}{E_{\text{laser}}} \approx f(\rho_R). \] (1.19)

The fusion energy output \( E_{\text{fusion}} \) depends on the burn fraction and hence the fuel \( \rho_R \) (it has a weak dependence on the hot-spot temperature, as discussed in Sec. 1.2.1). The gain does not explicitly depend on hot-spot parameters, although for isobaric configurations the fuel and shell \( \rho_R \) are related as previously discussed. Different target gains can be achieved given a pair of \( \rho_R \) and \( T \) that satisfy the self-heating or ignition requirements. For example, two marginally igniting targets can have different shell \( \rho_R \) and hence the burn fractions, resulting in different target gains. Igniting targets will generally have large gains, since the ignition condition (Eq. 1.18) requires moderate \( \rho_R \) and temperature, but a specific target gain is not required for ignition. For these reasons, the target gain is considered an engineering parameter that is important for energy production rather than an ignition physics parameter.

1.2.3 Compression and Hot-spot Generation

Several schemes have been proposed for the compression and inertial confinement of spherical targets. Two of these are the direct-drive and indirect-drive approach. In the case of the former, the laser irradiates the target directly, whereas in the latter an intermediate cylindrical cavity, known as a hohlraum, absorbs and converts laser energy into x-rays that irradiate the target. In principle, direct-drive has the ultimate advantage of increased target gains because the direct coupling of laser energy to the capsule is much higher than the energy deposited by x-rays. The reason for this is the inefficient conversion of laser light to x-rays in the hohlraum. Indirect-drive, on the other hand, provides more uniform target irradiation, which is an important parameter for hot-spot formation and thus ignition, as discussed in the following section. The indirect-drive approach to ICF is currently being pursued at the National Ignition Facility (NIF)\(^{11}\) (see Sec. 1.3.1), where the physics of fuel assembly and hot-spot formation are being studied with the ultimate goal of igniting a DT fuel pellet.

In this thesis, we restrict the discussion to the direct-drive approach to ICF, which is an active area of research at the Omega Laser Facility (see Sec. 1.3.2). The underlying physics of target compression and ignition can be studied using relatively low-energy lasers such as OMEGA, which are incapable of generating the conditions required for ignition. This is because the self-similarity principle of fluid mechanics allows the study of hydro-equivalent target implosions at a much smaller scale. The areal densities and other pertinent results from these experiments can, in principle, be scaled to ignition targets.

For the direct-drive approach, a number of methods have been proposed for generating the central hot spot and initiating a burn wave. Among these are conventional central hot-spot ignition, fast ignition and shock ignition, to name a few. Here, we restrict ourselves to the conventional and

![Figure 1.6. Compression and hot-spot generation in (a) central hot-spot ignition and (b) fast ignition. For central hot-spot generation, the converging shell heats the gas at the center of the target, forming the hot spot. In fast ignition, the hot spot is generated by an external driver.](image-url)
fast ignition approach.

Conventional Hot-spot Ignition

The conventional hot-spot ignition approach (Fig. 1.6a) uses a set of lasers to uniformly irradiate a layered cryogenic target (Fig. 1.7a). The outermost layer of the target is the shell, which is comprised of either plastic or glass, and an inner layer of solid fuel. The outer layer, which sometimes includes high-Z dopants, absorbs the incident laser energy. The inner fuel layer is composed of DT-(or D$_2$-) ice to maximize the initial density of the fuel and hence minimize the adiabat (see Eq. 1.11).

In some cases, the fuel is replaced with material identical to that of the shell’s outer layer. These “surrogate” targets (Fig. 1.7b) are easier to manufacture and are often used to study implosion physics (without fuel burn). The center of the capsule typically consists of vapor from the solid fuel (e.g., DT vapor) or in the case of surrogate targets, manually filled with DT, D$_2$ or D$_3$He gas. The overall size of the target is determined by the amount of available laser energy, as discussed below.

A typical direct-drive laser-pulse is shown in Fig. 1.8a. The pulse consists of several high-intensity, short-duration “pickets” that launch shocks into the shell which set the adiabat. The scaling between laser intensity and the pressure applied to the target (sometimes called the ablation pressure) is given by: 12

$$p_a \propto \left( \frac{I}{\lambda L} \right)^{2/3},$$  \hspace{1cm} (1.20)

where $I$ is the laser intensity and $\lambda L$ the laser wavelength. The picket intensities and timing thus prescribe the pressure and coalescence of the shocks and hence the shell adiabat. A large adiabat is desirable on the outside of the shell for stability considerations while a low adiabat is required in the fuel for effective compression, as discussed previously. To this end, the shocks launched by the pickets decay in strength as they travel through the shell and fuel, and thereby shape the adiabat appropriately. 14 Following the pickets, the laser intensity is slowly raised not to shock the fuel any further. For implosions where a low-adiabat is not desired (e.g., for studies that exclude the physics of high-density fuel assembly), the laser pickets and ramp in laser intensity are not required and a square laser pulse (Fig. 1.8b) is instead often used.

As the laser intensity is increased, the outer layer of the shell is ablated and expands outward. For this reason the shell’s outer layer is often called the “ablator.” The outward expansion of the ablated shell accelerates the remaining shell mass inward due to momentum conservation. The laser then reaches peak intensity and acceleration continues until the laser is turned off. For efficient laser absorption, the maximum intensity must not exceed $\sim 10^{15}$ W-cm$^{-2}$. This is because high intensity laser-light couples energy to relativistic electrons that escape the target, resulting in poor ablation and heating of the shell, as discussed further in Sec. 1.4. The pulse duration is thus largely
determined by this intensity and the total available laser energy. Typical pulse durations are several nanoseconds. Once the laser turns off, the shell coasts at a constant implosion velocity \( V_i \) until it decelerates due to internal pressure buildup from spherical convergence, and eventually stagnates to a final fuel \( \rho R \). As the fuel converges, the vapor trapped within it is compressed and heated to form the central hot spot. Stagnation occurs when the pressure of the hot spot equals that of the fuel, resulting in an isobaric hot-spot configuration. This process is not efficient, because much of the laser energy is coupled to the ablated plasma, rather than to the kinetic energy of the imploding shell. The hydrodynamic efficiency, defined as the ratio of shell kinetic energy to laser energy, is typically of order 10\%.\(^6\)

Both the fuel \( \rho R \) and the implosion velocity are essential parameters because they determine the energy gain of the target:\(^6\)

\[
G \approx \frac{73}{I_{15}^{0.25}} \left( \frac{3 \times 10^7}{V_i (\text{cm/s})} \right)^{1.25} \left( \frac{\Phi(\rho R)}{0.2} \right) \left( \frac{0.35 \lambda_L (\mu\text{m})}{0.35} \right)^{0.5}, \quad (1.21)
\]

where \( I_{15} \) is the laser intensity in units of \( 10^{15} \) W-cm\(^{-2} \), \( V_i \) is the implosion velocity in cm-s\(^{-1} \), \( \Phi(\rho R) \) is the burn fraction, and \( \lambda_L \) the laser wavelength in microns. For a fixed laser energy and fuel adiabat, the \( \rho R \) and hence the fuel fraction is fixed (see Eq. 1.11), the target gain depends only on the implosion velocity. In this case, lower implosion velocities can result from the use of more massive targets that release more energy, resulting in increased gains, as shown by Eq. 1.21. For conventional (isobaric) hot-spot configurations, a minimum implosion velocity of \( \sim 3.5 \times 10^7 \) cm-s\(^{-1} \) is required for ignition.\(^{15}\) This is because the hot-spot \( \rho R \) and temperature at stagnation are largely determined by the implosion velocity and adiabat. The hot-spot averaged ion temperature, in particular scales with the implosion velocity and fuel adiabat as:\(^6\)

\[
\langle T_h \rangle \approx \frac{2.96}{\alpha^{0.15}} \frac{V_i}{3 \times 10^7}
\]

where \( V_i \) is the implosion velocity in cm-s\(^{-1} \). This equation can be combined with the expression that relates laser energy to fuel \( \rho R \) (Eq. 1.11) and the ignition condition that relates \( T_h \) to fuel \( \rho R \) (Eq. 1.18) to obtain the total laser energy required for ignition:\(^6\)

\[
E_{\text{L}}^{\text{ign}} = 0.64 I_{15}^{-0.26} \alpha^{1.9} \left( \frac{3 \times 10^7}{V_i (\text{cm/s})} \right)^{6.6} \left( \frac{\lambda_L (\mu\text{m})}{0.35} \right)
\]

where \( I_{15} \) is the laser intensity in units of \( 10^{15} \) W-cm\(^{-2} \), and the other symbols have their usual meanings. Eq. 1.23 is a function of the adiabat and the implosion velocity (and a weak function of intensity), and has values in the range of several MJ. Targets are thus sized to achieve the required

---

**Figure 1.8.** Typical (a) adiabat-shaping (\( \alpha \sim 1-3 \)) and (b) square laser pulses (\( \alpha \sim 10-20 \)) used at the Omega Laser Facility to irradiate the target shell. The adiabat-shaping pulse uses short-duration, high-intensity “pickets” to launch shocks into the target and set the shell adiabat.
implosion velocity.

In addition to these implosion constraints, the formation of a hot spot requires tight control over laser illumination and target uniformity. High-mode non-uniform irradiation seeds fluid instabilities (e.g., Raleigh-Taylor (RT)) that can cause the cold dense fuel to mix with the hot spot and quench the burn. Laser smoothing techniques, including spatial beam-shaping and removal of laser hot spots, are employed at the Omega Laser Facility (see Sec. 1.3.2) to ensure irradiation uniformity. In addition, low-mode irradiation asymmetries cause $\rho R$ asymmetries and possibly inefficient conversion of shell kinetic energy to hot-spot thermal energy.

Another present challenge for direct-drive hot-spot ignition is efficient compression of the target to the high $\rho R$ required for ignition and burn. As discussed previously, this requires a low-adiabat throughout the compression phase of the implosion. Although a low-adiabat is initially prescribed by the laser drive, it can be altered due to heat sources and sinks as the target is compressed. During the shell acceleration phase, heating of the cryogenic fuel by energetic electrons or x-rays can raise the adiabat and thus degrade compression. Energetic electrons can be produced throughout the shell acceleration phase (during which the target is irradiated) from laser-plasma interactions, as discussed in Sec. 1.4. The impact of these energetic electrons on implosion performance has been studied in this thesis (see Chapter 4).

**Fast Ignition**

Fast ignition is an advanced fusion concept that has been described extensively in the literature.\(^{16,17}\)

In this scheme, separate lasers are used to compress and subsequently heat the core (Fig. 1.6b). The decoupling of compression and hot-spot formation have a number of advantages. First, constraints on laser illumination and target uniformity are significantly relaxed, since the hot spot (isochoric configuration) is formed by an external driver at peak fuel compression. Additionally, the implosion velocity can be reduced significantly, either by lowering the laser energy or by using more massive targets, resulting in increased target gain. Finally, core heating in fast ignition is in principle more efficient as compared to central hot-spot ignition. This is because the central hot-spot ignition approach relies on imploding shells with low hydrodynamic efficiency to form the hot spot. The advantages associated with the fast ignition approach outweigh the higher hot-spot $\rho R$ and temperature required for isochoric hot-sport configurations, but fast ignition has other challenges associated with hot-spot formation as discussed at the end of this section.

The laser configuration used to drive the shell for fast ignition is similar to that of conventional ignition. The laser intensity, pulse duration and shape are similar since in both cases a high-density fuel assembly is required. Heating of the core is achieved using an additional laser that generates either protons, heavy ions, or electrons. These charged-particles are transported to the dense core where they deposit energy, raise the temperature and generate a hot spot. Transport of these charged-particles to the core can be achieved by laser hole-boring\(^{17}\) or by using cones\(^{18}\)

In this work, we restrict ourselves primarily to cone-guided fast-ignition, as shown in Fig. 1.6. These targets consist of cones made of high-Z materials embedded inside spherical shells. Shells are conventionally imploded (with nanosecond laser pulses, or “long-pulse” lasers). At peak compression, a high-intensity laser ($\sim 10^{19}$ W-cm\(^{-2}\)) with a pulse-duration of a few picoseconds is used to generate multi-MeV electrons from the tip of the high-Z cone material (e.g., gold or copper). These electrons are stopped within the core, where they deposit energy, raise the temperature and generate a hot spot. For core-heating to be effective, the electrons must have a range that is comparable to the hot-spot $\rho R$ required for ignition. For $\rho Rh \sim 0.4$ g-cm\(^{-2}\) and $T_h \sim 8$ keV (see Fig. 1.5), this requires hot electrons with energies of 1-3 MeV. In addition, the total energy of the hot electrons (number of hot electrons) is prescribed by the hot-spot mass and the required hot-spot temperature. Since $\rho hR_h$
is fixed, the total hot-spot mass is inversely proportional to the hot-spot density \( M_h \propto 1/\rho_h^2 \). For higher core-densities, and hence smaller cores, a lower hot-spot mass and hence lower total electron energy is required. The choice of density and total energy is then an engineering tradeoff that must consider the available short-pulse laser energy, the energy coupling efficiency of the short-pulse laser to hot electrons and the size of the generated hot electron beam, which cannot be larger than the intended core size. These considerations typically lead to required core density of 300 g-cm\(^{-3}\), hot-spot radius of \( \sim 15 \) \( \mu \)m, short-pulse laser intensity of \( \sim 10^{19} \) – \( 10^{20} \) W-cm\(^{-2}\) and laser energy of \( \gtrsim 40 \) kJ.\(^{19}\)

As in the case of central hot-spot ignition, a minimum implosion velocity is required for the fast-ignition approach, although for different reasons.\(^{6}\) Velocities of \( \sim 10^7 \) cm-s\(^{-1}\) are required to achieve the desired core density (\( \sim 300 \) g-cm\(^{-3}\)). This velocity nevertheless lower than the \( \sim 3.5 \times 10^7 \) cm-s\(^{-1}\) required for central hot-spot ignition, and hence higher target gains are possible for comparable \( \rho R_s \) (and by Eq. 1.11, comparable long-pulse laser energies). As in the case of central hot-spot ignition, the overall target is sized based on the required implosion velocity while the total required long-pulse energy for fast ignition is determined by the desired \( \rho R \) (the desired gain).

The current challenges for electron fast ignition are the generation and transport of the multi-MeV electrons to the dense core. The generated electron beam (see Sec. 1.4.2) must be collimated so that it can efficiently heat a localized region of the dense core to form the hot spot as discussed earlier. Proper collimation can be hindered by relativistic laser-plasma interactions. These undesirable laser-plasma interactions can also alter the mean energy of the generated electrons and lead to poor energy coupling to the dense core. In this context, aspects of hot-electron coupling to the dense core of a compressed target have been studied in this thesis (see Chapter 5).

An alternative approach of using protons and heavier ions to heat the core has also been considered.\(^{20,21}\) Core heating with protons has some advantages over electrons. Ion beams can be more readily focused and made quasi-monoenergetic, resulting in more localized power deposition. The main drawback of using protons is that they are less penetrating, and hence must be more energetic than electrons for the same core density. The generation of monoenergetic multi-MeV protons and heavier ions is an area of active research, and is one of the topics explored in this thesis (see Chapter 6).

### 1.3 Laser Facilities for ICF

The study of different ICF fusion concepts requires different types of lasers that can irradiate targets with a range of intensities, pulse durations and wavelengths. Several lasers were used throughout this work, and their capabilities are summarized here. Sec. 1.3.1 presents a brief overview of the National Ignition Facility (NIF), where ignition-scale experiments are currently being conducted. Sec. 1.3.2 and 1.3.3 discuss features of the Omega and Multi-Terawatt laser facilities that were used throughout the course of this work.

#### 1.3.1 The National Ignition Facility (NIF)

The NIF,\(^{11}\) located in Lawrence Livermore National Laboratory (LLNL), is the largest and most energetic laser built to date. It consists of 192 laser beams and was designed to deliver a total energy of approximately 1.9 MJ in the ultra-violet (\( \lambda_L = 0.33 \) \( \mu \)m). The NIF has been used as part of the National Ignition Campaign (NIC), which has the primary goal of compressing and igniting a DT-fuel pellet using an indirect-drive laser configuration.
Operation of all 192 beams at total energy of nearly 1.9 MJ (UV) is routinely conducted. Currently, NIC involves tuning campaigns used to optimize implosion parameters such as the shell velocity, fuel adiabat, shape (symmetry) and fuel-shell mix.\textsuperscript{22}

1.3.2 The Omega Laser Facility

The Omega Laser Facility, located at the University of Rochester’s Laboratory for Laser Energetics, is comprised of the OMEGA\textsuperscript{23} and OMEGA EP\textsuperscript{24} lasers, as shown schematically in Fig. 1.9 (Note that Omega refers to the facility, comprised of both lasers, whereas OMEGA refers to that laser system alone). OMEGA and OMEGA EP each have dedicated target and laser bays, consisting of a target chamber and several beamlines.

In this thesis, the OMEGA laser was used to study the impact of hot-electron generation in ICF implosions, as discussed previously. The OMEGA EP laser was used to study the coupling of hot electrons to the dense fuel cores as relevant to fast ignition. The capabilities of these lasers are discussed below with particular emphasis on laser parameters such as contrast, smoothing and polarization that can have an impact on laser-plasma interactions, hot-electron generation and hence fast-ion acceleration.

OMEGA is a 60-terawatt laser, consisting of 60-beams capable of delivering 30 kJ of energy in the ultraviolet (UV) in a few nanoseconds. Three independent laser drivers allow for a variety of high-fidelity pulse shapes. These pulses are amplified with a neodymium-glass (Nd:glass) amplifier ($\lambda = 1053$ nm) then frequency-tripled ($\lambda = 351$ nm) with an efficiency of 75\% and finally focused onto a target inside the OMEGA target chamber. Distributed Phase-Plates (DPP),\textsuperscript{25} used for super-Gaussian spatial-shaping of the laser beam (so that the laser intensity is more uniform in

\begin{figure}[h]
\centering
\includegraphics[width=\textwidth]{Omega_Laser_Facility.png}
\caption{The Omega Laser Facility’s OMEGA and OMEGA EP lasers. Each laser has a dedicated target chamber and several beamlines. The OMEGA EP laser is capable of stand-alone operation or may be transported and focused into the OMEGA target chamber for simultaneous use with the OMEGA laser [Source: C. Stoeckl et al., Diagnostics for Fast-Ignitor Experiments on OMEGA/OMEGA EP].}
\end{figure}
space), set the on-target laser spot-size at best focus (100 μm diameter at best focus). One consequence of this technique is the laser-interference pattern produced due to the corrugated surface of the DPP. These laser speckles are smoothed out using Smoothing by Spectral Dispersion (SSD), whereby different frequency components of a laser pulse are dispersed spatially along a phase-plate so as to mitigate spatial interference. The presence of laser speckles can seed undesirable laser-plasma instabilities that are detrimental to the performance of ICF implosions. In this context, laser coherence and polarization are also of concern since absorption mechanisms depend on these laser parameters. OMEGA uses Distributed Polarization Rotators (DPR) to achieve a random mix of s- and p- polarization, since a high-intensity of either polarization can result in undesirable absorption mechanisms. The DDP also play an important role in mitigating unwanted laser-plasma coupling by breaking the coherence of the laser beam. Laser-plasma interactions in the context of ICF are discussed in Sec. 1.4.

The OMEGA EP (extended performance) laser was qualified for operation in 2009 to complement the OMEGA laser by adding petawatt capability in addition to stand-alone operation. The OMEGA EP laser consists of a four beamlines: two short-pulse (∼ ps) and two long-pulse (∼ ns) laser beams. The short-pulse beamlines are amplified using optical parametric chirped-pulse amplifiers (OPCPA), whereby pulses are stretched in time, amplified (with a Nd:glass amplifier), and compressed in time again. One of the two short-pulse beams may be transported into the OMEGA target chamber and subsequently focused on-target using an off-axis parabola. This complements the OMEGA long-pulse capability and opens the door for unique experiments, such as fast ignition. Alternatively, both long-pulse and short-pulse OMEGA EP beams may be used simultaneously within the OMEGA EP target chamber. The short-pulse and long-pulse beams are capable of delivering maximum on-target energies of 2.6 kJ and 6.5 kJ and minimum pulse durations of 1 ps and 100 ps, respectively. For 1-10 ps pulses, the maximum energy of the short-pulse beams is a function of the pulse duration because of damage thresholds of the optical coatings. At present, approximately ∼ 2 kJ of energy can be delivered to a target in 10 ps within a focal spot size of 20 μm (defined as the radius which encircles 80% of the incident energy), resulting in an on-target intensity of ∼ 10^{19} W-cm^{-2}. Higher intensities are possible for shorter pulse durations (e.g., 1 kJ in 1 ps, or a petawatt). Near-future upgrades of the laser are expected to increase the laser intensity by reducing the focal spot size in half. The laser contrast, defined as the ratio of peak laser power to the power in the laser prepulse, is an important parameter for short-pulse systems. The contrast determines the amount of pre-ionization of the target material before the main drive reaches the target, and this can be important for several applications. The OMEGA EP contrast is 10^6 at present, and expected to increase to 10^8 in the near-future.

1.3.3 The Multi-Terawatt (MTW) Laser Facility

The Multi-Terawatt (MTW) laser (Fig. 1.10) at LLE is a single-beam, ultrahigh-intensity short-pulse laser. The MTW laser served as a prototype front-end for the OMEGA EP laser, and continues to function as a stand-alone laser, capable of delivering on-target energies as high as 10 J of IR light (1053 nm) in 1 ps. On-target intensities as high as 5 × 10^{19} W-cm^{-2} with a contrast of 10^8 are possible at best focus, where 50% of the laser energy is focused to a 5 μm diameter spot.

Laser pulses (temporal Gaussian, spatial super-Gaussian) are generated from a single oscillator and amplified using the chirped-pulse amplification (CPA) technique. The pulses are then focused onto a target inside the MTW target chamber using an off-axis parabola (f/3). The time-duration of the laser pulse is measured on-line for each system shot using a time-expanded-single-shot-autocorrelator (TESSA), whereby the main beam is sampled, converted to second harmonic light and measured using a streak-camera. The laser contrast and beam spot-size are not measured
on-line, though they have been well-characterized.

In this work, the MTW laser was used to irradiate thin-foil targets at high-intensities to study the generation of fast ions (Chapter 6) for potential use in ion fast-ignition schemes. Generation of fast ions are also of interest for imaging and medical applications, including charged-particle radiography and proton cancer therapy.

1.3.4 Summary of Laser Specifications

These laser facilities provide flexibility for tests and validation of a variety of inertial fusion concepts. The OMEGA laser is routinely used for both direct- and indirect- drive spherical implosions. The petawatt capability added to OMEGA by the EP laser has been used in the last several years to test advanced fusion concepts such as fast ignition. The MTW facility complements the Omega lasers by facilitating studies of ultra-intense laser-matter interactions in connection with advanced fusion concepts and other HED applications such as medicine. A summary of the OMEGA, OMEGA EP and MTW laser parameters are presented in table 1.1. Not all parameters are relevant to all laser types (e.g., contrast in long-pulse lasers).

<table>
<thead>
<tr>
<th></th>
<th>OMEGA</th>
<th>OMEGA EP</th>
<th>MTW</th>
</tr>
</thead>
<tbody>
<tr>
<td>Wavelength</td>
<td>351 nm</td>
<td>1053 nm</td>
<td>351 nm</td>
</tr>
<tr>
<td>Total Energy</td>
<td>30 kJ</td>
<td>2.6 kJ</td>
<td>6.5 kJ</td>
</tr>
<tr>
<td>Pulse Duration</td>
<td>1 ns</td>
<td>1-100 ps</td>
<td>0.1-100 ns</td>
</tr>
<tr>
<td>Intensity Contrast</td>
<td>-</td>
<td>$10^4$</td>
<td>-</td>
</tr>
<tr>
<td>Laser Diameter</td>
<td>$&gt; 100$ μm</td>
<td>$&gt; 20$ μm</td>
<td>$&gt; 100$ μm</td>
</tr>
<tr>
<td>Max. Intensity</td>
<td>$&lt; 10^{17}$ W-cm$^{-2}$</td>
<td>$1 \times 10^{19}$ W-cm$^{-2}$</td>
<td>$1 \times 10^{19}$ W-cm$^{-2}$</td>
</tr>
</tbody>
</table>
1.4 Laser-plasma Interactions (LPI) and Fast-Ion Generation

Laser-plasma interactions (LPI) can generate significant populations of suprathermal electrons in ICF implosions and laser-foil experiments. Depending on the confinement and heating scheme, LPI-generated hot electrons are either a byproduct (as in the case of central hot-spot ignition) or integral to the ICF concept (in the context of electron fast ignition). The study and characterization of these electrons is thus important in diagnosing implosion performance.

Direct spectral measurement of the hot electrons is generally difficult. From the standpoint of diagnostics, absolute measurements of emitted electron spectra can be made, but direct measurements of the initial hot-electron spectrum are non-trivial. This is because as hot electrons begin to expand, they initially leave behind a net positive charge that then acts to modify the energy spectrum of the faster escaping electrons. The physics of this process must be unfolded and depends on the details of electron interaction with the target. Since the generation of hot electrons inevitably accelerates ions, spectral measurements of these ions can be used to deduce quantitative information about the hot electrons and to diagnose the performance of ICF implosions. The study of fast ions as they relate to hot-electron production is the main subject of this thesis, and is discussed further in Sec. 1.5. Over the last decade, much work has been done to link characteristics of the emitted proton spectrum to the initial hot-electron distribution.

The nature of the LPI and the temperature of electrons generated depend on the physical mechanism for laser absorption. This in turn depends on laser parameters such as the wavelength, intensity and pulse duration, and on target parameters such as the materials, thickness and densities. Sec. 1.4.1 discusses laser absorption mechanisms relevant to ICF implosions, including LPI that results in the production of (undesirable) suprathermal electrons. Sec. 1.4.2 presents LPI in the short-pulse regime that are integral to advanced fusion concepts such as electron and proton fast ignition. Finally, Sec. 1.4.3 presents an overview of ion acceleration due to charge-separation of hot electrons from the bulk plasma. In addition to a brief review of laser-plasma interactions, particular emphasis is given in the following sections to laser-absorption mechanisms that generate the fast ions measured throughout this work.

1.4.1 Laser Absorption and Hot-Electron Production in ICF Implosions

As laser light impinges the target, it irradiates the outermost layer of the shell and breaks down the shell material into a plasma with a density scale length of order several hundred microns. Light (electromagnetic) waves that have a normal angle of incidence with respect to the density gradient can propagate through this coronal plasma up to the critical density, at which point the plasma electrons screen the laser light and prevent it from propagating further. The critical number density is found by setting the laser frequency equal to the local frequency for electrostatic plasma oscillations ($\omega_{pe}$):}

$$n_c = \frac{4\pi^2 c^2 m_e e^2}{\epsilon_0 \omega_{pe}^2 \lambda_L^2}$$

where $m_e$ is the electron mass, $\epsilon_0$ the permittivity of free space, $e$ the electron charge, and $\lambda_L$ the laser wavelength ($\lambda_L = 0.35 \mu m$ for direct-drive implosions on Omega). At the critical density, the laser light is partially reflected. A fraction of the incident light can tunnel further into the overdense plasma. This wave decays rapidly in the overdense plasma and does not carry any energy and is hence said to be evanescent.

Obliquely incident light waves do not reach the critical surface but reflect at a turning point within the underdense plasma. Obliquely incident light waves reflect when the local plasma fre-
frequency is a fraction of the laser frequency:

\[ \omega_{pe} = \omega \cos \theta, \]  

where \( \omega_{pe} \) is the local plasma frequency at the turning point, \( \omega \) is the laser frequency, and \( \theta \) is the angle of incidence between the incident light wave vector (\( \mathbf{k} \)) and the density gradient (\( \nabla n \)). The turning point thus depends on the density profile of the underdense plasma (e.g., for a linear density profile, the turning point is given by \( L \cos^2 \theta \), where \( L \) is the density scale length). The local density scale length of the underdense plasma plays a different but important role in each of the absorption mechanisms discussed below.

**Collisional Absorption**

Laser light propagating in the underdense (\( \omega > \omega_{pe} \)) coronal plasma can be absorbed by electrons through inverse bremsstrahlung. The energy of these electrons is then transferred to the ions through electron-ion collisions. The fraction of laser light absorbed for collisional absorption is given by:

\[ f_A = 1 - \exp \left( -\frac{8 \nu_{ei}^* L}{3 c} \cos^3 \theta \right), \]

where \( \nu_{ei}^* \) is the electron-ion collision frequency evaluated at the critical density, \( L \) is the plasma density scale length, \( c \) is the speed of light, and \( \theta \) is the angle of incidence between the light wave and the density gradient of the coronal plasma. The absorbed fraction depends on the coronal temperature (through \( \nu_{ei}^* \)) since a hotter plasma is less collisional and results in less energy transfer from electrons to ions. The absorption also depends on the details of the density profile, since a higher plasma density is more collisional. Eq. 1.26 was computed for an exponential profile. The angle of incidence enters into the absorbed fraction because for a given density profile and temperature, a more oblique angle of incidence will result in the wave turning around before it can reach sufficiently high densities where the absorption is greater. A final subtlety in Eq. 1.26 is the implicit dependence on laser wavelength. Eq. 1.26 is expressed in such a way that the collision frequency must be evaluated at the critical density, which scales with the laser wavelength as \( 1/\lambda^2 \). Shorter wavelengths are thus more penetrating and allow access to higher densities and hence greater absorption fractions. This is partly the reason why high-intensity Nd:glass lasers such as OMEGA are frequency tripled to achieve wavelengths of 0.33 \( \mu \)m. In general, the absorption is time-dependent and must be calculated self-consistently, since the initial coronal temperature is cold and becomes less collisional as the corona absorbs energy and heats up. Typical corona temperatures are \( \sim 1-2 \) keV, and the absorption-fraction is \( \sim 70\% \) for plastic shells irradiated at an intensity of \( 10^{14} \) W-cm\(^{-2} \). This fraction includes s- and p- polarized light and decreases with increasing laser intensity (\( \sim 30\% \) at \( 10^{15} \) W-cm\(^{-2} \)) because of higher coronal temperatures that result from the use more intense laser light.

**Resonance Absorption**

Resonance absorption (RA) occurs when laser light is converted to electron plasma waves in a resonant layer near the critical surface. This occurs only for p-polarized light, where the component of the electric field normal to the density gradient drives electrons back and forth. As the p-polarized light wave propagates towards the critical surface, the electric field amplitude grows and reaches a maximum near the critical surface. Within this layer, electrons resonantly respond to the laser electric field. For warm coronal electrons (finite \( T_e \)) these electrostatic oscillations can propagate
due to the electron pressure gradient, and the dispersion relation of these electron plasma waves is given by:  
\[ \omega^2 = \omega_{pe}^2 + 3k^2v_{te}^2 = k^2v_{te}^2 \left( \frac{1}{k^2\lambda_D^2} + 3 \right), \]  
where \( k \) is the wave number, \( v_{te} \) is the thermal velocity of electrons and the other symbols have their usual meanings. The last equality makes use of the relationship between the electron Debye length, plasma frequency and thermal velocity (\( \lambda_D = v_{te}/\omega_{pe} \)). These waves propagate down the density (and temperature) gradient of the coronal plasma. At some region determined by Eq. 1.27 and the detailed density and temperature profiles, the local Debye length becomes comparable to the wavelength (\( k\lambda_D \sim 1 \)). Within this region, the wave phase velocity is sufficiently low enough (but faster than the thermal velocity of the electrons, \( v_{ph} \equiv \omega/k \sim 2v_{te} \)) and the wave can be damped without collisions. In this scenario a small fraction of electrons in the high-energy tail of the thermal distribution, which have velocities comparable to the phase velocity of the wave, see a constant wave electric field and gain energy and thereby damp the wave.

For typical OMEGA implosions, detailed hydrodynamic and wave simulations incorporating density and temperature profiles have shown that the electrostatic pressure of the wave is small compared to the kinetic pressure of the electrons (\( \sim 10-20\% \)). Thus, the electric field amplitudes associated with these waves are small compared to threshold values required for nonlinear damping (e.g., wavebreaking). The linear theory of collisionless damping therefore captures the essential physics and the resonant electron energies have been estimated from the wave’s phase velocity to be \( \sim 5 \text{ keV} \). This number is somewhat higher than typical coronal temperatures (\( \sim 2 \text{ keV} \)), and enhances collisional absorption. Langmuir wave excitation by resonance absorption adds an additional 2-15\% to the total absorption fraction for intensities of \( 10^{14}-10^{15} \text{ W-cm}^{-2} \).

In general, nonlinear processes can be important and need to be taken into account to accurately compute energies of hot electrons generated by damping of electron plasma waves. These processes can include wavebreaking, wave-particle trapping, and higher-order mode-coupling. The latter effect leads to a spectrum of wavenumbers and hence a spectrum of hot electrons. These effects result in an electron distribution with a hot component characterized by a single temperature that also scales with the laser intensity (\( T_h \propto I^{1/3} \)). In typical OMEGA implosions, these nonlinear effects are negligible for electron plasma waves excited by RA, but they can be important for electron plasma waves excited by parametric instabilities, as discussed next.

**Parametric Decay Instabilities and Hot-Electron Production**

Parametric laser-plasma instabilities are a class of LPI that can couple an incident wave to electromagnetic waves. Several instabilities are important for ICF, including Stimulated Raman Scattering (SRS), Two-Plasmon-Decay (TPD) and Stimulated Brillouin Scattering (SBS). SRS is the decay of an incident electromagnetic wave into a scattered wave and an electron plasma wave. The wave matching conditions allow SRS to occur in the underdense coronal plasma up to the quarter-critical density (\( n_e/4 \)). TPD is the decay of an incident wave into two electron plasma waves, and occurs in a narrow region near the quarter-critical density. Finally, SBS involves the decay of an incident light wave into an ion acoustic wave and a scattered light wave and occurs throughout the underdense plasma. These parametric instabilities arise from the same underlying physical picture. An incident electromagnetic wave that propagates through a density perturbation will induce a transverse current (\( \mathbf{J} \)) due to the wave electric field. Since ions are massive and hence fixed on timescales of light-wave propagation, the current is due to electron quiver motion in the laser’s electric field (\( \mathbf{J} = ne\mathbf{v}_{os} \), where \( \mathbf{v}_{os} \equiv e\mathbf{E}/m_e\omega \)). This current directly sources a decaying
wave (either electrostatic or a scattered light wave). The electric fields of the incident ($E_0$) and decay wave ($E_1$) can enhance the density perturbation through the classical ponderomotive force that arises from the quiver motion of electrons in the laser electric fields ($-e^2/4m_e\omega^2 \nabla (E_0 \cdot E_1)$), leading to an instability. In this case, the incident wave loses energy to the decay waves that grow in time.

The three-wave coupling for SRS, TPD and SBS described above must satisfy energy and momentum conservation:

$$\omega_0 = \omega_1 + \omega_2 \quad (1.28a)$$
$$k_0 = k_1 + k_2, \quad (1.28b)$$

where the indices 1, 2 refer to the two decay waves and $i$ the incident (pump) wave. These matching conditions can be satisfied in a number of ways, which lead to different types of decay waves. Here we consider only the Two-Plasmon-Decay (TPD) instability\textsuperscript{35,30} due to its low threshold and predominance in direct-drive implosions on OMEGA.\textsuperscript{36} By definition, this instability occurs when $\omega_1 = \omega_2 \approx \omega_0/2$. Since the decay waves are electron plasma waves, it follows that $\omega_{1,2} \approx \omega_{pe}$ (the local plasma frequency) which can only be satisfied at the quarter-critical surface where $\omega_0 = 2\omega_{pe}$. This instability has both absolute and convective modes,\textsuperscript{37} and is capable of producing electron plasma waves with phase velocities of the order of the speed of light. Due to plasma inhomogeneity, the matching conditions (Eqs. 1.28a and 1.28b) can only be satisfied within a narrow region near the quarter-critical surface.\textsuperscript{37} This results in a threshold for the TPD instability that can be expressed by the parameter:\textsuperscript{35}

$$\eta = \frac{I_{14} L(\mu m)}{230 T_c(keV)} \cdot \frac{\lambda_L(\mu m)}{0.351}, \quad (1.29)$$

where $I_{14}$ is the laser intensity (in units of $10^{14}$ W-cm$^{-2}$), $L$ is the density scale length at $n_{cr}/4$, $T_c$ is the coronal electron temperature and $\lambda_L$ is the laser wavelength. The instability occurs for $\eta > 1$ (for typical OMEGA implosions, $L \sim 150 \mu m$, $T_c \sim 2$ and the instability occurs when $I_{14} \gtrsim 4$-6 W-cm$^{-2}$). The dependence of $\eta$ on these parameters can be understood qualitatively in terms of three-wave phase matching and the feedback mechanism discussed above. Higher intensities result in larger wave electric fields, greater coherence and more feedback through the ponderomotive force, enhancing the instability. Similarly, a flatter density profile at $n_{cr}/4$ (larger $L$) results in a more homogeneous plasma and thus better phase matching. Higher electron temperatures introduce more random motion and break coherence, disrupting the feedback process and mitigating the instability. The wavelength dependence is more subtle. Although shorter wavelengths can more readily phase match in the presence of an inhomogeneous plasma, larger wavelengths result in a greater ponderomotive force which ultimately dominates and reinforces the instability. Further details of TPD-generated hot electrons are discussed in Chapter 4.

### 1.4.2 Ultra-Intense, Short-Pulse Laser-Plasma Interactions

Absorption mechanisms of short-pulse (fs - ps pulse duration) ultrahigh-intensity ($\gtrsim 10^{17}$ W-cm$^{-2}$) lasers are inherently different from the various mechanisms described above for two fundamental reasons. First, “short pulse” refers to laser pulse durations that are short compared hydrodynamic timescales, and for femtosecond pulses, shorter than thermal equilibrium timescales as well. One consequence of this is that the ionized plasma formed at the vacuum-solid interface where the laser interacts does not have time to expand and form sufficiently long scale lengths for many of the
mechanisms described above. In practice, short-pulse lasers have a finite contrast and the consequential low-intensity prepulse that lasts several nanoseconds can ablate the target and produce an underdense plasma before the main short duration pulse arrives. Depending on the severity of this preformed plasma the above mechanisms can contribute to the absorption. Relativistic effects such as filamentation and self-focusing of the laser can also become important, which is further discussed in the context of fast ignition in Sec. 5.5. Second, the higher intensities achieved by these lasers, partly due to shorter pulse durations, means that the quiver motion of electrons in the laser electric fields have a significant impact on light absorption. In particular, ultra-intense is defined as the regime where the normalized momentum of an electron in the laser electric field is unity \( a_0 \equiv p_{osc}/m_e c = 1 \). In addition, the radiation pressure in this regime is large enough to allow laser light to propagate up to several times the classical critical density.

The short-pulse laser regime is important for fast ignition, where an ultrahigh-intensity laser is used to generate (relativistic) electrons to heat the dense core of a compressed shell (see Chapter 5). It is also relevant for proton and heavy-ion fast-ignition, where a similar laser is used to generate hot electrons which then accelerate ions from a conversion foil onto a target (see Chapter 6). Depending on the laser intensity, one of two absorption mechanisms is dominant (although both mechanisms occur simultaneously), as described below.

**Vacuum Heating**

The RA mechanism described in the context of ICF implosions is no longer applicable when the electron quiver motion is relativistic and the scale length of the underdense plasma is small in comparison to the wavelength (e.g., the plasma is an overdense slab). In this regime, often termed vacuum heating, the laser cannot excite electron plasma waves near the critical surface. Instead, electrons see a fraction of the full period of the incident \( p \)-polarized wave, and they are driven by the normal component \( (\mathbf{E} \cdot \nabla n) \) of these large electric fields into vacuum and then into the overdense plasma with a random phase. It has been shown that these electrons are well-described by a single temperature. Background electrons that do not participate in this process can have a different temperature, and a typical system in this regime is characterized by two-temperatures (hot and cold). These hot electrons have been experimentally characterized by Beg et al., who found that the temperature scales with laser intensity and wavelength as:

\[
T_h = 215 \left( I_{18} \lambda_{\mu m}^2 \right)^{1/3} \text{keV},
\]

where \( \lambda_L \) is the laser wavelength in microns and \( I_{18} \) the intensity in units of \( 10^{18} \text{ W-cm}^{-2} \). A relativistic model by Haines et al. has been used to explain the scaling of Eq. 1.30 with laser intensity and wavelength. This mechanism is dominant in short-scale-length plasmas for \( I_\lambda^2 \) in the range of \( 10^{16}-10^{19} \text{ W-cm}^{-2} \mu m^2 \). For typical OMEGA EP short-pulse operating parameters, this mechanism can result in the production of electrons with temperatures as high as 500 keV.

**J × B Heating**

At laser intensities of \( 10^{18}-10^{19} \text{ W-cm}^{-2} \), the relativistic quiver velocity of electrons in the laser fields become large enough to make \( J \times \mathbf{B}_L \) motion significant (where \( J \equiv n e \upsilon_{osc} \) is the current associated with quiver motion of electrons and \( \mathbf{B}_L \) is the laser magnetic field). Electrons thus oscillate due to the laser electric field (in the direction transverse to laser propagation) and due to \( \upsilon_{osc} \times \mathbf{B}_L \) (in the longitudinal direction) at twice the laser frequency \( (2\omega) \). Electrons that have the proper phase will be ejected into the overdense plasma at the critical surface where the laser fields cutoff. This stochastic nature of dephasing from the wave results in a largely Maxwellian distribution of hot
electrons, as confirmed by particle-in-cell (PIC) simulations\textsuperscript{42} and experiments.\textsuperscript{43} The temperature of these electrons is given by equating the thermal energy of electrons with the energy of relativistic quiver motion in the fields (the ponderomotive potential, $U_P$):\textsuperscript{38}

$$T_h = U_P = m_e c^2 (\gamma - 1) \approx \left[ \left( 1 + \frac{I_{18} \lambda_{\mu}^2}{2.38} \right)^{1/2} - 1 \right] \times 511 \text{ keV},$$

(1.31)

where $\gamma$ is the Lorentz factor (associated with quiver motion in the laser fields) and the other symbols have their usual meanings. This mechanism occurs for either s- or p- polarized light. It is more important than vacuum heating when the $J \times B$ driving term is larger than the p-polarized electric field of vacuum heating.\textsuperscript{38} In practice, this occurs at intensities of a few $10^{18}$ W-cm$^{-2}$.

### 1.4.3 Fast-Ion Acceleration

Hot electrons produced from LPI in ICF implosions or planar targets can generate fast ions with energies from hundreds of keV to tens of MeV. Although the source and temperature of the hot electrons depends on the mechanism as discussed above, there are several common elements among ions accelerated by these electrons, as discussed here.

The ions considered in this work are primarily accelerated from a target because of charge separation. In general, the most energetic hot electrons produced escape the target first and leave behind a positive charge that traps the remainder of the hot electrons within the planar or spherical target. This charge is eventually depleted either by acceleration of ions surrounding the target (in the corona) or by return currents, and the remaining electrons can then escape. In this work, electron energies ($50$ keV - $1$ MeV) and typical target dimensions ($\sim 1$ mm) lead to charging of targets over fast time-scales ($\lesssim 1$ ps-100 ps).

In the case of ICF implosions, ions are accelerated radially outward from the coronal plasma surrounding the target (Chapter 4). For short-pulse laser interactions with planar targets (Chapters 5 and 6), a thin layer of ions are accelerated from the rear of the target due to forward-directed energetic electrons, before the majority of the hot-electron space-charge is depleted. The latter is sometimes referred to as Target Normal Sheath Acceleration (TNSA) to distinguish acceleration due to charge-separation from other ion-acceleration mechanisms (not considered in this work) in the context of short-pulse LPI. The resulting energy distribution of these ions has been predicted using fluid models, kinetic theory and simulations. An important feature of fast-ion physics is that the hot electrons and associated fast ions expand faster than the bulk plasma, and that the energy of the most energetic ion (the maximum ion energy) is directly proportional to the hot-electron temperature as shown next.

Consider an electron-ion plasma in a 1D planar slab geometry that is free to expand into vacuum. Initially, the plasma occupies the half-space $x < 0$ with a uniform density of $n_0$ (there is a sharp boundary between the plasma and vacuum). This plasma has an initial background temperature ($T = T_i = T_e$) and a hot-electron component characterized by a single temperature ($T_h \gg T$). From fluid theory, it is well known that the bulk plasma will expand into vacuum at several times the ion sound speed (for an isothermal plasma, $c_s = [(T_i + ZT_e)/m_i]^{1/2}$). The hot component however will accelerate a fraction of the ions from the bulk plasma to higher velocities. The fast ions and the hot electrons that drive the ions expand in vacuum much faster and are spatially separated from the bulk plasma. The problem can then be decomposed into two parts where quasi-neutrality is maintained between the hot electrons and fast ions and separately by the remaining bulk plasma. The fluid equations for the hot electrons and the fast ions then become analogous to that of the bulk plasma, except that the initial temperature of the ions is negligible.
compared to the hot electrons \((T_i \ll T_h)\). Neglecting electron inertia \((m_e/m_i \approx 0)\), the electron momentum equation relates the electric field (created by escaping hot electrons) to the electron density gradient:

\[
E = -\frac{T_h}{n} \frac{\partial n}{\partial x},
\]

where \(n (= n_e = Z n_i)\) is the density of hot electrons. The equations of motion for ions are:

\[
\begin{align*}
\frac{\partial n_i}{\partial t} + v_i \frac{\partial n_i}{\partial x} &= -n_i \frac{\partial v_i}{\partial x}, \\
\frac{\partial v_i}{\partial t} + v_i \frac{\partial v_i}{\partial x} &= -\frac{Ze}{m_i} E.
\end{align*}
\]

These coupled PDEs have analytic solutions for the density \((n)\), velocity \((v_i)\) and electric field \((E)\) given by:

\[
\begin{align*}
n &= n_0 \exp(-x/c_h t - 1) \\
v_i &= c_h + x/t \\
E &= T_h/c_h t.
\end{align*}
\]

It is evident from these equations that the plasma scale length is a function of time \((L \equiv c_h t)\) and the fast ions thus expand at the “hot” sound speed \((c_h \equiv [Z T_h/m_i]^{1/2})\), which is faster than the bulk expansion speed by \((T_h/T)^{1/2}\). For hot electrons in OMEGA implosions, this factor is 5-30 and for short-pulse LPI (vacuum heating or \(J \times B\) heating), it is \(\gtrsim 10\).

In general, analytic solutions such as the one presented above exist for cylindrical and spherical

\begin{figure}[h]
\centering
\includegraphics[width=\textwidth]{figure11.png}
\caption{(a) Initial and (b) late-time structure of the ion front, including the electron and ion densities and potential. In general, the target potential \((V_T)\) is the sum of the voltage drop across the quasineutral region \((\phi_q)\) and across the sheath \((\phi_s)\). Initially, the ion density gradient is steep and the potential is due entirely to charge-separation effects. Later in the expansion, the quasineutral region has a density gradient and a plasma scale length that is given by the product of the hot sound speed and time \((c_h t)\), whereas the ion front has a scale length on the order of the hot-electron Debye length.}
\end{figure}
geometries in the quasineutral limit, and yield the same qualitative results as the planar solution. The limit is useful for the study of collective electric fields and bulk fluid expansion velocity, but it ignores charge separation for spatial scales on the order of the Debye screening length (in this case, the local hot-electron Debye length is given by $\lambda_{D,h} \equiv [\epsilon_0 T_h / ne^2]^{1/2}$). Exact solutions for both the quasineutral region and the expansion front where charge-separation effects are important can be found numerically. Simulated density profiles of the hot electrons and fast ions are shown in Fig. 1.11, along with the relevant scale lengths of the quasineutral region and charge-separation region. The exact values depend on the hot-electron temperature and number density, but are on the order of $L \sim 1$-10 mm and $\lambda_{D,h} \sim 0.1$-1 $\mu$m for the experiments considered in this work. Fig. 1.11 also shows the total potential drop across the expanding plasma ($V_T$) that arises due to both collective effects ($\phi_f$, from the quasineutral region) and charge-separation effects ($\phi_s$, from the expansion front).

The analytic solutions for the density, velocity and electric field (Eqs. 1.34a-1.34c) are invalid when the local density scale length ($n/\nabla n$) becomes smaller than the local Debye length (for the simple planar example presented here, the local density scale length is the same as the plasma scale length, but this is not generally the case). Quasineutrality is violated everywhere for early times in this example (e.g., initially $\lambda_{D,h} \gg L = 0$) and just before the ion expansion front ($r = r_q$ in Fig. 1.11) for all times. The ion front ($r = r_f$) is the edge of the expansion and is characterized by strong charge separation between the fast ions and hot electrons. Within this region, the Poisson’s equation applies and must be solved together with the fluid equations to obtain solutions for all space. Numerical simulations of this problem confirm that the electric fields are largest within the ion expansion front and we thus expect that the most energetic ions can be found in the neighborhood of $r = r_f$. The relative importance of collective effects and charge separation effects on ion acceleration depends on the initial density profile of the ions. If the initial ion density gradient is steep (as shown in Fig. 1.11a), as typically found in short-pulse laser-plasma interactions (Chapters 5 and 6), then charge-separation effects are significant or even dominant. If the density profile varies slowly in space then collective effects are dominant (typical of the coronal plasma in ICF implosions, see Chapter 4). In general, exact analytic solutions that account for both effects are not possible and a numerical treatment must be applied with the appropriate (or approximate) initial ion density profile. Nevertheless, the simple example presented here provides insights and illustrates the underlying physics of ion acceleration.

The condition for violation of quasineutrality ($c_h t / \lambda_{D,h} \approx 1$) can be used to obtain $r_q$ for all times, which can be used to estimate the maximum ion velocity as a function of time as follows. Combining this condition with Eq. 1.34a yields:

$$\frac{c_h t}{\lambda_{D,h}} = \frac{c_h t}{(\epsilon_0 T_h / n_0 e^2)^{1/2}} \exp[(-x/c_h t - 1)/2] = \omega_{pi} t \exp[(-x/c_h t - 1)/2] \approx 1,$$  \hspace{1cm} (1.35)

where $c_h / \lambda_{D,0,h} \equiv \omega_{pi}$, and $\omega_{pi} \equiv (n_0 e^2 / m_i \epsilon_0)^{1/2}$ which is the (fast) ion plasma frequency. It follows from Eq. 1.35 that:

$$r_q = c_h t (2 \ln(\omega_{pi} t) - 1).$$  \hspace{1cm} (1.36)

If $\lambda_{D,h} \ll L$, then $r_q \approx r_f$ and the expression above can be used with Eq. 1.34b to obtain the velocity and hence energy of ions at the expansion front:

$$E_{i,\text{max}} = \frac{1}{2} m_i v_{i,\text{max}}^2 \approx 2 Z T_h \ln(\omega_{pi} t)^2,$$  \hspace{1cm} (1.37)

which is linearly proportional to the hot-electron temperature and the charge ($Z$) of the ion and...
very weakly proportional to $\omega_{pi}$ and acceleration time. In an actual experiment, the maximum ion energy grows logarithmically with time until hot-electron production ceases. The acceleration time is thus taken as the duration over which hot electrons are produced; in some cases this corresponds to the laser pulse duration.

The isothermal assumption made here is not always a good approximation and its applicability depends on the details of the experiment and the LPI mechanism responsible for hot-electron production. Models have been developed by numerically solving the fluid equations for different cases, including adiabatic expansion\(^{46}\) and so-called “two-phase” expansion.\(^{47,48}\) In the former, an initial hot-electron temperature cools over time as fast ions are accelerated while for the latter the temperature is assumed to remain constant while the laser is on and then decays after the laser turns off. The results of these models have the effect of changing the logarithmic term in Eq. 1.37 and hence preserve the linear scaling of the maximum ion energy with hot-electron temperature.

In addition, these fluid models are not applicable when the hot electron distribution is non-Maxwellian. This can occur when the trapped electrons lose energy to surrounding cold matter or if the expansion is collisionless so that hot electrons cannot thermalize. Deviations from a Maxwellian are addressed on a case-by-case basis in the ensuing chapters. In the absence of collisions, a kinetic theory of ion expansion by hot electrons has demonstrated that an initial Boltzmann distribution of hot electrons (Maxwellian hot electrons that are in Boltzmann equilibrium with the potential) is conserved throughout the expansion even as electrons cool and lose thermal energy to ions.\(^{49}\) Since the hot-electron generation mechanisms discussed in Sec. 1.4 are well-approximated by a Maxwellian distribution (they are well-characterized by a single hot temperature), the kinetic results suggest that a fluid model is adequate for modeling the expansion unless other mechanisms such as electron stopping are significant.

Finally, the velocity distribution of fast ions is in practice roughly exponential (for initial Boltzmann electron distributions) with a slope that inversely scales with the hot-electron temperature (for an isothermal expansion the slope is $(ZT/2)^{-1/2}$, where $Z$ is the charge state of the ion).\(^{44}\) The spectrum has a sharp cutoff given by Eq. 1.37, and hence either the cutoff energy or the slope of the spectrum can be used to infer the hot-electron temperature. The exponential nature of the spectrum implies that there are far fewer ions in the tail of the fast-ion distribution in comparison to the bulk. A subtle consequence of this result is that the fastest ions act as test particles and their energy (the maximum ion energy) represents the peak path-integrated electric field (voltage) on the target.

Measurements of fast ions can thus be used in a number of ways. They can be used to infer the temperature of hot electrons that drive them, the peak target voltage, the (sometimes undesirable) energy loss to fast ions, as well as the number of hot electrons that escape the target (by charge balance). These properties are exploited in chapters 4-6, as discussed in the thesis overview below.

### 1.5 Thesis Overview

In this thesis, fast-ion measurements have been used to study various aspects of the central hot-spot ignition and fast-ignition ICF schemes. Spectral measurements of these ions were conducted using existing and new diagnostics that were in part developed and calibrated at the MIT Linear Electrostatic Ion Accelerator (LEIA).

Chapter 2 discusses recent upgrades to LEIA, which produces DD and D\(^3\)He fusion products used to develop and calibrate nuclear diagnostics. Hardware upgrades include the implementation of a new ion source, multi-channel analyzer with pre-amplifiers and surface-barrier detectors (for measurements of charged-particles) and imaging diagnostics (for characterizing the fusion products
source size). Software enhancements include beam-target physics simulations (for characterizing the energies of charged-fusion products) and reusable modular control software. The chapter concludes with a discussion of how these improvements have been essential to the development of charged-particle and neutron diagnostics, excluding new diagnostics developed for use in this thesis, which is deferred to the next chapter. Appendix A presents detailed design and performance specifications of LEIA, and appendix B discusses the development of a reusable modular control software package that was developed for use on LEIA and other scientific experiments.

Chapter 3 first presents an overview of existing fast-ion diagnostics and their limitations. The limitations of these diagnostics motivated the development of two Thomson parabola spectrometers, which are the main subject of the chapter. These spectrometers were developed in collaboration with Los Alamos National Laboratory (LANL) and State University of New York (SUNY), Geneseo to measure heavy fast-ions at Omega and MTW. The design and operating principles are discussed, followed by practical implementation details. Data reduction and analysis techniques of these instruments are discussed, and a software package that was developed to this end is summarized. In addition, a detailed study of the nuclear track detectors used in these diagnostics is presented appendix C.

Chapter 4 details spectral measurements of fast ions from direct-drive implosions on Omega. These ions, measured using existing and newly developed Thomson parabola spectrometers outlined in Chapter 3, are accelerated by TPD-generated hot electrons and indicate that ICF targets charge to megavolt potentials. The large potential traps energetic electrons that can preheat the fuel, raise the adiabat and degrade compression. An estimate of the amount of preheat in OMEGA targets was obtained using previous measurements of the total hot-electron energy and the fast-ion energy measured in this work. Refined preheat estimates will require simulation of the hot-electron dynamics, which will require knowledge of the temporal evolution of the target potential. The fast-ion measurements were used in conjunction with radiographic images of the target to build an empirical model of how the target potential evolves in time.

Chapter 5 discusses the first measurements of fast protons from integrated fast-ignition experiments that were conducted at Omega. Existing magnetic and compact proton spectrometers and were used to simultaneously measure protons at several locations around the implosion. These protons are accelerated by hot electrons generated by short-pulse LPI. Since the hot electrons are intended for core heating of fast-ignition targets, any energy coupling to ions is a loss mechanism. Spectral measurements of the protons were used to determine the amount of energy coupled to protons. In addition, measurements of the maximum proton energies were used to infer the temperature of the hot electrons. It was shown that the electrons are hotter than predicted by theoretical scalings, leading to poor coupling of hot electrons to the dense core of fast-ignition targets.

Chapter 6 presents measurements of fast protons and heavy ions from thin-foil targets irradiated at the MTW Facility. Generation of fast ions is integral to advanced fusion concepts such as proton fast ignition, and also has medical applications (e.g., proton therapy for cancer). The charge states and maximum energies of these ions were measured using a Thomson parabola spectrometer. It was shown that a wide range of heavy-ion species are accelerated from these targets and that the mean and maximum energies of the ions scale with the charge state. In addition, a two-temperature electron distribution was inferred from the proton spectrum and used to explain the observed discrepancies between the maximum proton and heavy-ion energies. These results have implications for fast ignition and medical applications, which require specific ion species for localized power deposition.

Chapter 7 concludes with a summary of the results presented in this thesis.
1.6 References

Upgrade of a Linear Accelerator for Development of Nuclear Diagnostics for use on Omega, Z and the NIF

2.1 Introduction

The development of nuclear diagnostics, including charged-particle diagnostics for fast-ion measurements, are essential for diagnosing ICF implosion performance. The development of these diagnostics at national laser facilities such as Omega or the NIF are costly in terms of the time and money required to debug and develop an instrument. To this end, the MIT Linear Electrostatic Ion Accelerator (LEIA) has been upgraded to support the development of diagnostics for Omega, the Sandia Z-machine and the NIF. This includes, but is not limited to, the diagnostics used throughout this work.

This chapter is organized as follows: Section 2.2 gives an overview of the various hardware components of the accelerator, along with a discussion of recent improvements; section 2.3 discusses the software development undertaken for both data acquisition, control and simulation capability. Section 2.4 discusses some recent development work on diagnostics for in use at Omega and the NIF. Section 2.5 summarizes LEIA upgrades and discusses the future direction of LEIA.

2.2 Overview and Hardware Development

The MIT LEIA\textsuperscript{1–4} (see Fig. 2.1) generates DD and D\textsuperscript{3}He fusion products for the development of nuclear diagnostics. In recent years, significant improvement have been made to the system. Fusion reaction rates, as high as \(10^{7}\) s\(^{-1}\) and \(10^{6}\) s\(^{-1}\) for DD and D\textsuperscript{3}He, respectively are now well regulated with a new ion source and electronic gas control system. The implementation of a new ion source and gas control system has allowed for better control of the spatial and temporal characteristics of the fusion-product source, which is essential for the calibration of a number of nuclear diagnostics. In addition, the development of a CCD-camera based target viewing system (TVS), combined with modeling of neutron scattering within the target chamber, has allowed for detailed characterization of the fusion-product source; this has further facilitated the development of neutron yield diagnostics. Finally, a newly-implemented multi-channel analyzer (MCA) with advanced signal-processing algorithms has been used to measure charged-fusion products more accurately and with greater precision; these benefits directly propagate to the energy calibration of charged-particle spectrometers.
A complete schematic of the upgraded LEIA is shown in Fig. 2.2. The upgrades include a new radio-frequency (RF) driven positive ion source, an open-air high voltage deck, reconfigured beamline, and target chamber. The newly implemented ion source, manufactured by National Electrostatics Corp. (NEC),\(^5\) is capable of producing 200 \(\mu\)A deuteron beams and 170 \(\mu\)A \(^3\)He ion beams. A terminal voltage of 150 kV, generated using a Cockroft-Walton Multiplier (not shown in the schematic), accelerates these ions onto a target downstream. The target, composed of a copper substrate with a thin film of ErD\(_2\), is loaded with either D or \(^3\)He to allow production of either DD or D\(^3\)He fusion products.

### 2.2.1 Implementation of a New Ion Source

A new ion source, schematically shown in Fig. 2.2, was acquired and implemented to enhance beam control and stability. The DD (or \(^3\)He) plasma discharge is generated within the source using a capacitively-coupled 300 watt RF oscillator operating at 100 MHz. The source also includes permanent magnets and an Einzel lens assembly. The plasma is first compressed by magnetic fields at the entrance to an aluminum canal (1 or 2 mm diameter) where it is electro-statically extracted and focused to a downstream target. The source may be biased using up to three power supplies for operation, excluding the RF oscillator power supply: one each for the focus, the extractor and the probe electrodes; these supplies are referenced to the terminal voltage (deck bias). Note the probe and extractor electrodes, shown in Fig. 2.3, located upstream and downstream of the plasma bottle, respectively. In these types of sources, the purpose of the probe power supply is to drive ions out of the source by maintaining a potential difference between the extractor and the probe electrodes; for this reason it is typically referenced to the extractor bias using an isolation transformer. This approach was not taken as it does not allow a single digital electronic controller to readily control and monitor the supplies because they have large (several kV) DC offsets between them. The approach taken here was to reference all three supplies to the terminal voltage and to compensate for this by adding an offset voltage to the probe output equal to the instantaneous extractor bias. This output tracking between supplies was achieved in software. Nevertheless, any reference to the
Figure 2.2. Top and side view of LEIA. Shown are the target chamber, beamline and ion source along with supporting components. Electrostatic components include a focusing lens and a linear accelerating tube. Beam diagnostics include a Faraday cup, beam profiler and a residual gas analyzer. The vacuum system is comprised of two turbopumps and three dry vacuum (scroll) pumps.

Probe bias in this work refers to a potential difference between the extractor and probe electrodes. Three supplies were obtained from Glassman HV Inc. for this application: two MK-series supplies, with 15 kV and 20 kV outputs for the focus and probe, respectively, and an MJ-series power supply with an output of 5 kV for the extractor. These supplies are housed in the source supply box shown in Fig. 2.2. That box also incorporates an Acromag ES2152 fiber-optic-coupled controller, which drives and monitors the analog interfaces of the supplies while providing a single digital, fiber-isolated interface to the control computer. The bias conditions depend strongly on the electrostatic optics of the entire system. For this system, an extractor bias of 0-5 kV, a focus of ~3-10 kV and a probe of ~3-10 kV proved to be sufficient for extracting and focusing a beam of ions down to a diameter of 3 mm approximately 2.5 m downstream of the ion source. It was found that the optimal probe and focus values depend on the extractor bias. The extractor, despite its name, does not actually drive ions out of the source. It biases the entire source relative to the terminal
Figure 2.3. Schematic of the NEC RF Positive Ion Source, including the probe, extractor and focusing lens assembly. Shown below the source is a sample bias profile. The allowable range of values for each of the bias supplies are indicated and may be adjusted independently.

The stability of the fusion reaction rate depends strongly on the stability of the ion beam current; the latter is a function of the ion source bottle pressure. The nominal flow rate of gas into the source is approximately 0.02 sccm/s with a nominal fill pressure in the range of 10 – 30 mTorr. It was found that momentary fill pressures greater than 50 mTorr were often required for proper startup of the plasma and that the optimal fill pressure during operation range extended out to approximately 40-60 mTorr. These required flow rates of D are quite low for modern thermal mass-flow controllers. Thus, ion sources of this type often utilize manual or motor-driven variable-leak valves for gas pressure control, whereby the valve orifice is manually set to a fixed position (and adjusted manually as required). Since some diagnostic applications require reliable and steady beam currents over extended periods of time (as discussed in Sec. 2.4), a fast-response feedback-controlled gas-control system was designed and implemented.
The gas control system consists of a Horiba STEC piezoelectric flow control valve\(^8\) and an MKS 626-series Baratron Capacitance Manometer\(^9\) for direct measurement of the ion source fill pressure. The same fiber-optic-coupled controller was used to readout the fill pressure and drive the piezoelectric valve based on commands it received from a software implemented proportional-integral-differential (PID) control loop. The loop allows control of the pressure to within 0.5 mTorr. With proper tuning of the loop, the step response time of the system is < 2 seconds. Since the maximum valve orifice leads to over-pressurization of the source, the software package implements a valve calibration feature to prevent the valve from being opened too far. This feature is essential since over-pressurization of the source slows down the response time of the gas control system.

The new gas control system has helped to stabilize the beam current and hence the fusion reaction rate. The fusion reaction rate, as inferred by a fixed detector with a finite solid angle, is a function of beam current, beam energy, on-target beam position and the extent to which the target is loaded. We observe, as shown in Fig. 2.4, small variations in the DD fusion reaction rate over fine timescales. These variations are within the statistical uncertainty of the count-rate measurement ($\sigma = \pm 2\%$). More importantly, no significant long-term (of order 10 - 20 minute) trends are seen in the data when feedback control of the gas pressure is utilized. Various parts of the system, including the ion source canal and target, undergo thermal expansion as the source and beam are operated for long duration; any long-term changes in the count-rate associated with system components reaching steady-state operating parameters are mitigated with the use of feedback. For these data, the ion source was allowed to warm-up for several hours per NEC’s standard operating procedure; a residual gas analyzer (RGA) was used to verify that the beam predominantly consisted of deuterium. Characterization and control of the fusion reaction rate for long time scales is particularly important for development of neutron yield diagnostics based on activation of materials, as discussed in Sec. 2.4.

### 2.2.2 Targets and Target Viewing System

The targets used for LEIA were manufactured by Sandia National Laboratories.\(^10\) The active layer of the targets consists of erbium-deuteride (ErD\(_2\)) with a diameter of 1 cm and a nominal thickness of 5 $\mu$m, corresponding to the range of a 150 keV \(^2\)H\(^+\) ion. Although the thickness is sufficient for stopping 150 keV deuterons, it was found that these thin targets deteriorate after approximately 20 hours of beam-on-target time for ion beam powers in the range of 15-20 Watts. The thin layer of ErD\(_2\) is ablated off, exposing the copper substrate. Future targets will incorporate thicker active layers for increased lifetime and durability. Other target materials have also been tested, such as titanium and it was found that these targets produced DD fusion rates significantly lower than...
Figure 2.5. Images of a 140 keV deuteron beam incident on an ErD$_2$ target (a) with ambient lighting and (b) without ambient lighting. Shown in these images are the water-cooled target holder (constructed of copper) and the circular target itself (approx. diameter of 1.1 cm). The ion source was biased with a probe voltage of 8.7 kV, extractor of 5 kV and focus of 9 kV to achieve a focal spot size of approximately 3 mm (dashed circle) on a target 2.5 meters downstream of the ion source.

A new target viewing system (TVS) diagnostic was implemented for in-situ measurements of the on-target beam profile. The target camera, shown in Fig. 2.2, is a CCD-based network camera (Axis Communications Model 221), which provides a more accurate measurement of the fusion product source size relative to that of the beam profiler. The beam profiler samples the beam cross-section and position well upstream of the target (see Fig. 2.2) where it is typically broader as it is converging onto the target. The visible light self-emission of energetic deuterons exciting the target medium is sufficient to generate an image without the need for background lighting. The camera utilizes a CCD with a sensitivity of 0.65 lux, which provides sufficient sensitivity for this low-light application. Images of a 140 keV deuteron beam incident on target were taken with the TVS, with and without ambient lighting as shown in Figs. 2.5a-b. These images were taken during a single run where the beam was focused to a diameter of 3 mm; the source was operated at a fill pressure of 40 mTorr and biased with an extractor of 5 kV, focus of 9 kV and probe of 8.7 kV. Such in-situ measurements of the beam profile allow the operator to point and focus the beam for each run with greater accuracy than one might achieve with a beam profiler alone. Precise knowledge of the fusion source size and position with respect to target chamber center (TCC) is also essential for the calibration of several diagnostics, as discussed in Sec. 2.4.

2.3 Software and Simulation Development

2.3.1 Control, Monitoring and Logging

The upgrades and modifications to accelerator hardware require a software-based control solution, which is scalable and modular, allowing components of the control software to be re-used, modified or removed entirely in response to hardware changes. To this end, a novel modular and extensible toolkit was developed for control of the various accelerator subsystems. This Modular Control Toolkit (MCT) was written in C++ and uses open-source libraries for its graphical user interface; it currently supports both 32-bit and 64-bit UNIX-like systems. The toolkit itself consists of a central console, a module manager, an interlock engine and a shared library with templates for building modules. The shared library implements the base code which is common to all modules and is hence shared between all running modules. Using the toolkit, one only needs to write a module with the minimal code required to communicate with hardware specific to a given experiment. The development, architecture and usage of the MCT is further discussed in Appendix B.

Several new modules were written to control the accelerator, including the ion source controller, the vacuum valve controller, the turbopumps, ion gauge controller and other system-level compo-
2.3 Software and Simulation Development

Full electronic control, monitoring and logging of the system parameters (e.g., voltages, currents, pressures and temperatures) is now possible.

Improvements were also made to the way system and run parameters are logged. An SQL-based\textsuperscript{14} database was implemented to hold run data along with useful information about charged-particle detectors, targets and other system-wide parameters. The terminal voltage, source bias, fusion count rates, fill pressures and charged-particle data are stored for each shot and may be queried over a web interface using any number of fields. This new capability is extremely useful for quickly retrieving run data. The database also serves as an essential aid when debugging the system or resolving anomalies in data, as one has a reference of pertinent system parameters for each experimental run.

2.3.2 Charged-Particle Diagnostics Suite

The primary data-acquisition system consists of a newly implemented NIM-based setup with four signal chains for the measurement of charged-particle spectra. Surface barrier detectors (SBDs) are routinely used for direct measurements of the energy spectra of charged fusion products. This data acquisition system consists of a model N1728B multichannel-analyzer, obtained from C.A.E.N.,\textsuperscript{15} a four-channel preamplifier, and custom software developed in-house, hereon referred to as the C.A.E.N. MCA Application (CMA). The open hardware specifications and implementation details provided by C.A.E.N. have enabled us to acquire charged-particle spectra with greater accuracy and better precision. All aspects of the signal processing chain, from initial analog-to-digital conversion to post-processing and deconvolution of the preamplifier impulse response function (IRF) are controlled using custom software written in Java. Deconvolution of the preamplifier response is done in real-time in hardware (as originally implemented by C.A.E.N.) using a well-known algorithm.\textsuperscript{16} The MIT-developed software package allows the user to specify parameters for the algorithm, trigger and acquisition; simultaneous acquisition of energy spectra and oscillograms of single-particle events are also possible. Furthermore, one may use the CMA to record count rates over time (i.e., software scaler mode), to perform in-situ analysis, including data fitting, and to calibrate and store channel-to-energy mapping data; the stock open-source Java-based software furnished by C.A.E.N. does not implement these latter features, though it was helpful in developing this software suite.

Of particular importance in the new system is the deconvolution of the preamplifier IRF. Ideally, the SBD-preamplifier combination will have a fast-rising “impulse” with an amplitude that is linearly proportional to the energy of the incident particle; the SBD depletion depth ($\sim 2000\mu m$) is sufficient to stop the charged-particle. In practice, the SBD takes a finite amount of time to sweep out the charge (the electron-hole pairs) generated by the incident particle. The charge sweep-out time increases with incident particle energy, leading to an impulse amplitude that is systematically lower for higher energy particles. Thus, the incident particle energy is underestimated for more energetic particles. This is problematic for 14.7 MeV $^3$He protons in the laboratory since calibration of the SBD itself is typically accomplished using low-energy $\alpha$-particles from the decay of heavy isotopes, such as $^{226}$Ra.

The energy calibration of the MCA is defined by a linear mapping between channel and incident particle energy. Four $\alpha$-particles from a $^{226}$Ra source, with energies in the range of 4 - 8 MeV\textsuperscript{17}, are measured using the MCA. The energies of these particles as a function of measured channel, as shown in Fig. 2.6, are then fit to the form

\[ E_M = a_1 + a_2 \times N_C \]  

where $N_C$ is the MCA channel, $E_M$ is the energy of a particle as measured by the MCA, and $a_1$
and $\alpha_2$ are the fit coefficients. This calibration fit is extrapolated to both lower energies (e.g., 3 MeV DD-p) and higher energies (e.g., 14 MeV D$^3$He-p). The linearity of the MCA is quantified by computing the residual of the fit, defined as

$$R \equiv E_I - E_M$$  \hspace{1cm} (2.2)$$

where $E_I$ is the known energy of the $\alpha$-particle and $E_M$ is the energy of the particle as measured by the calibrated MCA. The residual is thus a measure of how much a measured particle energy deviates from the expected particle energy.

These residuals, along with residuals of the 14 MeV and 3 MeV fusion products, are shown in Fig. 2.7 for two cases: (a) the old MCA system, which does not deconvolve the preamplifier IRF, and (b) the new MCA, which performs hardware deconvolution in real-time. The deconvolution algorithm essentially linearizes the impulse height output of the SBD with respect to incident particle energy, resulting in more accurate measurements (lower residuals). The residuals of the fusion products at 3 MeV and 14.7 MeV are computed using the difference between simulated and measured values of the proton energies at 90° with respect to the ion beam, as shown in Fig. 2.8. For such charged-particle measurements, a thin aluminum filter (6 $\mu$m) is used to block SBD from the large flux of low-energy (< 140 keV) scattered beam ions to prevent the SBD from damage. The measured energy spectra are corrected for the charged-particle ranging through this filter. One significant component of the residuals at 3 MeV, and to a lesser extent, at 14.7 MeV, are the uncertainties of the expected energy of fusion products incident on the SBD (at 90°). These energies

Figure 2.7. Comparison of the new and old MCA systems, with and without deconvolution of the preamplifier IRF, respectively. Shown are the residuals of the calibration fit as a function of incident particle energy. The MCA is calibrated to the $^{226}$Ra-$\alpha$’s (4-8 MeV) and this calibration is then extrapolated to both lower and higher energies. The residuals of the calibration for 3 MeV and 14.7 MeV fusion-protons are significantly lower when the preamplifier IRF is deconvolved in the measurement.
are computed by adjusting the birth energies of the fusion products for beam kinematics (discussed in the next section) and ranging down the particle from its birth energy through several microns of the target material, ErD$_2$. The latter correction is complicated by the fact that one must know where reactions occur within the target. This depth is computed in a simulation in which particles are ranged through the appropriate amount of ErD$_2$ as they leave the target at a given angle. The energies are simulated to within an accuracy of 10-20 keV. These simulations are discussed in the following section. Another component that adds to the residuals of the fusion products is the precision of the energy measurement, which depends on experimental factors (e.g., how well the SBD position is known); these factors lead to a measurement precision of 20 keV. The major component of the residuals of the 3 MeV and 14.7 MeV fusion products are attributed to these two components, and the total systematic uncertainty of charged-particle energy measurements with this system is ±50 keV.

### 2.3.3 Beam-Target Physics Simulation

Predictive capability of the energy and fluence of fusion products is important for verifying the calibration of the charged-particle diagnostics suite and the associated-particle technique, where one relies on in-situ measurements of the fluence of 3 MeV protons to determine the fluence of the associated 2.45 MeV neutrons from the DD fusion reaction. To this end, a code was developed to simulate beam-target physics, including the slowing down of charged-particle reactants and products, relativistic reaction kinematics, and the differential angular cross-section for the fusion reactions of interest. The simulation takes beam energy, current, species and cross-sectional area as inputs. The beam cross-sectional area may be approximated using either the beam profiler or more accurate CCD data of beam area on target (obtained from the TVS) for a given run. The beam current and area are used to determine fusion reaction rates; the former may be measured using the Faraday cup. Note that in lieu of a beam velocity selector, a residual gas analyzer is used to verify that the measured beam current is dominated by the ions of interest (typically $^3$He$^+$ or D). Outputs of the simulation for practical cases are shown in Fig. 2.9. These simulations were made for 100 keV, 120 keV, 140 keV and 170 keV deuteron beams incident on an ErD$_2$ target. The outputs are normalized to the total number of fusion counts per second to illustrate relative differences between beam energies.

As shown in Fig. 2.9a-c, the fluence of DD protons and neutrons peak towards the forward
Figure 2.9. Simulated results for a beam of deuterons (with four different energies) incident on an ErD$_2$ target. The target forms an angle of 45° with respect to the beam, consistent with the lab setup shown in Fig. 2.8. Shown are (a) normalized DD-p counts per steradian as a function of laboratory angle (b) normalized DD-n counts per steradian as a function of laboratory angle (c) the DD-n/DD-p count ratio as a function of laboratory angle (d) normalized DD fusion reaction rate per $\mu$m as a function of depth into target and (e) Orientation of the beam and target inside the target chamber for these simulations.

Figure 2.10. Measured energy spectra for (a) an Americium-241 source (b) the LEIA DD proton source and (c) the LEIA D$^3$He proton source. The LEIA proton data were acquired with a 6 $\mu$m Al filter placed in front of the SBD. The amount of line broadening due to this filter is small ($\sim 20$ keV) but the downshift is significant and must be taken into account for characterization of the mean proton energies. The americium source is used to infer the broadening due to the SBD, preamplifier and MCA. This instrument linewidth is then deconvolved from the measured linewidth of the fusion protons, shown here, in order to determine the linewidth of the fusion products source itself.
beam direction (0° laboratory angle). The forward peaks in these plots are the result of the forward momentum introduced into the system by the deuteron beam. The ratio of DD-n to DD-p is essential for any associated-particle measurements. For these types of measurements, the DD neutron fluence at a given location in the target chamber is inferred by measurement of DD protons using an adjacent SBD. Correction factors must be applied between different laboratory angles to infer neutron yields properly. Corrections must also be applied for neutron scattering in the target chamber, as discussed in the next section. This technique has been used for neutron diagnostic development, as discussed in Sec. 2.4. Fig. 2.9d shows the normalized DD reaction rate as a function of target depth; as expected, reactions occur within a few microns corresponding to the range of the deuterons in ErD₂.

The beam-target physics code has also been used to obtain the expected energies of the charged fusion products at a given location in the target chamber and to understand the sources of line broadening for the 3 MeV and 14.7 MeV fusion-protons. The former is used to verify the MCA calibration, discussed in the previous section, while the latter is important for characterization of the instrumental linewidth of diagnostics being developed. The linewidth as measured by the SBD is due in part to the instrument response of the charged-particle diagnostics suite (SBD + preamp + MCA combination) and also due to beam-target physics. These two components must be quantified to determine what the actual incident linewidth is. This has been achieved by exposing the SBD to α-particles from a ²⁴¹Am source, as shown in as shown in Fig. 2.10a. The amount of broadening introduced by the charged-particle suite is ∼70 keV (FWHM). Shown in Figs. 2.10b-c are the energy spectra of DD and D³He protons, respectively, as measured by the SBD at a distance of 14.5 cm from the fusion product source and an angle of 90° with respect to the 140 keV deuteron beam.

The mean energies of DD and D³He protons, after they exit the target at 90° with respect to the beam, have been measured to be 3.06 ± 0.04 MeV and 14.58 ± 0.04 MeV, respectively. This was achieved by fitting the measured spectra shown in Figs. 2.10b-c and then up-shifting the mean energies from the fit through the 6 μm Al filter (placed in front of the SBD for these measurements). The mean values predicted by the beam-target simulation are 3.04 MeV and 14.67 MeV, which are within the statistical and systematic uncertainties associated with the measurement and calibration of the charged-particle suite. The actual linewidth of the fusion product source is the measured linewidth (Figs. 2.10b-c) after the instrumental linewidth has been deconvolved (Fig. 2.10a). After correcting for the instrumental width of ∼ 70 keV (FWHM), the linewidth of the DD and D³He fusion-products source are ∼150 keV and ∼140 keV (FWHM). The beam-target simulation is able to account for approximately half of the source linewidth. Though it simulates kinematics and ranging of beam ions and fusion products in the target, the simulation does not include energy straggling or finite source size effects. We attribute the remainder of the unaccounted linewidth to these effects.

2.3.4 MCNP Simulations

The development of neutron diagnostics, in particular yield diagnostics, requires a thorough characterization of scattering effects in the target chamber. The Monte-Carlo simulation code MCNP was used for this purpose. An accurate model of the target chamber itself, including all significant sources of scattering within the chamber, was used for this simulation; materials and geometry were specified from solid models. Simulations were conducted assuming an isotropic point source of DD neutrons with a birth energy of 2.45 MeV in the target. Figures 2.11a-b show a top and side view of the target chamber, with contours of normalized neutron fluence scaled by 4πR²; deviations from unity thus represent a divergence from a 1/r² scaling. The black lines shown in the figure
**Figure 2.11.** MCNP simulation of neutron scattering within the LEIA target chamber showing (a) top view and (b) side view of the chamber. Shown are contours of normalized neutron fluence scaled by $4 \pi r^2$; deviations from unity represent a divergence from a $1/r^2$ scaling. The black lines indicate the water-cooled target. Significant correction factors must be applied to neutron yields inferred from diagnostics. At a distance of 10 cm from the target a correction factor of approximately 0.80 needs to be applied to the measured fluence to correctly infer the neutron yield from a local measurement.

It is clear from these simulations that scattering corrections are significant, even for a large cylindrical target chamber (60 cm × 15 cm). The target, which is angled at 45°, casts a shadow along the entire chamber where the neutron fluence is significantly lower; the fluence generally increases as one approaches the target chamber wall, where it is enhanced by nearly 50%. On the non-shadow side, where neutron diagnostics are typically placed for calibration, significant corrections still need to be applied to experiments. Even at a distance of 10 cm from the source, the neutron fluence is enhanced by 20%. Neutron diagnostics are calibrated a few centimeters from the source. This choice minimizes the scattering correction, but is sufficiently far from the source that uncertainties in the absolute source position, typically 2 mm, are insignificant.

These detailed simulations of neutron scattering have been essential for the development and calibration of a CR-39-based neutron yield diagnostic used at both Omega and the NIF, as well as an indium activation-based yield diagnostic developed by Sandia National Laboratories for use at Omega, Z and the NIF.

### 2.4 Diagnostics Development

A number of nuclear diagnostics have been developed, tested and calibrated using LEIA. In the past, LEIA has been primarily utilized for characterization of CR-39 solid-state nuclear track detectors, which form the basis for many nuclear diagnostics; these works have led to a number of publications...
on the CR-39 response to charged-particles\textsuperscript{21–23} and coincidence counting with CR-39.\textsuperscript{24} Some aspects of the CR-39 response to protons are also discussed in Sec 3.2 and in Appendix C. Over the last several years, an increasing amount of time has been spent calibrating diagnostics to better precision and developing other types of diagnostics; each of the diagnostics are briefly outlined in what follows.

### 2.4.1 Wedge-Range-Filter Spectrometers

Wedge-Range-Filter (WRF)\textsuperscript{25} proton spectrometers are composed of an aluminum or zirconia wedge positioned onto a piece of CR-39. The WRF spectrometers, capable of measuring proton spectra in the energy range of 4 MeV - 20 MeV, have been in use at Omega\textsuperscript{26,27} for a number of years and have more recently been used on the NIF.\textsuperscript{28} These spectrometers are routinely used to measure primary and secondary fusion yields, shell $\rho R$ from the downshift of charged fusion products and fuel $\rho R$ from scattered fuel ions (“knock-ons”).

After manufacture, the WRFs are calibrated and performance tested using the accelerator before they are sent out for use at Omega and the NIF. WRF proton energy spectra acquired on LEIA is shown in Fig. 2.12. WRFs are also periodically tested for surface degradation between shots on the NIF. The aforementioned upgrades to the charged-particle diagnostics suite have recently allowed for more precise calibrations of the WRFs. Uncertainties in the energy calibration of WRFs are in the range of $\pm 150$ keV and this has enabled new physics studies. For example, using multiple WRFs, it is now possible to study the $P_2/P_0$ $\rho R$ asymmetry mode in ICF implosions, in-flight, to a precision in $P_2/P_0$ of $\pm 0.07$. Techniques are also being developed to measure the fuel ion temperature ($T_i$) using the line-width of charged-fusion products. The line width, as measured by the WRF, consists of several components, one of which is the Doppler broadening due to finite $T_i$. To extract this component from the WRF data, the instrumental broadening due to the WRF itself must be well characterized. These studies are being carried out on LEIA. Further details of the operation and limitations of WRF spectrometers are discussed in Sec. 3.2.1.
2.4.2 CR-39-based DD-n Yield Diagnostic

A novel, CR-39-based DD-neutron yield diagnostic, developed on LEIA and tested on OMEGA, measures absolute DD yields as well as the fraction of neutron backscatter at a given location. The diagnostic (Fig. 2.13a) consists of a piece of CR-39 partially covered by a foil of 100 μm thick polyethylene (C₂H₄), which serves to enhance the neutron-induced proton signal on the CR-39 directly behind the foil. The uncovered CR-39 serves as a background region for subtraction of both intrinsic background and recoil protons produced in the CR-39 itself. The background-subtracted signal from the polyethylene-covered region of the CR-39 therefore measures only recoil protons generated in the polyethylene, and is directly proportional to the absolute neutron fluence. Furthermore, because only polyethylene-generated protons are detected, the detector is sensitive only to DD neutrons incident from the front, a distinct advantage over other neutron detectors that are susceptible to backscatter. By placing polyethylene behind as well as in front of the CR-39, it is possible to measure the relative amount of neutron backscatter (by comparing the polyethylene-produced signal on the back side to the polyethylene-produced signal on the front side).

The neutron detection efficiency and background subtraction methods used for this diagnostic were developed on LEIA using the associated particle technique, whereby the absolute DD-neutron fluence at the detector is inferred from the measured DD-proton fluence at an adjacent SBD. Sample data taken on LEIA is shown in Fig. 2.13b. In the top right corner of the figure is an image of the CR-39 detector, with the left side covered with 100 μm polyethylene. In this image, darker pixels represent a greater number of recoil proton tracks. Alongside the image is a lineout of proton fluence across the detector, which shows the enhancement of recoil protons due to the polyethylene. The neutron yield corresponding to this data is inferred from the adjacent SBD measurement, which must be corrected for neutron scattering, finite source size and kinematic differences in the DD-
p and DD-n yields. Thus, precise characterization of the fusion products source, as determined by the beam-target physics simulations, MCNP characterization of neutron scatter in the target chamber and the TVS diagnostic are central to development of this diagnostic. Several aspects of the diagnostic are still being developed, including directionality of incident neutrons, effects of prolonged exposure to vacuum, and variability of CR-39 response to protons.

2.4.3 Indium Activation Neutron Yield Diagnostic

The indium activation neutron yield diagnostic relies on the activation and subsequent decay of the activated material to infer neutron yield at Z and the NIF. Indium “slugs” are activated during experiments by an unknown number of DD or DT neutrons through the reaction $^{115}$In$(n,n')^{115m}$In. The meta-stable $^{115}$In will then decay by emitting 336 keV gammas (the threshold energy for activation), which are measured using a high-purity germanium (HPGe) detector. A calibration factor, known as the $F$ factor, relates the measured gamma yield to the total neutron yield.

The $F$ factor is generally obtained through experimental calibration and encompasses detector efficiencies. Note that on Z, one must also consider the competing reaction, which stems from the strong x-ray background: $^{115}$In$(\gamma,\gamma')^{115m}$In. Calibration of this diagnostic on LEIA is thus unique since $F$ factors may be obtained for an x-ray free environment.

Preliminary experiments on LEIA were used to determine the $F$ factors for a number of indium samples. Samples of various sizes were activated within approximately 30 minutes at DD reaction rates of about $10^6 \text{ s}^{-1}$. The samples were then removed from the target chamber within minutes after the shot; the gammas were then counted using the HPGe. As in the case of the CR-39-based DD-n Yield Diagnostic, and as discussed in Secs. 2.2.2 and 2.3.4, the target viewing system and simulation of neutron scattering were essential for precise determination of the $F$ factor. In addition, stability of the fusion reaction rate is particularly important over these timescales. This is because the indium is continuously bombarded with neutrons for extended periods of time on LEIA. If the neutron flux is not stable, it is difficult to assign an $F$ factor to the diagnostic. As discussed in Sec. 2.2.1, recent improvements to the gas control system have stabilized the reaction rate significantly and will benefit future $F$ factor experiments.

2.4.4 Particle Time-of-Flight (pTOF) Diagnostic

The particle time-of-flight diagnostic (pTOF) is used at OMEGA and the NIF to measure the D-$^3$He shock-bang and DD compression-bang time by resolving fusion protons and neutrons to an accuracy of ±150 ps. The diagnostic consists of a chemical-vapor-deposited (CVD) diamond, biased to several hundred volts, and a filter in the front of the diamond to reduce the large x-ray background present in indirect-drive implosions at both OMEGA and the NIF. These measurements, when combined with the shock $\rho R$ as measured by WRFs, strongly constrain implosion models. Improvements to the overall accuracy of the diagnostic are therefore essential to the ignition effort.

One improvement that will be implemented at the NIF is to increase the bias voltage on the detector. In addition to improving the response-time of the diagnostic, saturation effects that may be caused by the large hohlraum x-ray signal will be reduced. Various bias voltages are being tested on the accelerator before such a capability is implemented at the NIF.

Critical to the interpretation of the pTOF data is a thorough understanding of the instrument response function and sensitivity to protons. Single-particle and integrated charge studies are being conducted on the accelerator by ranging down 14.7 MeV protons to energies of interest. The charged-particle diagnostics suite is essential in calibrating pTOF to the SBD during integrated charge studies. Moreover, the accelerator does not generate hard x-rays, allowing these studies to
be carried out without an x-ray background.

## 2.5 Summary

The MIT Linear Ion Accelerator (LEIA) has undergone several upgrades, which allow it to develop advanced diagnostics for Omega, Z and the NIF. Implementation of a new ion source and custom gas control system now provides better regulation and improved stability of the fusion-product source. The addition of a new charged-particle data acquisition system has allowed for more precise energy characterization of the fusion products and therefore better calibration of charged-particle diagnostics. In-situ measurements of the on-target beam profile, made with a CCD camera, together with simulations of neutron scattering within the target chamber, have facilitated the development of neutron yield diagnostics. These improvements allow better support of existing diagnostics and the development of new diagnostics in aid of the national program.

## 2.6 References

9. MKS Instruments, see http://www.mksinst.com/.
13. GIMP Toolkit, see http://www.gtk.org/.
15. C.A.E.N., see http://www.caen.it/.
19. Monte Carlo N-Particle (MCNP) Transport Code, see http://mcnp-green.lanl.gov/.
3

Charged-Particle Diagnostics for Fast-Ion Measurements on Omega and MTW

3.1 Introduction

Charged-particle diagnostics have been used throughout this work to measure energy spectra of fast ions generated in ICF implosions and laser irradiation of flat-foil targets. The fast ions diagnosed in this work have a wide range of energies, charge states and yields. Thus, different types of diagnostics are required for these measurements depending on the experimental goals.

In this chapter, existing charged-particle diagnostics and their limitations are discussed first in section 3.2. The limitation of these diagnostics motivated the implementation of Thomson parabola spectrometers, outlined in section 3.3.1. Two Thomson parabolas were developed, calibrated and fielded in collaborations with State University of New York, Geneseo and Los Alamos National Laboratory, as presented in sections 3.3.2 and 3.3.3. These diagnostics were used in experiments where the existing ones were deficient.

3.2 Existing Diagnostics and Limitations

Proton, deuteron, triton and alpha-particle energy spectra are routinely measured on OMEGA using both magnet-based charged-particle spectrometers (CPS1 and CPS2) and Wedge-Range-Filter (WRF) spectrometers. These instruments utilize CR-39 solid-state nuclear track detectors (SSNTD).

As charged-particles pass through CR-39, they leave behind a trail of damage. To reveal these trails of damage, the CR-39 is etched in a heated alkali solution (such as NaOH), which etches away the damage sites faster than the bulk material. A lot of information can be extracted from the characteristics of a track, as discussed below. The response of CR-39 to neutrons has also been characterized\cite{1,2} and it has been shown that the detection efficiency of D-D and D-T neutrons is significantly lower (10^{-4} and 10^{-5} for D_{2} and DT neutrons, respectively) than that of charged-particles. In addition, CR-39 is immune to electro-magnetic-pulse (EMP) and to some extent x-rays, making it ideal for measurements of charged-particles in the presence of harsh neutrons and x-ray environments, such as those presented in Chapters 5-6.

It is well-known that CR-39 itself contains information about the energy and species of the detected charged-particles. In particular, the diameter and contrast of a particle’s track are routinely used to identify low-Z particles (protons, deuterons, tritons and alpha-particles). In addition, for a given particle species, the track diameter at the surface of the detector is a function of the incident
particle energy. This is because the rate of damage left behind as a particle penetrates and travels through the CR-39 detector is related to the particle’s instantaneous energy loss. Since a particle’s energy loss in CR-39 is related to the energy of the particle \( (dE_p/dx \sim 1/E_p) \), the CR-39 by itself can in some situations be used for coarse spectral measurements.

It has been shown, however, that there is a large amount of piece-to-piece variability in the dependence of track diameter on particle energy,\(^3\) as shown in Fig. 3.1 for protons. The curves were acquired using \( \text{D}_2 \) and \( \text{D}^3\text{He} \) protons on LEIA, which were ranged down to the energies shown here (see Appendix C for more information). Eleven pieces of CR-39 were characterized using this method (Fig. 3.1a), alongside the mean and standard deviation of these curves (Fig. 3.1b). Given the rather large variations shown, it is difficult to utilize CR-39 by itself for direct spectral measurements.

Further complicating the matter is the fact that the CR-39 detection efficiency of protons falls from 100% at 4-5 MeV to zero at approximately 10 MeV, because the energy loss of light energetic particles is too low for optical detection of tracks.\(^1\) For these reasons, CR-39 alone cannot be used for accurate measurements of charged-particle spectra above 4-5 MeV and must be paired with additional particle ranging and dispersion mechanisms for accurate spectral measurements of ions.

### 3.2.1 Wedge-Ranged-Filter (WRF) Spectrometers

The Wedge-Range-Filter (WRF) Spectrometers are passive instruments consisting of CR-39 overlaid with a wedge-shaped material.\(^1\) The wedge, constructed of either aluminum or zirconia ceramic (\( \text{ZrO}_2 \)), disperses particles along the CR-39 based on the ranging of the particle in the wedge medium (see Fig. 3.2). The combination of wedge-shaped material and CR-39 set limits on the minimum and maximum detectable charged-particle energies at each location along the wedge. Along this dispersion direction, particles must be energetic enough to penetrate the wedge filter, but not overly energetic so that the track can be registered onto the CR-39. Thus, at each “bin” along the dispersion direction, there exists a narrow energy range over which particles are detected. Within each bin, the diameter of a track, combined with knowledge of the local wedge thickness, is used to infer that particle’s energy. The minimum and maximum energies at each bin are known (by design) and correspond to the largest and smallest track diameters. Thus, one only requires
knowledge of the shape of the CR-39 response curve, and not the absolute value, to interpolate between these points. This feature makes the WRFs relatively immune to piece-to-piece variations in the absolute CR-39 response, as the absolute response curve is inferred from the data itself.

As discussed briefly in Chapter 2.4, these spectrometers are routinely fielded at OMEGA for the measurement of fast ions, primary and secondary charged fusion products, and scattered fuel and shell ions (“knock-ons”), which are elastically scattered by primary fusion products. Aluminum WRFs are more robust than their zirconia counterparts, but cannot be manufactured thin. This limits the proton low-energy cutoff to between 4-7 MeV. Zirconia WRFs, on the other hand, are fragile but can be made much thinner. As a result, the low-energy instrument cutoff for measurement of protons is approximately 3 MeV. The high-energy cutoff of the WRFs is typically around 20 MeV (protons) and is set by both the CR-39 detection efficiency of protons and thickness of the thick-end of the wedge.

The WRFs are inexpensive, compact (5 cm across) and thus ideal for probing targets at close distances and several locations simultaneously. Several (either 3 or 5) WRF modules, each consisting of two WRFs (or a single WRF in the case of aluminum), can be used at a single measurement location to obtain good statistics. WRFs are fielded in the OMEGA target chamber within interchangeable diagnostic manipulators referred to as ten-inch-manipulators (TIMs). The TIMs both define a space envelope that houses a compatible diagnostic and provide a power and control interface to the diagnostic if necessary. In contrast to diagnostics fixed on the OMEGA target chamber, the distance and pointing of TIM-based diagnostics, such as the WRFs, can be adjusted from shot-to-shot.

Although WRFs spectrometers are well-suited for measurements of protons and deuterons, they cannot be used for spectral measurements of heavy ions for a couple of reasons. The WRFs rely on charged-particle track diameter and contrast on the CR-39 for the discrimination of ion species. Though the differences in contrast and diameter are drastic and measurable between protons and heavy ions, they are subtle among heavy ions. Thus, it is difficult, if not impossible, to discern between different heavy-ion species using this approach. In addition, the thin end of the wedge readily stops energetic heavy ions, thereby significantly increasing the low-energy measurement limit for these measurements.

### 3.2.2 Magnetic Charged-Particle Spectrometers

The magnetic Charged-Particle Spectrometers (CPS1 and CPS2), implemented on OMEGA, utilize a high-field magnet and CR-39 for momentum-dispersion and detection of charged-particles, respectively (see Fig. 3.3). An interchangeable vertical slit (slit width 0.1 - 2.0 mm) is used to aperture down the incident flux of particles to accommodate a wide range of charged-particle yields.
(10^8 - 10^{15}). Two of these spectrometers, CPS1 and CPS2, are permanently mounted on fixed diagnostic ports on the OMEGA target chamber. These spectrometers provide nearly orthogonal views of the target (101°), thereby allowing an assessment of yield anisotropy.

These magnetic spectrometers are routinely used on OMEGA for diagnosing low-energy fast protons and deuterons, primary charged fusion products and scattered fuel ions. The CPS complement the WRFs in their ability to measure spectra down to 100 keV (proton equivalent energy) with high resolution. Selected ions are deflected through the magnetic field, sometimes through large-angles, according to their gyro-radii,

$$r_g \propto (A E)^{1/2}/Z$$

(3.1) expressed here in terms of the atomic mass A, charge number Z and energy of the particle, E. Note that the position of a given particle along the dispersion plane is inversely proportional to the particle’s gyro radius. Nevertheless, it is clear from Eq. 3.1 that there exists a degeneracy among different ions species. For example, a 1 MeV deuteron will have the same gyro radius as a 2 MeV proton, and hence land at the same location along the dispersion plane. For low-Z ions (Z ≤ 2), though, discrimination between species is possible using track contrast and diameter.

The CPS cannot be used for quantitative spectral measurements of heavy ions (Z > 2). As in the case of the WRFs, discrimination of heavy-ion species is difficult because the variation of track contrast and diameter are weak among heavy ions with different charge states.

Spectral measurements of low-Z ions in the presence of heavy ions, however, is possible with the CPS. This can be difficult if the (undesired) heavy-ion yields are high, since the CR-39 will saturate much quicker because of the additional particles captured by the detector. To alleviate this problem, the CPS are typically run with filters constructed of either mylar, aluminum or tantalum. These filters, placed directly on the CR-39, are sized to the appropriate thickness required to
stop the undesired particles (e.g., carbon) while simultaneously allowing lighter ions (protons and deuterons) to pass through.

3.3 Development of Thomson parabolas for Heavy-Ion Measurements at Omega and MTW

3.3.1 The Thomson parabola Principle

A combined electric-magnetic spectrometer has advantages over magnetic spectrometers since it can break the charge-to-mass (q/m) degeneracy associated with magnetic dispersion (see Eq. 3.1). These two dispersion mechanisms allow simultaneous discrimination of ions based on momentum and energy. In the context of this work, Thomson parabolas allow simultaneous spectral measurements of a variety of heavy ions with different charge-states. Such a spectrometer was first used by J. J. Thomson in 1913 for the discovery of electrons. The spectrometer, which now bears his name, utilizes a parallel (or anti-parallel) arrangement of static electric and magnetic fields.

The dispersion of ions using this field arrangement results in parabolas on the detector plane, as shown in Fig. 3.4. Parabolas are discrete contours of constant q/m ratios. Each position along a fixed parabola corresponds to a unique incident particle velocity (or energy, if the materials used in an experiment are known). The Thomson parabola thus allows simultaneous spectral measurements of a number of heavy-ion species. Two Thomson parabolas were built, one each for Omega and MTW. In this section, we present an overview of Thomson parabolas in general, followed by design details of each spectrometer in Secs. 3.3.2 and 3.3.3.

In general, the field arrangement can be such that the magnetic field precedes the electric field or vice-versa. In some cases, designs have utilized overlapping field regions, resulting in a compact design. In terms of decoupling the two dispersion mechanisms from one another, it is best for the magnetic field to precede the electric field. The alternative field arrangement results in a situation where the electric field imparts additional energy and hence velocity to the particle while it is being deflected by the magnetic field. This causes additional $v \times B$ coupling where the velocity is no longer the initial velocity of the particle. Although this effect is not detrimental to the operation of the Thomson parabola, it does make calibration of the instrument more difficult.

In such an arrangement, with the magnetic field preceding the electric field, the magnetic and electric deflection scale q, m, energy of the particle ($E_p$), magnetic field strength (B), and the electric field strength (E) as

$$\Delta M \propto q B / (E_p m)^{1/2}$$
$$\Delta E \propto q E / E_p.$$  \hspace{1cm} (3.2a) \hspace{1cm} (3.2b)

The constants of proportionality depend on the geometry of the Thomson parabola (i.e., geometry of the field-generating magnet and electrodes). Derivations of the exact formulations for a simple, uniform field geometry are provided in Appendix D. Typical values for the magnetic and electric deflection are a few centimeters. The scalings presented in Eqs. 3.2a-3.2b can be combined with the instrument-specific parameters in table 3.1 and certain operating conditions to estimate the deflection for a given ion-specie. Fringe fields will modify the deflections calculated using the above formulas. Combining Eqs. 3.2a-3.2b (and eliminating the particle energy in the process) gives an
equation for the parabola slope

\[
\Delta E = \alpha \Delta^2 M
\]

\[
\alpha \propto \frac{m E}{q B^2}
\]

The slope is a function of the field strengths and geometry of the Thomson parabola, and the q/m ratio of the particle. Thus, the Thomson parabola also has a degeneracy in q/m space, although this is less prevalent in practice. Each position along that parabola corresponds to a unique velocity, which can be converted to an energy value. Note that once the species has been identified by its slope, the magnetic (or electric) deflection by itself is sufficient to determine the energy of the particle. Since the electric field is a tunable parameter that can be varied from one experiment to the next, it is convenient to use the magnetic deflection for energy calibration.

The resolution of a Thomson parabola is determined by a combination of the amount of particle deflection due to the fields and the optics of the instrument. Typically, the instrument is run with a collimator or aperture. The projection of this aperture onto the detector plane is determined by the aperture size \(a\), the length of the instrument \(d\) and the distance from the source of charged-particles to the aperture \(D\). For the Thomson parabolas used in this work, the fractional energy resolution for any given particle scales as:

\[
\frac{\Delta E_p}{E_p} \propto \frac{a(m E_p)^{1/2}}{qB} \left( 1 + \frac{d}{D} \right).
\]

These scalings can be used with the instrument-specific design parameters, found in table 3.1 and a certain set of operating conditions to estimate the resolution for a given particle species. For the Thomson parabolas used in this work, typical values of the fractional energy resolution are just a few percent. More information on the applicable limits of Eq. 3.4, including a derivation, can be found in Appendix D.

For a given ion species and charge-state, the upper-bound on the measurable ion energy is determined by the acceptable energy resolution for a given measurement. It is evident from Eq. 3.4 that as the energy of a particle increases, the fractional energy resolution gets worse \(\Delta E_p/E_p\) increases). Thus, the energy corresponding to when the fractional resolution becomes unacceptably large determines the high-energy limit. For the measurements presented in this work, an energy resolution of 5% - 15% is acceptable.

The lower bound on the measurable ion energy is determined by the maximum deflection set by the fields and the size of the detector. If an ion is deflected outside the detector region because the magnetic field is too large, then the lower bound is determined by Eq. 3.2a, where \(\Delta M\) takes on the value of the detector size in the magnetic dispersion direction. Similarly, if an ion is deflected
outside the detector region because the electric field is too strong, Eq. 3.2b applies where \( \Delta E \) takes on the value of the detector size in the electric-dispersion direction. Instrument-specific values for these limits are given in the following sections.

These are the basic elements behind the operation of Thomson parabola Spectrometers. Additional details, including filtering schemes, dynamic range and detecting medium are instrument specific. In the next two sub-sections, specific implementations of two Thomson parabolas, currently operating at the MTW and Omega laser facilities are presented.

### 3.3.2 Thomson Parabola Ion Spectrometer (TPIS)

The TPIS\(^7\) was built to resolve heavy ions at the MTW Facility. Since the facility consists of a single low-energy short-pulse laser, the spectrometer was designed to measure moderate fluence of multi-MeV protons and heavy ions. As shown in Fig. 3.5, the main components comprising the TPIS are a circular aperture, a permanent magnet, electrostatic deflector plates, and a stacked detector assembly consisting of a CR-39 detector and an image plate (IP). An interchangeable tantalum aperture, with a 400 \( \mu \)m diameter is used to control the field of view and to partly set the energy resolution of the diagnostic. The permanent dipole magnet features a peak field of 5.3 kG, a pole gap of 1 cm and 5.1 cm \( \times \) 5.1 cm square pole surfaces, resulting in a uniform magnetic field over a large region. The electrostatic deflector plates are 2 cm apart, and are constructed from stainless steel and affixed to the diagnostic housing using a dielectric; with proper conditioning, they can be operated to potentials as high as 80 kV.

The detector pocket can accommodate either a 10 cm \( \times \) 10 cm, 1.5 mm-thick piece of CR-39 and a Fuji TR IP, or a stacked detector assembly consisting of both. The Fuji IPs are capable of detecting protons with energies above 20 MeV, though the sensitivity is a strong function of energy. An absolute calibration of the IP response to protons has therefore been recently conducted at the MTW for proton energies in the range of 1-8 MeV.\(^7\) Since the Fuji TR lacks a protective layer of Mylar found on other IPs, carbon ions can be detected as well. Heavier ions are readily stopped before they can deposit energy into the active layer of the IP, and are thus not detectable. IPs are therefore ideal for measurements of energetic protons whereas CR-39 is essential for the detection of high-Z ions.

The TPIS was implemented to study protons and heavy ions accelerated from LPI with planar targets. It was designed to measure proton equivalent energies in the range of 0.8 MeV to 20 MeV.
The low-energy limit is set by the size of the detector and magnitude of the fields, and the upper limit of 20 MeV (protons, with a resolution of 16%) is set by the maximum proton energy expected to be generated by the MTW. The TPIS is permanently fixed to a diagnostic port on the MTW target chamber. The distance from the aperture to the target chamber center, where targets are typically positioned, is 50 cm.

The first consideration for heavy-ion measurements using CR-39 in the TPIS is the acceptable range of ion fluence on the CR-39. This range is set by the intrinsic noise floor at the low end and saturation (track overlap) of the detector at the high end. The on-detector fluence requirements are a minimum of about $10^4$ proton tracks per cm$^2$ for good statistics relative to the background intrinsic noise, and no more than proton $10^6$ tracks per cm$^2$ to avoid track overlap. After accounting for the standard 400 μm aperture and 50 cm target-to-aperture distance, this requirement translates to an acceptable incident fluence range of 10 - $10^3$ ions/μsr for a an on-detector area that is defined by the projection of the aperture onto the detector. Since the latter defines the energy-resolution, this incident fluence range is essentially for a mono energetic beam of particles. Thus, for an arbitrary (normalized) spectral shape given by $S(E_p)$ spanning between energies $E_{p1}$ and $E_{p2}$, the maximum total measurable incident fluence ($F$, in units of ions/μsr) is given by:

$$F = 10^3 \int_{E_{p1}}^{E_{p2}} \frac{S(E_p)}{\Delta E_p(E_p)} dE_p,$$

where $\Delta E_p$ is given by Eq. 3.4. For LPI-generated fast ions, the spectra are typically continuous with an exponential dependence on energy. Thus, large incident ion fluence (relative to that of charged-fusion products) are readily possible.

The second consideration for heavy-ion measurements with the TPIS is the instrument response to these ions. It is not clear that the scalings for the deflection of ions in the TPIS due to uniform magnetic and electric fields, as given by Eqs. 3.2a-3.2b, hold for a range of heavy ions. For accurate identification of the particle species and precise determination of ion energies, it is important to verify these scalings for a range of heavy ions.

The initial TPIS calibration was obtained through experiments at SUNY Geneseo. These experiments were performed with 0.6-3.4 MeV protons to obtain the magnet energy calibration, while additional experiments using 0.8-3.4 MeV protons were conducted to calibrate the line-integrated electric field strength. Although these calibrations are in excellent agreement with the nominal scalings, it is not clear that they may be extrapolated and applied to more energetic ions with a large mass-to-charge ratio, such as C, Cu or Al. These heavier ions, which traverse a different part of the magnet and electrodes, may experience non-uniformities and fringe fields that lighter ions do not. For densely packed parabolas (e.g., lower operating voltages or higher-Z targets) a calibration accurate to within just a few percent is required to correctly identify the particle species and charge state. A charged-particle particle-tracking code was developed at MIT and used to compute the trajectories through the field maps for a range of ions, charge states and energies. The magnetic field map was supplied by Dexter Magnetic Technologies Inc., and the electric field map was generated by a finite-element method (FEM). The FEM approach used detailed models.

Figure 3.6. (a) Magnetic field map supplied by Dexter Magnetic Technologies, Inc., Inc. (b) Electric potential map, from which the electric field is calculated by using a finite element analysis and CAD models for the TPIS.
of the electrodes including the Ultem dielectric ($\kappa \sim 3$) to which they are affixed, as well as the instrument housing. The potentials on these surfaces were solved for all locations using the appropriate boundary conditions. The electric fields were computed using a finite-difference of the three-dimensional potential map using a fine grid. The simulated results for the magnet energy scaling of protons and carbon ions are shown in Fig. 3.7a. The magnetic field map had to be scaled in strength by approximately 3% to match the experimentally determined energy scaling. This systematic difference may be attributed to the position accuracy of the field map (since there is likely to be a systematic discrepancy between the simulated and constructed magnet assembly), as well as small offsets in the alignment of the magnet relative to the housing of the as-built TPIS. The absolute energy calibration is accurate to within 1%, as inferred from small deviations between simulations and experiments that persisted after scaling the magnetic field map. The magnetic energy calibration has been combined with simulations of the electrostatic deflection of protons in the TPIS to obtain a derived calibration for $\alpha$ as a function of electrode voltage. This is illustrated in Fig. 3.7b alongside the SUNY Geneseo calibrations.\(^7\)

The calibrated TPIS was used to study the acceleration of protons and heavy ions from the rear of thin foils at the MTW laser facility. These included several charge states of carbon, nitrogen, oxygen, aluminum, and copper, as discussed in Chapter 6. Table 3.1 summarizes the TPIS design parameters.

### 3.3.3 Thomson Parabola Ion Energy Analyzer (TPIE)

The TPIE Analyzer was originally designed and built by Los Alamos National Laboratory (LANL) for use on the OMEGA EP laser. As in the case of the TPIS, it was also used to used by LANL to study fast ions generated from short-pulse LPI with planar targets.\(^9\) In this work, the TPIE Analyzer was ported from the OMEGA EP target chamber to OMEGA for measurements of fast ions generated from long-pulse LPI in spherical implosions.

The OMEGA and OMEGA EP target chambers have similar diagnostic handling capabilities. In particular, the TIMs used on both systems are compatible, so that diagnostics used on one laser can be migrated to the other. The TPIE Analyzer, as shown in Fig. 3.8, is a TIM-based diagnostic. It features an interchangeable square tungsten aperture (100 $\mu$m, 250 $\mu$m and 1 mm), an interchangeable dipole magnet (with either 5.6 kG or 8.4 kG peak field strengths), parallel-plate electrodes, and a detector “goalpost.”
Since TPIE is fielded within a TIM, the distance from the aperture to TCC can be varied continuously between 37 cm and 150 cm. The near-TCC limit is set by the TIM space-envelope interference with incident laser beams. The far-limit is set by the range of motion of the TIM itself. In contrast to the TPIS, this flexibility, combined with the interchangeable aperture, allows the experimenter to simultaneously set the energy resolution and to optimize the system for a given particle yield.

The dipole magnets used on TPIE have a large pole region \((10 \text{ cm} \times 5 \text{ cm})\) and result in more dispersion for the same peak field strength relative to the TPIS. The interchangeable magnet assembly affords TPIE tremendous flexibility, since it is capable of probing protons in the energy range of 1.5 MeV - 100 MeV using the 5.6 kG and 8.4 kG magnets, respectively, with a fractional energy resolution of just a few percent. Lower-field magnets can be fabricated to probe less energetic particles.

The electrodes are long (21 cm) with a small separation gap (1 cm) that results in more electric dispersion for the same voltages relative to the TPIS. These electrodes are fastened to the rest of the diagnostic assembly using KEL-F dielectrics. Two compact DC-DC converters with adjustable voltage outputs, shown in Fig. 3.8, are used to convert the 24 VDC power supply from the TIM to the high-voltages required for operation. The voltage on each electrode can be independently set, although the opposing polarity on the voltage converters requires that the potentials be negative on one electrode and positive on the other with respect to ground. The maximum operating voltage for the electrodes is 20 kV (total), but the practical upper limit is set by low-energy ions striking the electrodes at a high enough voltage. For example, a 200 keV proton will strike the negative electrode for a total operating voltage of 2 kV. Equation 3.2b can be used to scale this value for other species or voltages.

The detector goalpost (see Fig. 3.8), accommodates a 10 cm \(\times\) 5 cm piece of CR-39 (1.0 mm thick), an IP, or a stacked assembly of the two. As in the case of the TPIS, the IP is required for
3.3 Development of Thomson parabolas for Heavy-Ion Measurements at Omega and MTW

Figure 3.9. Two null regions on a CR-39 detector from TPIE for a spherical implosion. The unfiltered region (left) shows the strong TPIE background from scattered low-energy ions, that is mitigated with the use of a 1.5 μm-thick layer of Mylar (right).

detection of high-energy (> 10 MeV) protons, while the CR-39 is required for detection of heavy ions.

In the process of porting TPIE from the OMEGA EP target chamber to OMEGA, it was discovered that a large fluence of low-energy scattered ions contaminate the CR-39 and hinder proper operation of TPIE. In fact, no measurable signals can be detected in the presence of this background. Fig. 3.9a shows a microscope image (40× magnification) of a null region on the CR-39 (after several minutes of etching in NaOH) from an OMEGA shot. The region is covered with particle tracks consistent with that of low-energy protons. To remove this background, a 1.5 μm-thick layer of Mylar was placed over the CR-39 (Fig. 3.9a) to stop the scattered ions. Thinner layers of mylar, including 0.7 μm and 0.9 μm, have also been tested and proven to be sufficient. These filters are thick enough to stop the large flux of ~ 100-200 keV protons generated from long-pulse LPI interactions with the capsule. Although one of the physics goals for TPIE is the measurement of such ions with energies up to the ~ 1 MeV range, the greater number of low-energy ions that are more prone to striking the electrodes and scattering off support structures within TPIE cannot be measured and thus set the low-energy cutoff of the spectrometer. Note that there are far greater numbers of low-energy ions since the energy-spectra of LPI generated protons is typically exponential in nature. The reason for the greater number of ions on OMEGA is likely due to the large difference in laser energy between OMEGA and OMEGA EP. OMEGA is likely to generate far more hot electrons through LPI and thus accelerate a greater number of ions. The goalpost was modified to accommodate a wide range of filters on OMEGA. A filter frame, containing the desired filter material and thickness is slid into the goalpost in front of the CR-39 / IP assembly. If necessary, filters as thick as 50 μm can be fielded.

The yield dynamic range for TPIE is set by the same limits as TPIS. The yield must be higher than the intrinsic noise floor but low enough to avoid track saturation (track overlap) of the CR-39. For a 100 cm target-to-aperture distance and a 250 μm square aperture, the acceptable incident fluence range for a monoenergetic beam of ions (with energies between \(E_p\) and \(E_p + \Delta E_p\), where \(\Delta E_p\) is given by Eq. 3.4) is in the range of \(10^2 - 10^4\) ions/μsr. Since TPIE can be used to measure charged-particles from implosions where the emission is typically isotropic (or nearly isotropic), this translates to total charged-particle yields of \(10^9 - 10^{11}\) over 4π for monoenergetic beams. As mentioned previously, these numbers do not take into account the spectral shape and are hence conservative estimates for the total measurable yields. Significantly larger yields (\(10^{15} - 10^{16}\)) are allowable for continuous charged-particle spectra (see Eq. 3.5).

The TPIE Analyzer was primarily calibrated using detailed simulations. A particle-tracking code developed by LANL, similar to the MIT code used for the TPIS, was used to compute particle trajectories through the TPIE fields. The magnetic field map was obtained from Dexter Magnetic Technologies Inc., while the electric field map was generated at MIT using FEM analysis of the electrodes. As shown in Fig. 3.10, it is necessary to incorporate the detailed geometry of the entire spectrometer to accurately simulate the potential everywhere. The proximity of the electrodes to grounding planes (e.g., the magnet, spectrometer housing) perturbs the potential and must be taken
Figure 3.10. Electric potential map for the TPIE Analyzer, generated using a finite element analysis and CAD models of system geometry. On the upstream (magnet) side of the electrodes, the magnet assembly acts as a grounding plane and forces the potential to drop much faster than on the downstream (goalpost) side of the electrodes.

into account. The energy calibration was checked experimentally using a stacked detector assembly with CR-39 on top of an IP. For this assembly, the CR-39 acts as a range-filter for protons, and hence only protons with enough energy to penetrate the 1.0-mm CR-39 (approx. 10.2 MeV) will leave a signal on the IP. This proton energy and its corresponding deflection due to magnetic fields was checked against Eq. 3.2a and found to be in agreement. The results of simulations for a range of ions are shown in Fig. 3.11. These simulations were run for the 5.6 kG magnet with a peak electric field of 12 kV-cm⁻¹. The experience gained from comparing simulations and experiments on the TPIS suggests (e.g., scaling magnetic fields maps to match experimental calibration campaigns) that at present, the absolute energy calibration of TPIE is accurate to within ~ 10%.

3.3.4 Analysis Techniques

A software program, known as the Thomson parabola Analysis Code (TPAC) was developed at MIT (see Appendix E) for reduction and analysis of data acquired using TPIS and TPIE throughout this work. The software package contains calibration information for both diagnostics and is equipped with analysis routines for identification of particles species and extraction of ion energy spectra. Detailed analysis procedures are discussed in Appendix E. Some of the salient features are summarized here.

As discussed in Sec. 3.2, processing of the CR-39 data includes etching in an alkali solution (typically diluted NaOH). After this process, latent tracks become visible and can be optically detected. An automated microscope system with 40× magnification is used to scan the entire CR-

Figure 3.11. Simulated parabolas on TPIE for protons, helium (³He and ⁴He), deuterium, and tritium. For these simulations, the peak magnetic field was 5.6 kG and the peak electric field was 12 kV-cm⁻¹.
Development of Thomson parabolas for Heavy-Ion Measurements at Omega and MTW

3.3

39 area, and to record location, diameter, eccentricity and contrast of each charged-particle track. This information is stored within a data file that is then analyzed by the TPAC.

The TPAC can be used to generate 2D histograms or “images” of parabolas. Charged-particle tracks are binned in two dimensions, and the value of each bin is mapped to a color to produce the figure shown in Fig. 3.12a. To ease identification of particle species, routines were developed to generate a histogram of particle tracks binned by parabola slope (see Fig. 3.12b). Furthermore, local detector noise can be correlated between Figs. 3.12a and 3.12b and excluded from the analysis. Fig. 3.12b can be used to determine the mean slope and establish the boundaries of each parabola. The mean slope is used to identify the ion species, while the boundaries are used to select tracks that pertain to an ion species of interest. These tracks are then binned according to their magnetic deflection. The bin size is determined by the projection of the aperture onto the detector. The magnetic deflection associated with each bin is then converted into a particle energy by using the magnet energy scaling. Each bin thus has an associated fixed mean energy, yet spans a small finite width in energy-space (the energy resolution). The result is an energy histogram, as shown in Fig. 3.12c for several ions.

At any point during the above analysis steps, contrast, diameter and eccentricity limits may be imposed or modified to filter undesired tracks or localized noise. In addition, sophisticated background subtraction algorithms were developed to generate both the N(α) and the Y(E) histograms. For the case of the latter, two subroutines are available and may be chosen based on the nature of the data set. The first subroutine performs a flat background subtraction, where a specified

Figure 3.12. (a) Two-dimensional histograms of the CR-39 detector from the TPIE Analyzer. Each pixel represents a number of heavy-ion particle tracks that were individually resolved using a microscope. The local track density was then mapped to a color value to produce these “images.” As shown here, several charge-states of aluminum from an flash-coated plastic (CD) shell (OMEGA Shot 65271) were measured (b) Corresponding histogram of particle tracks as a function of their parabola slope, α. The observed smaller line widths for decreasing α (increasing q/m ratio of the ions) are a result of magnetic focusing, which arises from spatial non-uniformities in the field strength. (c) Corresponding energy spectra for a number of aluminum charge states. For these data, the peak magnetic field was 5.6 kG and the peak electric field was 2.5 kV-cm⁻¹.
background region is sampled and used as a measure of the background for the entire detector. The second subroutine is adaptive, and for each energy bin along a parabola, looks just above and below to determine the local background level. The background is interpolated through the parabola in the direction of electric displacement. Although this adaptive algorithm is generally preferable, it fails when parabolas are not widely separated. In this case, the sampled background is large compared to the real background and particle tracks are under-counted. For such cases, the flat subtraction algorithm is used instead.

In addition to data analysis, the TPAC facilitates fielding of either the TPIS or TPIE by providing an interface where users can mock-up a setup, including voltage, and determine the range of measurable energies and the energy resolution for a given ion species.

### 3.4 Summary

Two Thomson parabola spectrometers have been developed in collaboration with LLE and LANL to measure fast heavy ions from ICF experiments. These spectrometers complement existing CPSs and WRFs that have been used for a number of years to measure light ions, including protons, deuterons, tritons and α-particles. The design specifications and operating regimes of the TPIS (implemented for use at the MTW Laser Facility) and the TPIE Analyzer (implemented for the Omega Laser Facility) are summarized in table 3.1.
3.5 References

Table 3.1. Specifications of the TPIE and TPIS Thomson parabolas on OMEGA, OMEGA EP and on the MTW, respectively.

<table>
<thead>
<tr>
<th></th>
<th>Thomson parabola Ion Spectrometer (TPIS)</th>
<th>Thomson parabola Ion Energy (TPIE) Analyzer</th>
</tr>
</thead>
<tbody>
<tr>
<td>Laser system</td>
<td>Multi-Terawatt (MTW) Laser</td>
<td>OMEGA, OMEGA EP</td>
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<tr>
<td>Detector</td>
<td>CR-39, Image Plate</td>
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<td>Detector size ($\Delta M \times \Delta E$)</td>
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<td>10 cm $\times$ 5 cm</td>
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<td>Energy range ($H^+$)</td>
<td>0.8-20 MeV</td>
<td>2-100 MeV*</td>
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<tr>
<td>Energy resolution ($H^+$)</td>
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<td>1.5% at 2 MeV†</td>
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<td>Aperture</td>
<td>400 $\mu$m (round)</td>
<td>100, 250, 1,000 $\mu$m (square)</td>
</tr>
<tr>
<td>Magnetic field (peak)</td>
<td>5.3 kG</td>
<td>5.6 kG, 8.4 kG</td>
</tr>
<tr>
<td>Magnet pole gap</td>
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<td>0.5 cm</td>
</tr>
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<td>Magnetic pole size</td>
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<td>Magnetic drift length</td>
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<td>45 cm</td>
</tr>
</tbody>
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* Applies for 5.4 kG magnet
† Applies for 5.6 kG magnet, 250 $\mu$m aperture and 100 cm source to aperture distance

3.5 References


4

Target Charging and Discharging in Direct-Drive Implosions on OMEGA

4.1 Introduction

The charging of ICF targets to several hundred kilovolts and the associated acceleration of fast ions due to this potential was first observed in electrically isolated planar targets nearly 35 years ago. In these early experiments, 0.3-cm-thick aluminum targets were irradiated with 50 ps laser pulses ($\lambda_L = 1.06 \mu m$). Electric potentials with peak amplitudes of order 1 kV were measured using a fast-response voltage probe. These potentials had relatively fast rise times (1 ns) and long decay times (10 ns). Charging of the target was attributed to the production of LPI-generated hot electrons. Electrons in the high-energy tail of the hot-electron distribution escape the target, leaving behind a positive potential. The remaining space charge accelerates ions from the coronal plasma surrounding the target and traps hot electrons that have energies less than the target potential.

After initial experiments in the 70’s, optical diagnostics were used to qualitatively characterize the currents flowing along the thin fiber or stalk that connects targets to the chamber. It was determined that ionization and subsequent heating of the stalk occur primarily because of the large electric fields across the stalk that ionize the stalk material into a plasma and drive an ohmic current through it.

Additional evidence for target charging was found in the 80’s on the six-beam ZETA facility and later on the 60-beam OMEGA Laser. On OMEGA, charging of plastic and glass shells to voltages of order 1 MV was observed (by Hicks et al.). In these experiments, the peak target voltage was inferred from time-integrated spectral measurements of fast protons. These fast protons, with total yields of $10^{13}-10^{14}$, are sourced from water vapor and other hydrocarbon contaminants typically present on target surfaces. Another indicator of target charging in these experiments was the energy upshift of charged fusion products (total yields of $10^{10}-10^{11}$). Charged fusion products gain additional energy as they escape and traverse the potential drop that surrounds the target.

Charging of ICF targets is of concern within the community because it can hinder the achievement of the $\rho R$ required for ignition and energy gain. The bulk of hot electrons are trapped within the target for a duration that is prescribed by the target potential decay time. During this time, these electrons deposit their energy within the fuel (ice layer of cryogenic targets). This electron preheat occurs during the acceleration phase of ICF implosions (the flat-top portion of the laser pulse) and raises the fuel adiabat, thereby reducing the achievable $\rho R$. For cryogenic OMEGA implosions, it takes about 30 J of preheat (0.1% of the laser energy) to reduce the $\rho R$ by 50%. At present, the measured $\rho R$ for cryogenic implosions falls short of the simulated $\rho R$ by about 15%. It is unclear whether the observed $\rho R$ reduction is due entirely to preheat or other effects as well,
as a robust predictive capability is lacking.

The amount of preheat and hence $\rho R$ degradation of a target can be addressed in detail using simulations of hot-electron stopping and scattering in the fuel. This requires knowledge of the hot-electron source characteristics (e.g. temperature, distribution) and target potential behavior (e.g. peak potential, time-dependence). The former has been addressed through experiments and simulations, whereas the temporal evolution of the target potential has not been previously modeled and is a major component of this chapter. The model developed here can be used to predict voltage decay times, which is useful for preheat calculations within hydrodynamic codes used to model ICF implosions. The hydrodynamic code LILAC,\(^8\) for instance, assumes at present that the target potential is large for all time and that all hot electrons are trapped within the target.\(^9\)

The temporal evolution of the target potential is determined by charge sinks, which include currents through the target stalk and acceleration of fast-ions from the corona. Since the hot electrons lose energy to these fast ions (and other sinks), measurements of the fast ions are essential for validating preheat calculations. First, fast-ion measurements place an upper bound on preheat, which is given approximately given by the difference between the total hot-electron energy and the energy lost to these ions. Second, the total charge of these ions is equal to the number of escaping electrons. A measurement of this charge thus places a constraint on the number of electrons that can escape the target (or equivalently, the number of electrons that are completely stopped in the shell). Fast-ion measurements can be used to support and benchmark the aforementioned preheat simulation.

This chapter is organized as follows: section 4.2 discusses hot-electron characteristics and presents fast-ion measurements as they relate to preheat and $\rho R$ reduction in OMEGA implosions. These measurements are used to estimate the amount of preheat in OMEGA implosions. Section 4.3 presents an empirical circuit model for the temporal evolution of the target potential. This model incorporates spectral measurements of fast ions and radiographic measurements of stalk currents and can be used in hydrodynamic simulations to more accurately predict the amount of preheat in OMEGA implosions. Sec. 4.4 concludes with a summary of these results.

## 4.2 Target Charging and Preheat in Direct-Drive Implosions

### 4.2.1 Hot-Electron Production in OMEGA Implosions

In direct-drive OMEGA implosions, the hot electrons responsible for target charging and preheat are produced by the TPD instability (see Sec 1.4.1). These hot electrons are produced for on-target laser intensities greater than \(4.6 \times 10^{14} \text{ W-cm}^{-2}\). The temperature and total energy of these electrons has been well characterized for 60 beam OMEGA-scale implosions, as summarized below.

The hot-electron temperature has been previously inferred using measurements of the hard x-ray (HXR) spectrum generated by the hot electrons as they stop on the dense shell.\(^10\) Time-resolved measurements of the hard x-ray spectrum are routinely conducted on OMEGA using a four-channel hard x-ray detector (HXRD).\(^11\) Each channel consists of a high-pass filter (composed of either aluminum or copper) and a fast scintillator that is coupled to a photomultiplier tube (PMT). A typical filter setup results in low-energy cutoffs at 20, 40, 60 and 80 keV. These channels are calibrated for the filter transmission and scintillator efficiency, resulting in a four-channel time-resolved measurement capability. Sample HXR signals are shown in Fig. 4.1 for warm and cryogenic implosions with different pulse shapes (from a single HXRD channel for photon energies > 40 keV). The time-averaged hot-electron temperature inferred using data from HXRD channels is shown in Fig. 4.3a as a function of the on-target laser intensity.
These HXR signals indicate when the electrons stop on the shell, rather than when they are produced. Onset of hot-electron production occurs slightly before hard x-ray production and has been previously investigated using measurements of the scattered light in these implosions. In particular, OMEGA laser beams can act as Thomson Scattering probe beams that interact with the electron plasma waves produced by the TPD instability. The scattering geometry is complex since many different OMEGA beams contribute to the scattered light spectrum. The resulting spectrum is thus an average over many Thomson-scattering configurations, but is characterized by a strong Thomson up-scattered component at $3\omega/2$. A sample time-resolved measurement of the $3\omega/2$ is shown in Fig. 4.2 alongside the laser pulse and hard x-ray signal. The $3\omega/2$ signal represents the onset of hot-electron production, and was previously found to increase with increasing laser intensity.\(^\text{12}\) At an intensity of $4 \times 10^{14}$ W-cm$^{-2}$, onset occurs roughly $\sim$ 400 ps into the flat-top portion of the laser drive. Similarly, at an intensity of $1 \times 10^{15}$ W-cm$^{-2}$, onset of hot-electron production occurs about $\sim$ 100 ps after the start of the main drive. In both cases, hot electrons are present for the remainder of the laser drive as evidenced by both the hard x-ray and $3\omega/2$ signals. These electrons are thus produced during the flat-top portion of the laser drive, when the outer part of the shell is being ablated away while the remaining shell mass is accelerated inward. At this time, the electrons can penetrate the shell and preheat the fuel, raise the adiabat and degrade target compression.

![Figure 4.1](image1.png)

**Figure 4.1.** Sample HXRD signals for implosions utilizing (a) a warm plastic shell irradiated with a square pulse (b) a warm plastic shell irradiated with a triple-picket low-adiabat pulse and (c) a cryogenic target irradiated with a triple-picket low-adiabat pulse. Shown alongside the HXRD signals are the laser pulse shapes for each shot (dashed lines) [HXRD signals courtesy of C. Stoeckl].

![Figure 4.2](image2.png)

**Figure 4.2.** Sample $3\omega/2$ signal from coherent Thomson scattering of OMEGA beams ($\omega$) with electron plasma waves ($\omega/2$), alongside the HXRD signal and the laser pulse for OMEGA shot 64684.
The total amount of laser energy coupled to hot electrons has also been previously investigated. For a 60 beam OMEGA implosion, the total amount of hot-electron energy was inferred by irradiating spherical molybdenum (Mo) targets with a CH overcoat. The overcoat characteristics are identical to the CH ablator of typical targets so that the absorption physics is accurately reproduced. The Mo sphere absorbs the generated hot-electrons and emits K-α line radiation (17.5 keV), which is related to the total hot-electron energy and can be readily measured using the HXRD. For a given hot-electron temperature and hence electron distribution function, the total amount of K-α energy was related to the total hot-electron energy using Monte-Carlo simulations of hot-electron-induced shell transitions in Mo. The total hot-electron energy found from these experiments is shown in Fig. 4.3 as a function of the incident on-target laser intensity.

For OMEGA implosions (∼ 20-30 kJ laser energy), no more than ∼ 200 J (1%) of the laser energy is coupled to hot-electrons for an on-target laser intensity of about 1 × 10^{15} W-cm^−2. This is well over the ∼ 30 J estimate required to degrade ρR by 50% and is not observed in experiments. The total hot-electron energy is lost to several sinks: (i) preheating the fuel (the ice layer in cryogenic implosions), (ii) heating the outer part of the unablated shell (the CD ablator in cryogenic implosions), (iii) hard x-rays from bremsstrahlung radiation and (iv) acceleration of fast ions from the corona. Losses from (iii) have been estimated using the HXRD, and were found to be negligible (∼ 100 mJ) for typical implosions. Measurements of the total energy lost to fast ions (iv) can thus be used to estimate the amount of hot-electron energy available for preheat and are presented in the next section.

### 4.2.2 Fast-Ion Acceleration

Previous spectral measurements of fast ions from OMEGA implosions were limited to protons and deuterons from warm plastic and glass shells. Results indicated that approximately ≤ 100 J (0.3%) of energy was carried by fast protons. These measurements were extended to include all fast ions and a greater variety of laser and target configurations.

The spherical targets used in this study had diameters ranging from 400-860 μm, were composed of either plastic (CD or CH) or glass (SiO₂) with wall thicknesses ranging from 2-27 μm. These targets were filled with either DT, D₂ or D³He gas. In some cases, the targets were flash-coated with 100 nm of aluminum to mitigate diffusion of ³He. Cryogenic targets, with a plastic (CD) shell and an inner layer of DT or D₂ ice (∼ 100 μm), were also studied. The on-target laser energy (∼ 20 kJ) and pulse duration (∼ 1-3 ns) were varied to achieve on-target intensities in the range of 10^{14}-10^{15} W-cm^−2. The laser pulse shapes used in this study were square or shaped and in some cases preceded by either a single or triple picket.
Fast proton and deuteron energy spectra were measured using the magnet-based charged-particle spectrometers (CPS1 and CPS2, see Sec. 3.2.2). A Thomson parabola (TPIE, see Sec. 3.3.3) was used for spectral measurements of heavier fast ions (C, Si, O, Al).

Data recorded by CR-39 detectors in TPIE are shown in Fig. 4.4. Glass and plastic (CD) targets show several strong lines of heavy ions, including silicon, oxygen and aluminum. In some cases several weaker lines not visible in these images, including charge states of carbon were detected and recorded. It was found that the majority of energy is carried by highly ionized charge states of silicon (for glass) and aluminum (for flash-coated CD targets). Sample energy spectra of these fast heavy ions are shown in Fig. 4.5 alongside proton spectra acquired using CPS1. The maximum proton and heavy-ion energies scale linearly with the ion charge state, as shown in Fig. 4.6 (the maximum ion energies are not always visually clear in Fig. 4.5 due to the truncated vertical axis). Heavy-ion spectra and scalings of the maximum ion energies for a number of other shots can be found in Appendix F.

For each shot, spectra similar to those shown in Fig. 4.5 were integrated to obtain the total energy carried by fast ions. Fig. 4.7 shows the percentage of laser energy carried by fast ions as a function of the peak voltage on the target, which was determined by measuring the energy of the fastest proton. These data incorporate protons from CH and glass targets, protons and deuterons from CD targets (solid symbols) and the total energy carried by all ion species for a few CD and glass targets (open symbols). The latter data show that the total energy carried by fast ions is comparable between CD and glass targets for a given potential. Neither laser pulse parameters nor target metrology affect the shape or the magnitude of the curves shown here, as each of them

![Figure 4.4](image)

**Figure 4.4.** Data recorded by a CR-39 detector in the Thomson parabola (TPIE) spectrometer is for a glass-shell (a) and an aluminum flash-coated plastic (CD) shell (b). Each pixel represents a number of heavy-ion particle tracks. The local track density was then mapped to a color value to produce these “images.” Several lines of highly ionized silicon and oxygen (from glass shells) and aluminum and oxygen (from flash-coated CD shells) were measured. TPIE was run with peak magnetic and electric fields of 5.4 kG and 2.8 kV-cm$^{-1}$, respectively.

![Figure 4.5](image)

**Figure 4.5.** Energy spectra of fast ions, including several charge-states of aluminum alongside protons from a flash-coated plastic (CD) target (OMEGA Shot 65273). The maximum proton energy for this shot (corresponding to the peak target voltage) was 1.35 MeV. Spectra for a number of other shots can be found in Appendix F.
Figure 4.6. Scaling of the maximum proton and heavy-ion energies with the ion charge state (OMEGA Shot 65273). The maximum ion energies scale linearly with the ion charge state (the maximum ion energies are not always visually clear in Fig. 4.5 due to the truncated vertical axis). Scalings for a number of other shots can be found in Appendix F.

... incorporate a variety of target parameters (diameters and thicknesses), and laser parameters (laser energies, pulse duration and pulse shapes). This demonstrates that the fraction of laser energy carried by fast ions is a unique function of the peak target potential.

Based on the measurement of the total energy carried by fast ablators in glass and CD targets for voltages between 1-1.5 MeV, we estimate that heavy fast-ions carry nearly as much energy as their proton and deuteron counterparts. The comprehensive dataset for CD is described by a power law dependence on the target voltage (reduced $\chi^2 = 0.99$) and then scaled to match the total energy loss measurement for CD shells, resulting in the dashed curve shown in Fig. 4.7. From the fit, the fraction of incident laser energy carried by fast ions is given by

$$f_i = 3.7 \times 10^{-3} V_{0,\text{MV}}^{2.55},$$

(4.1)

where $V_{0,\text{MV}}$ is the maximum proton energy (in megavolts) that corresponds to the peak target voltage for each shot. The upper error bars of the total ion energy measurement and hence the uncertainty in the above fit are determined by the accuracy of the measurement. The low-energy cutoff of TPIE (1.6 MeV proton equivalent) precludes measurement of the spectrum at lower energies, resulting in an underestimate of the total energy by no more than 40-50%. This effect varies from one ion-species to the next and depends on the spectral shape, but for most ions shown here, the majority of the spectrum is measured. The lower error-bars for the total ion energy measurement are a result of instrument precision (25%) that is determined by uncertainties in the energy calibration of the TP. The magnetic spectrometers CPS1 and CPS2 measure proton and deuteron spectra down to 100 keV, resulting in precision-dominated error bars ($\pm 25\%$) for the measurements of the energy carried by the light ions shown in Fig. 4.7.

Since the energy loss to fast ions is parameterized by the peak voltage, it is important to examine what parameters dictate the initial voltage on target. A previous study showed that the target voltage scales with the on-target intensity, and that it was larger for thin $< 5 \mu\text{m}$ CH shells in

Figure 4.7. Percentage of incident laser energy carried by fast ions, including protons from CH and SiO$_2$ shells ($\lesssim 0.3\%$), protons and deuterons from CD shells ($\lesssim 0.6\%$), and heavier fast ions CD and SiO$_2$ ($\lesssim 0.3\%$). The total incident laser energy coupled to fast ions (dashed curve) is estimated by scaling the CD dataset (solid blue circles) to match the total ion measurements (open blue circles). Where not shown, the error-bars are comparable to the symbol size.
comparison to thick > 5 μm shells. We extend that study to include data from cryogenic and warm CD-shells (see Fig. 4.8) with a variety of shell thicknesses (3-10 μm-thick) in addition to CH-shells (15-27 μm-thick) and thin glass shells (2-3 μm).

Thin glass shells follow the scaling of thicker shells, resulting in lower peak target voltages for a given intensity than thin plastic shells. The observed discrepancy between thin glass and thin plastic shells is due to differences in shell absorption of laser light and hot-electron stopping power. Higher collisional absorption of incident laser light for higher-Z thin glass shells relative to thin plastic shells leads to higher coronal temperatures ($T_c$) and lower laser intensities at the quarter-critical surface. These conditions can mitigate the TPD instability, which has a threshold parameter $\eta \propto I/T_c$. In addition, the path-integrated hot-electron stopping power ($\Delta E = \int dE/ds ds$) of thin glass shells is greater than that of thin plastic shells (for 50-100 keV electrons, $\Delta E_{SiO_2}/\Delta E_{CD, thin} \approx 2-4$). Thus, any hot-electrons that are produced are more readily stopped by thin glass shells, resulting in lower peak target voltages. For thin glass and thick plastic shells, the scaling with intensity is given by (reduced $\chi^2 = 0.78$):

$$V_{0,MV} = 0.12 I_{14} - 0.4,$$  \hspace{1cm} (4.2)

where $V_{0,MV}$ is the peak target voltage in megavolts and $I_{14}$ is the laser intensity in units of $10^{14}$ W-cm$^{-2}$. For thin plastic shells, the peak voltage is systematically higher (reduced $\chi^2 = 0.89$):

$$V_{0,MV} = 0.13 I_{14} - 0.18.$$  \hspace{1cm} (4.3)

Shell absorption arguments can also explain the difference between thin and thick plastic shells. It has been shown for cryogenic targets that thicker shells prevent the laser from burning through to the lower-Z D$_2$ or DT ice-layer. For thinner shells that burn through, it has been suggested that the lower-Z deuterium plasma penetrates into the underdense region, resulting in lower collisional absorption of laser light. This creates conditions that are favorable for the TPD instability and hence the generation of hot-electrons.

For these data, Distributed Phase Plates (DPP) and Distributed Polarization Rotators (DPR) were used to shape the spatial profile of the laser beam and to randomize the laser polarization, respectively. On some shots, laser speckles were smoothed using Smoothing by Spectral Dispersion (SSD). The observed scatter in the cryogenic data is attributed to the wide range of pulse shapes used in these experiments. Scatter in warm-plastic and glass targets is largely due to uncertainties in the measurement of the peak target potential. For these shells, pulse shapes (1 ns, square) and SSD modulations were comparable.

The observed linear scaling of the maximum fast proton energy with laser intensity is expected
Figure 4.9. Target voltage for D$_2$ cryogenic implosions as a function of the hard x-ray signal (shown here as the charge collected on the hard x-ray detector). The hard x-rays are from bremsstrahlung radiation (photon energies $> 40$ keV) produced by hot-electrons stopping on the shell. The target voltage scales with hard x-ray emission, which in turn has been previously shown to correlate with TPD-generated hot-electrons.

because these ions are accelerated by TPD-generated electrons. The temperature of the hot electrons scales linearly with laser intensity as demonstrated above (Sec. 4.2), and there is a linear relationship between the maximum fast-proton energy (the voltage) and the hot-electron temperature (see Sec. 1.4.3). Further confirmation that the fast ions shown here are accelerated by TPD-generated hot electrons is the observed correlation between the target voltage and the hard x-ray signal that arises from electrons stopping on the shell (Fig. 4.9). In addition, it is worth noting that Hicks et al. previously used the fast-proton spectrum to infer a hot-electron temperature$^5$ that is consistent with the HXRD temperature measurements discussed above (Sec. 4.3).

4.2.3 Estimates of Preheat

As discussed in Sec. 4.2, the fast-ion measurements can be used together with measurements of the total hot-electron energy to estimate the amount of preheat. Eq. 4.1 for the scaling of the total fast-ion energy as a function of peak voltage was combined with the voltage-intensity scaling for thick shells (Eq. 4.2) to obtain an empirical expression for the fraction of laser energy that is coupled to fast ions as a function of on-target laser intensity:

$$f_i = 3.7 \times 10^{-3} (0.12 I_{14} - 0.4)^{2.55}$$  \hspace{1cm} (4.4)

This equation is shown alongside the total hot-electron energy in Fig. 4.10a. The difference between these two curves corresponds to the total energy coupled to the entire target. The amount of preheat (energy coupled to the ice layer) is a fraction of the total energy coupled to the target, since some of the energy is deposited in the plastic ablator. This fraction ($f_{PH}$) was estimated using the relative stopping power of the DT ice and the CD ablator:

$$f_{PH} \equiv \frac{\left[\frac{dE}{d(\rho L)} \times \rho L\right]_{DT}}{\left[\frac{dE}{d(\rho L)} \times \rho L\right]_{DT} + \left[\frac{dE}{d(\rho L)} \times \rho L\right]_{CD}} \approx 0.8 - 0.9$$  \hspace{1cm} (4.5)

where $dE/d(\rho L)$ is the stopping power of the material.$^{20,21}$ $\rho$ is the density (2 g-cm$^{-3}$ for solid DT and 1 g-cm$^{-3}$ for CD plastic), $L$ is the thickness (variable for DT, 10 $\mu$m for CD). Although stopping power depends on energy, $f_{PH}$ is a weak function of energy for electrons with energies between 10 keV and 1 MeV. The range of numerical values in Eq. 4.5 correspond to DT ice layers with thicknesses of 30-90 $\mu$m. The amount of preheat is thus the product of $f_{PH}$ and the maximum preheat (total energy coupled to the target) shown in Fig. 4.10.

The $\rho R$ reduction corresponding to the preheat is estimated as follows. The scaling of maximum $\rho R$ with the fuel adiabat ($\rho R \sim 1/\alpha^{0.54}$, see Eq. 1.11) can be combined with the adiabatic equation of state ($p \sim \alpha \rho^\gamma$ with $\gamma = 5/3$, see Eq. 1.9) and the ideal gas law to yield the following expression.$^6$
where the subscripts denote the areal density and fuel temperature without preheat. Eq. 4.6 indicates that the $\rho R$ reduction scales nearly linearly with the temperature increase of the fuel. To calculate the $\rho R$ reduction then requires knowledge of the initial and final temperatures. The initial fuel temperature can be obtained from hydrodynamic simulations which indicate a cryogenic fuel temperature of 15-20 eV during the acceleration phase of the implosion when the hot electrons are first produced. This temperature is determined by the adiabat of the laser pulse. For higher adiabat pulses, stronger shocks are launched at the beginning of the implosion and result in higher fuel temperatures during the acceleration phase of the target. The final temperature can be calculated using the estimate preheat and the number of ions in the ice layer. While Eq. 4.6 applies generally to implosions on OMEGA (and the NIF), the detailed calculation presented here is dependent on target metrology (e.g., thickness and outer diameter of the ice layer) and the laser pulse shape (e.g., the adiabat and onset of hot-electron production). Figure 4.10b shows such a calculation for typical 860-μm-diameter OMEGA targets with different ice layer thicknesses (30-90 μm) that are irradiated with a low-adiabat pulse (triple-picket, 1.5 ns duration). This calculation assumes that the laser energy is varied to achieve the shown intensities, although to achieve intensities of $0.8 - 1 \times 10^{15} \text{ W-cm}^{-2}$ requires the use of slightly smaller targets or a shorter pulse duration, in which case the preheat is slightly less than indicated by the curves.

This calculation is an estimate, since the full hot-electron dynamics are not taken into account, but it illustrates two important points. First, thicker ice layers are able to tolerate more preheat since the temperature increase is lower for the same amount of preheat energy. Second, preheat is insignificant for intensities below $6 \times 10^{14} \text{ W-cm}^{-2}$. The latter is an important point and suggests that any $\rho R$ degradation observed at lower intensities is likely the result of another process, either a different source of preheat or something different altogether.
4.3 Temporal Evolution of the Target Potential

4.3.1 Model Overview

The capsule discharging model is developed by first considering the acceleration of fast ions from the coronal plasma towards the target chamber wall. The ions expand together with the hot electrons (temperature of $T_h$) with a characteristic velocity given by the hot-electron sound speed ($c_h \equiv \sqrt{T_h/m_i}$), reducing the target potential (the strength of the sheath fields surrounding the target) as the target shell is imploding. This hot-electron and fast-ion expansion front is the sheath boundary ($R_n \sim c_h t$) indicated in Fig. 4.11a. Outside this boundary, the fields are screened and the potential is zero. The structure of the sheath is discussed further in Sec. 4.3.2. The decay of the potential is thus captured by a lumped circuit element model consisting of a capacitor ($C_T$) discharging through a resistor ($R_I$, see Fig. 4.11). The capacitor represents the space charge across the sheath surrounding the target, while the parallel time-dependent resistor allows for a flow of direct current across the sheath in the form of fast ions.

The stalk is composed of segments of silicon carbide (70 $\mu$m diameter, 1 cm long) and boron (140 $\mu$m diameter, 1 cm long) and connects the target to the electrically grounded chamber. It allows the flow of current to and from the target and is modeled using a series resistor-inductor combination (Fig. 4.11). It has been shown recently by Manuel et al. (see Ref. 22) that currents flow along a thin outer layer of the stalk that is ionized early on in the implosion. This annulus has been observed to be several tens of microns thick, which is comparable to the skin depth of the stalk ($\delta \sim 50 \mu$m) for the fast voltage pulses considered here. This thin layer of the stalk explodes outward as it carries current, resulting in an inductance that decreases with time, whereas the core of the stalk remains intact even after the implosion and does not carry current.

Given this circuit model, the goal is to determine the decay time of the system, or equivalently, the voltage on the capacitor (the target potential) as a function of time. Although the model itself can be applied to a wide range of targets, the parameters are empirically derived in the following section and have limited scope. We restrict the discussion to either warm or cryogenic spherical targets consisting of $\gtrsim 10$-$\mu$m-thick plastic or glass shells irradiated on OMEGA with intensities of
4.3 Temporal Evolution of the Target Potential

4-10 \times 10^{14} \text{ W-cm}^{-2}.

We assume that the target voltage rise time is fast relative to the decay time ($\tau_r \ll \tau_d$). The voltage rise time is characterized by the time it takes for the fastest electrons of the hot-electron population ($\gtrsim 50 \text{ keV}$) to traverse and escape the target ($\sim 1 \text{ mm}$ in diameter), which is on the order of $\sim 10 \text{ ps}$. Exact calculation of $\tau_r$ is difficult because hot electrons are produced continuously and the voltage peaks when the source of hot electrons equals the charge sinks (fast-ion acceleration and stalk currents).

The decay time of the potential depends on the magnitude of currents in the stalk and how quickly fast ions are accelerated from the target to deplete the space charge. As shown in this paper and first measured on OMEGA by Hicks et al., the potential falls soon after the laser turns off resulting in a decay time of $\lesssim 1 \text{ ns}$. We thus focus on modeling the decay time, and the initial conditions of our model are $V_T|_{t=0} = V_0$ and $dV_T/dt|_{t=0} = 0$, where $t = 0$ corresponds to a few tens of ps after the onset of hot-electron production.

This model is only valid for positive target voltages. For negative target voltages, the model no longer portrays the underlying physical picture. Once the target voltage crosses zero, the space charge no longer exists and current cannot flow in the reverse direction (the hot electrons and fast ions expand adiabatically). In this sense the resistor branch of the target behaves like a diode. In the following analysis, we constrain the target voltage to $V_T > 0$.

4.3.2 Empirical Circuit Parameters

The values of the lumped circuit elements have been determined empirically, using charged-particle spectroscopy and monoenergetic charged-particle radiography. Some of these measured parameters will be swept within their experimental uncertainties to obtain the best agreement with the measured temporal evolution of the target potential, as discussed in Sec. 4.3.3. In each of the following subsections, we use the results of Sec. 4.2 and those obtained by Manuel et al. (Ref. 22) to extract parameters relevant to the circuit model.

Target Capacitance

The capacitance across the sheath surrounding the target can be inferred from Eq. 4.1. The dashed curve in Fig. 4.7 can be thought of as the energy stored in a capacitor ($E_T = 1/2 CV_0^2$) that is initially charged to a peak voltage ($V_0$). Since the power-law fit of Eq. 4.1 is normalized to laser energy and has an exponent $> 2$, we expect the capacitance to scale with laser energy and exhibit a power-law dependence on the voltage. Taking $C \sim \beta E_L V^\alpha$, where $\alpha$ and $\beta$ are constants to be determined from the data, the total energy stored in the capacitor is given by:

$$E_T = \int_0^{V_0} \beta E_L V^{1+\alpha} dV = \beta E_L V_0^{2+\alpha}/(2 + \alpha). \quad (4.7)$$

Comparing this equation with Eq. 4.1, the target capacitance is inferred to be

$$C_T = C_0 E_L V_{MV}^\alpha \text{ pF}, \quad (4.8)$$

where $V_{MV}$ is the instantaneous target voltage in megavolts, $E_L$ is the incident laser energy in kilojoules, and the constants are $C_0 = 9.4^{+4.7}_{-2.4}$ and $\alpha = 0.55 \pm 0.03$. This scaling is valid for laser energies in the range of 20 – 30 kJ and target voltages of 0.1-1.6 MV. The target capacitance is an implicit function of time since the target voltage evolves over the laser pulse duration. For peak target voltages of 0.1 - 1 MV and an incident laser energy of 20 kJ, the initial capacitance is approximately 50-190 pF. This result can be compared to the capacitance of two concentric spheres
Target capacitance \( (C_T) \), in units of picofarads per kilojoule of laser energy, as a function of the instantaneous target voltage. The dashed curves indicate the absolute uncertainty of the measured capacitance.

\[ R_I = \frac{V_T}{J A}. \] (4.9)

The current density is defined as \( J \equiv e(n_i v_i - n_e v_e) \), where \( e \) is the elementary charge and \( n_{i(e)} \) and \( v_{i(e)} \) are the ion (electron) density and velocity, respectively. The surface area of the expanding sheath front is \( A = 4\pi r_f(t)^2 \), where \( r_f(t) \) is the time-dependent radius of the front. The voltage in Eq. 4.9 will be solved for self-consistently in Sec. 4.3.3, while the current can be estimated using a fluid model and Poisson’s equation as follows.

Consider the density profiles of hot electrons and fast ions shown in Fig. 4.13. Such profiles have been previously simulated by authors that treated the expansion of cold ions \( (T_i \approx 0) \) by a hot-electron population.\(^{26,27} \) The hot-electron density follows the ion density until the local density scale length becomes comparable to the local Debye length \( (\lambda_{Dq} \sim 0.1 \mu m) \). The target voltage drops over both the quasineutral and sheath regions as indicated in Fig. 4.13. The current flows in the sheath region and this is understood as follows. As electrons cool, they give energy to the ions and fall behind (the Debye length decreases). The density of electrons at \( r_f \) increases in time and this amounts to a current. To compute the current, we take the flow velocity of the ions and electrons to be equal \( (v_i \approx v_e) \) and estimate this velocity using the flow velocity evaluated at \( r_q \). Thus, \( J \approx e(n_i - n_e)v_q \), and we then estimate \( n_i - n_e \) using Poisson’s equation, \( \epsilon_0 \nabla^2 \phi_s \approx 3\phi_s/\lambda_{Dq}^2 = -e(n_i - n_e) \), where \( \phi_s = V_T - \phi_q \) is the potential in the sheath region and \( \lambda_{Dq} \) is the Debye length evaluated at \( r_q \). We furthermore estimate the area using \( r_q(t) \) in place of \( r_f(t) \), which is a reasonable approximation since \( r_f(t) = r_q(t) + \mathcal{O}(\lambda_D) \), and \( r_q(t) \gg \mathcal{O}(\lambda_D) \). The resistance is then given by
4.3 Temporal Evolution of the Target Potential

Figure 4.13. Structure of the ion front, including the electron and ion densities and potential. The target potential ($V_T$) is the sum of the voltage drop across the quasineutral region ($\phi_q$) and across the sheath ($\phi_s$).

$$ R_I = \frac{V_T}{12\pi\epsilon_0 v_0 (t)^2 (V_T - \phi_q)/\lambda_{D,q}^2} \quad (4.10) $$

To proceed, we require solutions for the hydrodynamic variables (density, temperature, velocity and potential) and knowledge of the position $r_q$ at all times, which are described next.

In a 1D spherical geometry, the potential ($\phi$), temperature ($T_h$), density ($n$), velocity ($v$) and characteristic size ($R$) of the quasineutral region are given by:26,27

$$ \phi(r,t) = T_h(t)[r/R(t)]^2 \quad (4.11) $$
$$ T_h(t) = T_{h,0}[R_0/R(t)]^2 \quad (4.12) $$
$$ n(r,t) = n_0 [R_0/R(t)]^3 \exp[-(r/R(t))^2] \quad (4.13) $$
$$ v(r,t) = (2 q r t T_{h,0})/[m_i (R_0^2 + 2 c_{h,0}^2 t^2)] \quad (4.14) $$
$$ R(t) = (R_0^2 + 2 c_{h,0}^2 t^2)^{1/2} \quad (4.15) $$

Here, $n_0$ is the initial density of the fast ions, $R_0$ is the corona edge when hot electrons are produced, $c_{h,0} \equiv (T_{h,0}/m_i)^{1/2}$ is the initial hot sound-speed and $R(t)$ is the radius of the where the fast-ion density profile has fallen to $n_0/e$ (hence $R(t) < r_q(t)$). In the derivation of Eqs. 4.11-4.15, it is assumed that the electrons expand adiabatically. Hot-electron production is a complex process involving both energy loss and gain. Electrons that are trapped within the target reflux and lose energy to the unablated part of the shell. They also gain energy from electron-plasma waves (generated by the TPD instability at the quarter-critical surface) as they pass through the corona. Furthermore, electrons in the sheath region (Fig. 4.13) can lose energy and become trapped, while trapped electrons that gain sufficient energy from the corona can escape and drive ion acceleration. Simulations suggest that this process results in a Maxwellian distribution of hot electrons.28 The time-averaged hot-electron temperature ($T_h$) was measured from the hard x-ray spectrum produced when electrons stop on the shell and was found to vary linearly from 25-50 keV for incident laser intensities of $4-10 \times 10^{14}$ W-cm$^{-2}$, as discussed in Sec. 4.2.10*

* An analytical error was identified in the analysis used in Ref. 10 and when corrected the hot-electron temperatures...
is important can be readily seen by comparing the refluxing time to the decay time of the potential. We will show in Sec. 4.3.3 that the latter is on the order of 1 ns, while refluxing time of a hot electron is \( R(t)/v_e \lesssim 100 \) ps. We thus expect the temperature to be uniform in space and use \( \gamma = 5/3 \) since scattering in the shell relaxes the distribution to a Maxwellian over three degrees of freedom.

Next, we find \( r_q(t) \) by setting the local density scale length equal to the local Debye length \( (\lambda_{D,q} \nabla n_q/n_q = 1) \). Making use of Eqs. 4.11-4.15, we obtain a transcendental equation for \( r_q(t) \):

\[
 r_q(t) = \sqrt{2} R(t) \left[ \ln \left( \frac{R_0^{5/2}}{2r_q(t)\lambda_{D,0}R(t)^{1/2}} \right) \right]^{1/2},
\]

where \( \lambda_{D,0} \) is the initial Debye length. This equation is solved numerically for a given time value. Eqs. 4.11-4.16 are then evaluated at \( r_q \) and combined with Eq. 4.10 to obtain the resistance:

\[
 R_I = V_T × \frac{m_i R_0^4}{384 \sqrt{2} \pi \epsilon_0 e T_{h,0} \left( R_0^2 + 2c_{h,0}^2 t^2 \right)^{3/2} (V_T - \phi_q)} \times \ln \left[ \frac{R_0^{5/2}/2r_q(t)}{R_0^2 + 2c_{h,0}^2 t^2} \lambda_{D,0} \right]^{-5/2}.
\]

The resistance depends on the initial values of the hot-electron temperature, size of the sheath and the fast-ion density (through \( \lambda_{D,0} \)). The temperature is routinely measured and the initial sheath radius is estimated \( (R_0 \sim 1 \) mm) based on hydrodynamic simulations. The initial fast-ion density is just \( N_i/V_i \), where \( N_i \equiv C_T V_0/e \) is the number of ions and \( V_i \) is the volume they occupy. Initially, these ions will be confined to a spherical shell around \( R_0 \) with a characteristic shell thickness of \( \lambda_{D,0} \sim 1 \mu \)m. The volume was left as a fit parameter but was constrained to scale with the Debye length for different laser intensities and hence hot-electron temperatures \( (V_i \propto \lambda_{D,0} + \mathcal{O}(\lambda^2)) \) for the case of \( \lambda_{D,0} \ll R_0 \). The fit parameter was determined by comparing the model against Hicks’ measurement. As discussed in Sec. 4.3.3, good agreement was obtained for \( R_0 \sim 1.5 \) mm and \( n_0 \sim 10^{14}-10^{17} \) cm\(^{-3} \). The scaling of \( n_0 \) with Debye length and capacitance produced reasonable results for other cases.

**Stalk Inductance and Resistance**

The resistance and inductance of target stalks has been previously inferred using time-gated proton radiography.\(^{22}\) This technique\(^{25}\) uses a monoenergetic proton backlighter, driven by several OMEGA beams, to infer path-integrated magnetic and electric field strengths of a subject that is driven by a separate set of OMEGA beams (Fig. 4.14). The DD-protons generated from the implosion of the backlighter were used to image the silicon carbide and boron stalks that supported the imploding ICF capsules. The resolution of this technique is set by the characteristics of the backlighter fusion burn region, which has a Gaussian radial profile with a FWHM of 45 \( \mu \)m and a burn duration of 150 ps.\(^{29}\) This resolution was sufficient for the study of target stalks initially 100 \( \mu \)m in diameter (several cm in length) that evolve over several nanoseconds. Since the stalk was at an angle of 40° with respect to the imaging axis, the current flowing along the stalk generated a magnetic field that deflected incident protons and resulted in fluence modulations on the CR-39 detector.

Electric fields surrounding the stalk changed the energy of the backlighter protons measured by the CR-39. Magnetic and electric fields were thus inferred by imaging variations in the fluence and

\[ \text{are 40% lower [Source: private communication with C. Stoeckl]}. \]
energy of the backlighter protons using CR-39 detectors. The relative timing between the laser beams that drive the proton backlighter and the target was varied to sample the stalk at different times. Using this technique, radiographs containing information about the electric and magnetic fields surrounding the stalk were produced at different times, and the temporal evolution of the stalk was studied over several identical shots. In particular, the temporal evolution of the stalk plasma and current were thus studied over several shots for a target with a peak target potential of $\sim 100$ kV (irradiated at $4 \times 10^{14}$ W·cm$^{-2}$, see Fig. 4.15).

Simulations indicated that a current flowing in a cylindrical annulus surrounding the stalk reproduced the radiographs and that portions of the annulus were charged. The time-dependent inner and outer radii of the annulus ($R_i \sim 70-500$ μm and $R_o \sim 110-1100$ μm) and hence the expansion speed of the annuli were also inferred (see table 4.1 for a summary of the measured values). The expansion speed is consistent with either electrostatic acceleration (e.g., coulomb explosion of the stalk) or the thermal expansion of a 500 eV stalk plasma. Initially, the ionized outer layer of the stalk is cold ($T_s \sim 10$ eV) and ohmic heating is not large enough to heat this plasma to 500 eV. In fact, a self-consistent heat transfer calculation with an ohmic source (neglecting expansion, convection and diffusion which only cool the stalk plasma) shows that the temperature increases by about a factor 2 over 1 ns). We thus conclude that the stalk blows apart due to electrostatic forces rather than ohmic heating and subsequent thermal expansion. Given this result, we expect the stalk to expand slightly faster for larger drive voltages (target potentials). As a result, the time-dependent annuli from Ref. 22 were scaled for different initial voltages ($R_i$ and $R_o$ scale with the expansion velocity that is proportional to $\sqrt{V}$) and then used to calculate the inductance. The inductance and resistance of a wire with length $l_s$, inner radius $R_i$ and outer radius $R_o$ are given by:

$$L = \frac{\mu_0}{2\pi} \ln \left( \frac{R_o}{R_i} \right)$$

$$R = \frac{\rho}{l_s}$$

An incorrect time constant ($L_s/R_s$) was previously published by assuming that the expansion of stalk plasma was thermal-driven ($T_s \sim 500$ eV was incorrectly inferred and used in the calculation of $R_s$).
Figure 4.15. (a) Schematic of the target support stalk, skin plasma and the sheath boundary (left). The target potential generates a skin plasma around the stalk and drives an ohmic current through it. The magnetic fields generated by this current deflect protons from a backlighter (not shown) which can be detected and imaged at various times (right). These images of proton fluence variation contain information about the geometry of the current annulus and illustrate the temporal evolution of the skin plasma surrounding the stalk. In particular, the magnetic field structure (darker halo around the solid stalk shown in white) indicates that the current-carrying annulus expand away from the stalk center as the plasma evolves. (b) The laser pulse used to drive the target (and hence the stalk) in these experiments. See Ref. 22 for a detailed analysis and discussion of these proton radiographs.

\[ L_S = \frac{\mu_0}{2\pi} R_o \left[ \frac{l_s}{R_o} \ln \left( \frac{l_s}{R_o} + \sqrt{\frac{l_s^2}{R_o^2} + 1} \right) - \sqrt{\frac{l_s^2}{R_o^2} + 1 + 1} \right] \]  
\[ R_S = \frac{\eta l_s}{\pi R_o^2 [1 - (R_i/R_o)^2]} \]  

where \( \eta \) is the Spitzer resistivity (\( \propto T_s^{-3/2} \) with \( T_s \sim 10 \) eV). The time-dependent inductance and resistance curves for these stalks are shown in Fig. 4.16 for several initial drive voltages. The precision of the inferred annuli and hence the inductance and resistance are high (±50 μm). The

Table 4.1. Measured values associated with the current-carrying plasma surrounding the stalk, including the inner and outer radii (\( R_i \) and \( R_o \)), the peak current (\( i_{s,m} \)) and the maximum magnetic field (\( B_{max} \)).

<table>
<thead>
<tr>
<th>OMEGA Shot</th>
<th>Time, ns</th>
<th>( R_i ), μm</th>
<th>( R_o ), μm</th>
<th>( i_{s,m} ), kA</th>
<th>( B_{max} ), Tesla</th>
</tr>
</thead>
<tbody>
<tr>
<td>51244</td>
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<td>70</td>
<td>110</td>
<td>2</td>
<td>3.64</td>
</tr>
<tr>
<td>51246</td>
<td>1.9</td>
<td>470</td>
<td>750</td>
<td>5</td>
<td>1.33</td>
</tr>
<tr>
<td>51247</td>
<td>2.4</td>
<td>500</td>
<td>850</td>
<td>7</td>
<td>1.65</td>
</tr>
<tr>
<td>51250</td>
<td>3.4</td>
<td>500</td>
<td>1100</td>
<td>6</td>
<td>1.09</td>
</tr>
</tbody>
</table>
Figure 4.16. Stalk (a) inductance ($L_S$) and (b) resistance ($R_S$) as a function of time for several values of peak target potential. For these plots, $t = 0$ corresponds to the time when the voltage has reached its peak value, at which point the skin plasma surrounding the stalk is well-established. The current carrying stalk plasma expands in time, resulting in larger plasma cross section and hence lower inductance and resistance values. In addition, larger peak target potentials cause the stalk to blow apart faster, resulting in lower inductance and resistance for larger peak potentials. The values shown here are for an effective stalk length of 2 cm.

Systematic uncertainty arises from the fact that the stalk length that contributes to these circuit parameters is not well known. The portion of the stalk that sees an electric field and hence carries current is limited to the region within the sheath boundary. As discussed in Sec. 4.3.2, the sheath boundary is the fast-ion and hot-electron expansion front that expands at the hot-electron sound speed. We estimate that the boundary starts at the target radius and expands out to $\sim 1-2$ cm.

For simplicity, we use a time-averaged effective stalk length in subsequent calculations, and scale the length with the hot-electron sound speed and hence laser intensity. It will be shown in Sec. 4.3.3 that agreement was obtained with experimental data for a time-averaged effective stalk length of $\sim 2$ cm, as determined by fitting the model to experimental data.

Equation 4.19 for the inductance is derived by applying the Biot-Savart law to a thin hollow current profile of finite length. In contrast to the infinite wire approximation, the integral in this case does not diverge and there is no need to cutoff the integrand at an arbitrary radius. The dominant uncertainty in the inductance then arises from the fact that the return path of the stalk current is ignored (see Fig. 4.17). The positive charge that leaves the target is spherically distributed around the outside of the sheath boundary (the capacitance is distributed around the target in the real system). There are two reasons why this seemingly crude calculation is a reasonable approximation. First, the return path of the stalk current is $1-10$ millimeters away from the stalk, whereas the radius of the current around the stalk is $0.1-1$ mm (these range of

Figure 4.17. The return path for the stalk current (orange arrows) in the target-stalk circuit. In reality, the capacitance is distributed and hence the length of the return path is a continuous variable. Charge that leaves the stalk is distributed around the outside of the spherical sheath boundary (other side of the capacitor). The diode in series with $R_I$ is shown in this schematic to emphasize that the expanding fast-ion current discussed previously is unidirectional. When calculating the inductance, the return current contribution is ignored on the basis of geometric arguments.
values correspond to initial radii and the radii when the voltage has decayed). The return path is typically at a distance of 10 times the profile of the current around the stalk and hence we expect its effect on the fields surrounding the stalk to be small. In addition, the radiography technique employed here is used to infer a field strength from which the current is deduced. In a sense, any contributions from the return path are included in our measurement. Second, the inductance of the return path itself is small compared to the current along stalk due to the same geometric arguments. In particular, the aspect ratio ($l_s/R_o$) of the current-carrying stalk plasma is approximately ten times greater than for the return path. The uncertainty of this calculation and the sensitivity of the model to $L_S$ is further discussed in Sec. 4.3.4.

Finally, measurements of the stalk current were used to correct the capacitance Eq. 4.8. For targets with a peak voltage of $\sim 100$ kV, ($I_L = 4 \times 10^{14}$ W-cm$^{-2}$), the amount of charge that leaves the stalk is measured to be comparable to the amount of charge expelled from the target in the form of fast ions. Since measurements for higher laser intensities and hence larger driving voltage are not available, we take the amount of charge leaving the stalk to scale with the initial voltage, since a larger voltage will to first order drive a proportionally larger current.

### 4.3.3 Voltage Decay Time

The dynamics of the voltage decay are described by a nonlinear differential equation. The governing equations are derived here in the time-domain and then solved numerically. The charge on the capacitor ($Q_T$) depletes over time according to:

$$\frac{dQ_T}{dt} = -(i_A + i_S),$$

where the fast-ion current ($i_A$) and the stalk current ($i_S$) are given by:

$$i_A = V_T/R_I,$n
$$i_S = \left( V_T - L_S \frac{di_S}{dt} \right) / R_S,$$

where $R_S$ and $L_S$ are the time-dependent resistance and inductance of the stalk, and where $R_I$ represents the time-dependent parallel leakage path that allows for fast-ion acceleration. Eq. 4.22 implies that the inductance varies slowly compared to the decay time \([L_S(dL_s/dt)^{-1} \ll \tau_d]\), which is evident by comparing a $\sim 1$ ns decay time to the inductance curves in Fig. 4.16a. Combining Eqs. 4.20-4.22, together with the continuity equation ($i_C = i_A + i_S$) results in the following differential equation for the charge on the capacitor:

$$\frac{d^2Q_T}{dt^2} + \frac{R_S}{L_S} \frac{dQ_T}{dt} + \frac{d}{dt} \left( \frac{Q_T}{R_I C_T} \right) + \frac{Q_T}{L_S C_T} \left( \frac{R_S + R_I}{R_I} \right) = 0.$$  

No assumptions have been made at this point regarding the nonlinearity or time-dependence of the circuit elements, with the exception of the inductor, which is taken to vary slowly compared to the decay time. It is evident from inspection of Eq. 4.23 that the stalk inductance and resistance contribute a time constant to the system ($L_S/R_S$), while the target capacitance and resistance ($R_I C_T$) contribute another. Furthermore, the charge-leakage resistance modifies the bandwidth of the system. Next, we incorporate the relation between capacitance and charge ($C_T \sim V_T^2$) to obtain a differential equation for the voltage on the capacitor:
For a linear capacitor, the dynamics of the charge describe the behavior of the voltage. It is clear from Eq. 4.24 that the nonlinearity of the capacitance adds a second-order damping term (second term) and modifies the $R_I C_T$ time constant and bandwidth by a scalar factor.

The numerical solution to Eq. 4.24 was found using the values discussed in Sec. 4.3.2. In particular, the stalk length and initial ion density were swept to match the constraints on total decay time and the shape of the experimental measurements, respectively. The experiments utilized 900 μm diameter thin-glass shells irradiated with 1 ns square laser pulses with a total energy of 26 kJ resulting in on-target laser intensities of $9 \times 10^{14}$ W-cm$^{-2}$, and a peak voltage of 660 kV. The results of the numerical calculation are shown in Fig. 4.18a (solid curve) alongside the measured values and the square laser pulse. The absolute uncertainty of the voltage and time error bars of each data point are also shown, along with an estimate of what the voltage rise might look like. Agreement was obtained for an effective stalk length of $\sim 2$ cm (reduced $\chi^2 = 2$).

The effective stalk length was constrained by the measured data and by the total decay time. The target cannot begin charging before hot-electron production starts, as indicated by the arrows in Fig. 4.18, and we expect the rise time to be $\sim 10 - 100$ ps. Further increase in the effective stalk length increases the inductance of the system, decreases the current flow through the stalk and hence increases the decay time. This leaves little time for the voltage to rise and violates that constraint. The peak magnitude of current flowing through the stalk for a laser intensity of $4 \times 10^{14}$ W-cm$^{-2}$ is comparable to measurements by Manuel et al. ($i_s \sim 7$ kA). In addition, the total charge flowing through the resistor $R_I$ is consistent with the total measured charge of fast ions emitted by the target (see Sec. 4.2.2). At higher intensities, the model predicts 1.5-2 times more charge flowing through $R_I$ than is observed. This shortcoming is reasonable given the lack of stalk measurements at higher intensities, and hence the poor estimate of the capacitance as discussed at the end of Sec. 4.3.2.

Fig. 4.18b shows the relative magnitudes of the terms in Eq. 4.24. The bandwidth term is the dominant term for the first several hundred ps after which the $R_I C_T$ term takes over, while the nonlinear and stalk damping terms of Eq. 4.24 are negligible. The circuit is thus an under damped system and the decay shown in Fig. 4.18a is a quarter-cycle of an oscillation for the first few hundred ps. The remainder is damped by the $R_I C_T$ term, which corresponds to damping by fast-ion acceleration. It is worth noting that the shape of Hicks’ data alone also suggests that both the stalk and fast-ion current must discharge the target. If only the $R_I C_T$ term were dominant, the shape of $V_T(t)$ would be concave (e.g., decaying exponential) and does not match the measurement. Similarly, if only the bandwidth term were dominant, the shape of $V_T(t)$ would be convex (e.g., sinusoidal). The competition between these two terms flattens out $V_T(t)$ and agrees better with Hicks’ measurement. This was verified by heuristically changing circuit parameters to make one term more dominant than the other, as discussed further below (Sec. 4.3.4).

The voltage decay time decreases with decreasing laser intensity (see Fig. 4.18c). For lower laser intensities, the onset of hot-electron production occurs later relative to the start of the pulse. Decay curves are shown in Fig. 4.18c for an 860 μm diameter target irradiated with laser energies of 12, 18 and 26 kJ, corresponding to on-target intensities of 4, 6 and $9 \times 10^{14}$ W-cm$^{-2}$. These curves have been time-shifted to account for onset of hot-electron production and show that the
Figure 4.18. (a) The modeled decay of the target voltage (solid curve) is shown together with measurements by Hicks for a laser energy of 26 kJ and intensity of $9 \times 10^{14}$ W-cm$^{-2}$. Also shown are the laser pulse shape used in the experiments (dashed curve) and the target voltage rise. (b) Magnitude of the damping and bandwidth terms of Eq. 4.23 (c) Voltage decay curves for 860 μm diameter targets irradiated with a 1 ns pulse with laser energies of 12, 18 and 26 kJ, corresponding to on-target intensities of 4, 6 and $9 \times 10^{14}$ W-cm$^{-2}$. Arrows indicate the onset of hot-electron production.
4.3 Temporal Evolution of the Target Potential

Table 4.2. Parameters used in the calculation of the voltage decay curves of Fig. 4.18c, including the laser intensity ($I_L$), laser energy ($E_L$), measured hot-electron temperature ($T_{h,0}$), measured peak voltage ($V_0$), modeled stalk length ($l_s$) and initial ion density ($n_0$). Model predictions for the peak stalk current ($i_{s,p}$), ratio of the number of predicted to measured fast ions ($N_{i,p}/N_{i,m}$) and decay time ($\tau_d$) are also shown.

<table>
<thead>
<tr>
<th>$I_L$ (W-cm$^{-2}$)</th>
<th>$E_L$ (kJ)</th>
<th>$T_{h,0}$ (keV)</th>
<th>$V_0$ (MV)</th>
<th>$l_s$ (cm)</th>
<th>$n_0$ (cm$^{-3}$)</th>
<th>$i_{s,p}$ (kA)</th>
<th>$N_{i,p}/N_{i,m}$</th>
<th>$\tau_d$ (ns)</th>
</tr>
</thead>
<tbody>
<tr>
<td>$4 \times 10^{14}$</td>
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<td>24</td>
<td>0.08</td>
<td>1.46</td>
<td>$2.7 \times 10^{14}$</td>
<td>6</td>
<td>1.25</td>
<td>0.78</td>
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<td>32</td>
<td>0.33</td>
<td>1.72</td>
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<td>0.87</td>
</tr>
<tr>
<td>$9 \times 10^{14}$</td>
<td>26</td>
<td>44</td>
<td>0.66</td>
<td>2.00</td>
<td>$6.6 \times 10^{17}$</td>
<td>55</td>
<td>2.11</td>
<td>0.96</td>
</tr>
</tbody>
</table>

voltage goes to zero shortly after the end of the laser pulse. The parameters used in the calculation of Fig. 4.18c are summarized in table 4.2.

4.3.4 Sensitivity to Circuit Elements

The circuit elements each have an uncertainty associated with them since they are based on a measurement (or on a model, for the case of the fast-ion resistance). As demonstrated in Sec. 4.3.3, the dominant terms in Eq. 4.24 involve $C_T$, $R_I$ and $L_S$. The uncertainty in $C_T$ comes from the fast-ion measurements, while uncertainty in $L_S$ arises from the way it is calculated (return currents are ignored and the stalk length that contributes to the inductance is not known with precision). Finally, the dominant uncertainty for $R_I$ arises from the fact that the electrons are assumed to cool adiabatically. This is strictly not true and the precise uncertainty associated with this assumption is difficult to estimate. However, it is alleviated by the fact that the hot-electron temperature used in these calculations is the time-averaged temperature, which implicitly accounts for the sources and sinks in our problem. To address the effect of these uncertainties, it is instructive to examine how sensitively the results of the model depend on each of these circuit elements.

Shown in Fig. 4.19a are results of the model for the baseline case (dashed curve, with $E_L = 26$ kJ, $I_{14} = 9 \times 10^{14}$ W-cm$^{-2}$) alongside model outputs when $C_T$ is artificially altered by 50% (the uncertainty of $C_T$ is well within this range, as discussed in Sec. 4.3.2). It is evident that increasing (decreasing) $C_T$ by 50% increases (decreases) the decay time by about 100 (200) ps. The shape of the curve is however preserved and this is expected since a change in $C_T$ equally impacts the two dominant terms in Eq. 4.24 (the bandwidth and $R_IC_T$ terms). In particular, a larger $C_T$ reduces the magnitude of these terms and thereby increases the overall decay time.

Sensitivity to the stalk inductance is illustrated in Fig. 4.19b. A ten-fold decrease in $L_S$ decreases the decay time and adds slightly more curvature to the decay. The increase in $L_S$ enhances the bandwidth term. More bandwidth results in faster decay (charge leaves through the stalk faster) and also results in a decay curve that is slightly more sinusoidal in shape. Conversely, a comparable increase in $L_S$ increases the decay time and slightly flattens out the decay curve.

The effect of changing $R_I$ is shown in Fig. 4.19c. As in the case of sweeping the inductance, changing the resistance modifies the shape of decay curve drastically, but does not significantly change the decay time. A smaller resistance enhances the $R_IC_T$ term relative to the bandwidth term, resulting in a decay curve that is more exponential in shape.

These results illustrate several important points. First, even though the decay time predicted by the model is somewhat sensitive to changes in $C_T$, the uncertainty in $C_T$ is not large enough to significantly impact the model results. Second, since the measured shape of the voltage decay
Figure 4.19. Sensitivity of the circuit model to variations in (a) the target capacitance ($C_T$), (b) the stalk inductance ($L_S$) and (c) the fast-ion resistance ($R_I$). The dashed curves are the baseline calculation (nominal circuit values) for an 860-μm diameter target irradiated with a 1 ns pulse with a laser energy of 26 kJ (on-target intensity of $9 \times 10^{14}$ W-cm$^{-2}$).

The curve is linear (neither exponential or sinusoidal) both the $R_I C_T$ and bandwidth terms must have comparable contributions. In particular, this constraint fixes the relative magnitude of $L_S$ and $R_I$, since they alter the two terms and hence the shape, as demonstrated in Figs. 4.19b-c. Finally, compared to $C_T$, a relatively large variation in $L_S$ or $R_I$ is required to significantly impact either the shape or duration of the decay curve. This relative insensitivity forgives the systematic uncertainties in $L_S$ and $R_I$ that arise from neglecting return currents and from assuming that the electron-ion expansion is purely adiabatic. In other words, since the shape of the measured decay curve is well-reproduced using this model, the uncertainties in $L_S$ and $R_I$ cannot be larger than about a factor of 2 or the shape would not be well-modeled.

4.3.5 Implications for IFE

Since IFE requires the implosion of targets with high-repetition rates ($\sim$10 Hz),$^{32}$ the use of a support stalk is impractical. Current designs make use of targets without supporting stalk fibers, which are launched into a chamber at which point they are tracked and subsequently shot.$^{33}$ Since
the stalk plays an important role in discharging the target, we expect IFE targets to discharge over a longer period of time.

The absence of a stalk amounts to setting Eq. 4.22 to zero, preventing current flow through the stalk. This is equivalent to analyzing Eq. 4.24 in the limit \( L_S \to 0 \) and \( R_S \to \infty \), which results in:

\[
\frac{dV_T}{dt} + \left[ \frac{1}{(1 + \alpha) R_I C_T} \right] V_T = 0. \tag{4.25}
\]

This equation is solved numerically since \( R_I \) is a nontrivial function of time. The results are shown in Fig. 4.20 for an 860-\( \mu \)m diameter target irradiated with laser energies of 12, 18 and 26 kJ (corresponding on-target laser intensities of 4, 6 and \( 9 \times 10^{14} \) W-cm\(^{-2}\)). The voltage decay time increases with decreasing laser intensity, but the dependence is weak. The absolute time when the potential goes to zero varies significantly because the onset of hot-electron production varies with intensity. The weak dependence of the decay time on laser parameters can be understood as follows. For a lower laser intensity and hence a lower peak voltage, the capacitance is smaller and there is less charge to remove. However, the resistance is roughly proportionally larger because there are less charge carriers and this results in a weakly-varying \( R_I C_T \) time constant. This is most easily seen from the analytic expression for \( \phi_q \). Although this expression does not include charge separation effects, it gives a feel for scaling of the potential decay time. The time it takes for the voltage to drop to 1/e of its initial value (the no-stalk decay time or \( \tau_{d,ns} \)) can be found by inverting Eq. 4.11:

\[
\tau_{d,ns} \sim R_0/T_{h,0}^{1/2} \tag{4.26}
\]

Since the geometry is fixed for OMEGA targets (\( R_0 \sim 860 \) \( \mu \)m) the decay time decreases with increasing laser intensity, albeit weakly, as observed in Fig. 4.20. The scaling of \( \tau_{d,ns} \) with \( R_0 \) is confirmed by our numerical solution to Eq. 4.25, although it is not readily apparent from the expressions for \( R_I \) and \( C_T \).

Fig. 4.20 illustrates the effect of removing the stalk for OMEGA targets, but current designs for IFE-scale targets (i.e., direct-drive ignition-scale targets) will have different longer discharge curves. These targets are larger than those used at OMEGA by a factor of 2-3. They are irradiated at comparable on-target laser intensities (\( \sim 7 \times 10^{14} \) W-cm\(^{-2}\)) but with higher total laser energy (\( \sim 2 \) MJ). Under these conditions, the hot-electron temperature is \( \sim 2 \) times higher for a given intensity and hot-electron production is driven to saturation at \( \lesssim 1\% \) of the laser energy.\(^{14} \) Since the discharge time scales roughly with the geometry and hot-electron temperature for targets without stalks, we expect the decay times for IFE-scale targets to be roughly two times longer than the curves shown in Fig. 4.20. Detailed studies can be carried out once the hot-electron production for these targets, and hence the capacitance has been thoroughly characterized.

A longer decay time in the absence of a stalk implies that preheat in these targets will be higher than in stalk-mounted targets, which depends on the detailed dynamics of trapped electrons passing through the target shell. This will be addressed in the future using simulations that utilize these voltage decay curves. The role of the stalk is thus significant but the underlying physics is subtle. The \( L_S/R_S \) time constant alone does not play a large role since it is longer than the decay time. However, it is evident that charging the inductor (stalk) with current (as represented by the bandwidth term in Eq. 4.24) results in faster expulsion of charge from the target.
Figure 4.20. Voltage decay curves for (hypothetical) OMEGA targets without stalks. These curves are for 860-μm diameter targets irradiated with a 1 ns pulse with laser energies of 12, 18 and 26 kJ, corresponding to intensities of 4, 6 and $9 \times 10^{14}$ W-cm$^{-2}$.

4.4 Summary

The first measurements of the total energy carried by fast ions on OMEGA have been conducted. It has been shown that the total energy carried by fast ions is $\lesssim 1\%$ of the incident laser energy, and can be parameterized in terms of the peak target voltage. The peak voltage scales linearly with laser intensity for a range of targets, with a dependence on shell thickness and to some extent, shell material. It was observed that the target voltage is correlated with the hard x-ray signal, previously shown to scale with TPD-generated hot-electrons, confirming the mechanism behind charging of these targets at OMEGA. These measurements were used together with previous measurements of the total hot-electron energy to estimate the maximum amount of preheat for cryogenic OMEGA targets. Results presented here indicate that preheat is unlikely to play a significant role below an intensity $6 \times 10^{14}$ W-cm$^{-2}$. At an intensity of $7 \times 10^{14}$ W-cm$^{-2}$, target preheat is less than 20 J, corresponding to at most a 20-30\% $\rho R$ reduction. These simple estimates do not take into account the full hot-electron dynamics. Since this is larger than the maximum 10-20\% $\rho R$ reduction seen in current experiments on OMEGA, a more detailed study, incorporating the dynamics of hot-electron preheat is necessary. Essential to this study is knowledge of the hot-electron characteristics as well as the temporal and spatial behavior of the target potential that traps hot-electrons.

To this end, we presented a circuit model that describes the voltage decay dynamics of spherical, stalk-mounted targets shot on OMEGA. The circuit elements of the model were determined from previous spectral measurements of fast ions and radiography of target stalks. This model captures the essential physics of the decay dynamics and provides valuable insight into the role of the target support stalk in the discharging process. Furthermore, it can be incorporated into hydrodynamic codes or future PIC simulations of hot-electron preheat.

Some of the circuit elements are approximate and require refinement for detailed quantitative studies. Measurements of stalk characteristics at higher laser intensities and hence higher drive voltage are required to correct the capacitance and to verify the physics of stalk plasma evolution in this regime. Time-resolved measurements of the fast-ion current can also be used to further validate the resistance ($R_I$) formalism employed here. Finally, our model assumes that electrons are generated in a burst (that the capacitor is initially charged to a voltage and subsequently discharges), which is an approximation since electrons are produced over a few hundred picoseconds or an appreciable fraction of the total decay time. This model can thus be refined by incorporating a detailed time-dependent electron source and additional charge sinks such as electron stopping in the shell.

It has been shown on the basis of this model that the target stalk plays a significant role in the
voltage decay dynamics. In particular, it was shown that the voltage decay time depends weakly on the laser intensity and that it is longer for targets without stalks (IFE targets), implying that preheat could be more severe for these targets.

4.5 References

9. Private communication with J. Delettrez.


Diagnosing Cone-in-Shell Fast-Ignition Experiments on Omega using Fast Protons

5.1 Introduction

The fast-ignition concept\textsuperscript{1,2} has been described thoroughly in the literature as one alternative to direct-drive hot-spot ignition. As discussed previously in Sec 1.2.3, a high-energy (10\textsuperscript{15} W-cm\textsuperscript{-2}) laser is used to compress a cold shell containing fusion fuel to high areal densities (\(\rho R \sim 1\) g-cm\textsuperscript{-2}). A short-pulse, ultrahigh-intensity laser (10\textsuperscript{19} W-cm\textsuperscript{-2}) is then used to generate megavolt electrons to heat the core of the dense fuel assembly in a time short compared to hydrodynamic timescales. The use of two independent laser drivers for compression of the fuel assembly and subsequent heating of the core allows for higher target gains, in principle, for the same amount of driver energy. This is because large fuel masses can be assembled with slow implosion velocities and ignition is achieved through efficient coupling of short-pulse beam energy to the dense core.\textsuperscript{1} In comparison to conventional hot-spot ignition, the symmetry requirement of the fuel assembly in fast ignition is not as stringent, which relaxes the illumination uniformity and power balance constraints of the driver.

The success of this approach relies on the effective energy coupling between the short-pulse laser and the pre-assembled dense fuel. A high coupling efficiency (CE) depends on the generation of hot electrons and their transport and energy deposition to the dense fuel core. A potential problem is that the generation of energetic electrons will inevitably accelerate ions as well. Any energy coupling to ions is a direct loss channel that must be examined.

The acceleration of ions, and in particular protons, by electrostatic fields set up by LPI-generated hot electrons has been observed in both direct-drive\textsuperscript{3} and indirect-drive\textsuperscript{4} configurations with \(\sim 10^{14}\) W-cm\textsuperscript{-2} long-pulse beams. Protons and heavier ions produced by ultra-intense (\(\sim 10^{18}-10^{19}\) W-cm\textsuperscript{-2}) short-pulse LPI have also been studied extensively using flat-foil and cone targets. In short-pulse scenarios, laser-to-proton energy conversion efficiency, angular emission of proton beams, and effects of plasma scale length on proton acceleration have been studied.\textsuperscript{5–7} Proton measurements have also been used, in conjunction with ion expansion models,\textsuperscript{8–10} to infer the temperature of the LPI-generated hot-electron distribution that accelerates these protons.\textsuperscript{11,12}

In this chapter, we present the first measurements of fast protons in surrogate cone-in-shell fast-ignition experiments conducted at the Omega Laser Facility.\textsuperscript{13,14} In these experiments, a short-pulse laser was focused into gold cones to generate hot electrons and subsequently heat a pre-assembled
dense D$_2$ core, with the aim to increase the DD-neutron yield by raising the ion temperature. The neutron yield enhancement due to core heating has been measured to be a factor of $\sim 4$. This corresponds to a CE of 3.5%, which was inferred using simulations that matched the measured neutron yield.

In the context of proton acceleration, these experiments differ from previous work with cone targets and short-pulse lasers in that protons have been used here as a diagnostic tool (1) to determine the energy coupling to protons, a loss mechanism in fast-ignition experiments, and (2) to assess effectiveness of fast-ignitor coupling to the dense core. In particular, it is shown that of the laser energy coupled to hot electrons, about 5% is lost to protons. In addition, proton-inferred hot-electron temperatures suggest that the observed low coupling efficiency is partly due to electrons that are significantly hotter than expected.

This chapter is organized as follows: section 5.2 presents an overview of the experimental setup of fast-ignition experiments at Omega and charged-particle diagnostics used to measure proton spectra. In Sec. 5.3, proton spectra and maximum energies are presented, followed by a discussion in Sec. 5.4 of where and how the protons are generated. In Sec. 5.5, protons are used to infer the hot-electron temperature for these experiments. Sec. 5.6 concludes with a summary of the results.

5.2 Integrated Fast-Ignition Experiments at Omega

Experiments were performed at Omega using both the OMEGA (long pulse) and OMEGA EP (short pulse) lasers. In these experiments, 54 OMEGA beams delivering 18 kJ of UV light to the capsule were used to compress the target along a low adiabat ($\alpha \approx 1.5$), which was achieved using a short single picket, followed by a main drive pulse with a duration of approximately 2.7 ns. A single short-pulse ($\sim 10$ ps) Gaussian-shaped OMEGA EP beam was then brought to focus inside the OMEGA target chamber. At best focus, 80% of the beam energy was contained within a diameter of approximately 50 $\mu$m, resulting in a maximum, beam-averaged on-target intensity of $\approx 6 \times 10^{18}$ W-cm$^{-2}$. For these experiments, the OMEGA EP power and energy contrast were of order $10^6$ and $10^4$, respectively.

The targets for these experiments (Fig. 5.1) were re-entrant gold cones inside 40-$\mu$m-thick deuterated-plastic (CD) shells with a nominal diameter of 870 $\mu$m. The cones were either 1.2 or 1.8 mm in length and had an opening half-angle of 17º. The cone tips had flat-top geometry with variable tip thickness (5-15 $\mu$m) and tip diameter of 10 $\mu$m. The cone walls were 10 $\mu$m thick inside the shell and 50 $\mu$m thick outside. The shells were not gas-filled, leaving only the CD shell and the ablated material from it to undergo fusion. Additional target details can be found in Ref. 15.

Proton energy spectra were measured using both the magnet-based charged-particle spectrometer (CPS1) and Wedge-Range-Filter (WRF) spectrometers. These instruments utilize CR-39

![Figure 5.1.](image-url)
solid-state nuclear track detectors (SSNTD), which are known to provide information about the energy and species of the detected charged-particles.\textsuperscript{16} It has been shown recently, however, that there exists CR-39 piece-to-piece variability in its response to charged-particles.\textsuperscript{17} Thus, CR-39 alone cannot be used for accurate measurements of charged-particle spectra and must be paired with an additional particle dispersion mechanism. CR-39 is immune to EMP and to some extent to x-rays, making it ideal for short-pulse experiments such as those presented here.

CPS1 features a 0.1-mm slit and a 7.6-kG magnet for dispersion of charged-particles onto CR-39 detectors. These spectrometers are capable of measuring proton energy spectra in the range of 200 keV to 30 MeV. The low energy limit is set by filtering (directly in front of the CR-39), which is required to mitigate a very large flux of low-energy charged-particles that would otherwise scatter within the diagnostic and saturate the detector. The high energy limit is set by the magnet dispersion and detector arrangement. CPS1 is fixed to the OMEGA target chamber as shown in Fig. 5.2. In practice, the exponential energy spectra of short-pulse accelerated protons results in a large on-detector proton fluence at lower energies. This may cause saturation of the CR-39 detector at these energies, effectively raising the low-energy limit of this diagnostic. It is worth noting that CPS1 cannot resolve heavy ions because of the degeneracy between charge-state, mass and energy that exists for magnetic spectrometers.\textsuperscript{16} Filters constructed of mylar and aluminum are overlaid on the CR-39 to filter out these ions. Furthermore, any energetic heavy ions that penetrate the filters are separated from protons on the basis of the contrast and diameter of the tracks they leave on the CR-39.

The WRF spectrometers use CR-39 overlaid with a piece of wedge-shaped zirconia ceramic (ZrO\textsubscript{2}), in which the minimum particle energy required to penetrate the wedge varies along the thickness (dispersion) direction. Since the zirconia wedge cannot be made thinner than 40 μm, the low-energy instrument cutoff for measurement of protons is approximately 3-4 MeV. The WRFs are compact (5 cm across) spectrometers that are ideal in probing the implosion at several locations. Several (either 3 or 5) WRF modules, each consisting of two WRFs, were used at a single measurement location to obtain good statistics. Fig. 5.2 shows the azimuthal projection of the location of these spectrometers in the OMEGA target chamber relative to the short-pulse beam and target. The coordinate system is defined such that the pole (0°) corresponds to the direction of the short-pulse laser.

The WRF proton spectrometers were the primary diagnostics. These were fielded on nearly every shot, while CPS1 was fielded on a handful of shots to corroborate the WRF measurements and provide additional details of the spectrum at energies below the WRF low-energy cutoff. The spectrometers subtend small solid angles (1 μsr for CPS1 and 100 μsr for the WRFs). They measured protons accelerated normal to the CD shell surface for the locations shown in Fig. 5.2.
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Figure 5.3. Proton spectra measured with CPS1 (80°) on a fast-ignition shot and reference shot. In both cases, a gold cone-in-shell target was compressed using 54 OMEGA beams (∼20 kJ) and a low-adiabat laser drive. For the fast-ignition case, an EP short-pulse laser was fired, at peak compression of the target, to generate hot electrons and heat the dense core. These energy spectra were background subtracted, although some residual background is observed in the 4-7 MeV range. The gaps in spectrum at ∼1 MeV and ∼2.3 MeV are due to the instrument.

In addition, when fielded at 80°, the spectrometers measured protons accelerated nearly normal to the cone surface since that surface is nearly parallel to the spectrometer aperture due to the 17° cone opening half-angle.

5.3 Proton spectra and maximum energies

A typical proton energy spectrum from integrated experiments, acquired using CPS1 (OMEGA Shot 56971) is shown Fig. 5.3. Alongside this spectrum is the proton spectrum for a reference implosion (OMEGA Shot 56976), where a similar target was imploded using the same long-pulse configuration (∼20 kJ, 54 OMEGA beams) without any short-pulse core heating. It is well established that long-pulse LPI generate protons up to about ∼1 MeV, consistent with the data shown for the reference implosion. Nearly all of the observed energetic protons, however, arise from short-pulse LPI. These spectra exhibit a high-energy cutoff corresponding to the maximum path-integrated electric fields seen by the ions.

Proton energy spectra were measured down to approximately 200 keV using the CPS. As proton emission was anisotropic, it was difficult to measure the total energy lost to protons with precision. On the basis of measurements such as the one shown in Fig. 5.3, we estimate that the total energy carried by these protons is about 10 J, or about 1% of the incident short-pulse laser energy. This number can be compared to the previously inferred 20% coupling efficiency of short-pulse laser energy to hot electrons. In other words, approximately 5% of the short-pulse laser energy coupled to hot electrons is lost to the acceleration of ions.

The fact that the observed ions were protons (and not deuterons or heavier ions) was confirmed by simultaneous charged-particle measurements using CPS1 and WRF spectrometers. Since CPS1...
uses magnetic fields for ion dispersion, it can be shown that the inferred energy of an ion depends inversely on the assumed ion mass. Thus, the CPS1-inferred energy of a deuteron mistakenly identified as a proton will be twice as large as the actual particle’s energy. The WRFs have an opposite energy-mass dependence, whereby the inferred energy of a deuteron mistakenly identified as a proton will be lower than the actual particle’s energy. Thus, it is possible to constrain the particle species using these measurement techniques on the same shot and same polar angle. In particular, CPS1 and the WRFs measure particles at the same polar angle ($80^\circ$) but different azimuthal angles.

Since the target is composed of a CD shell and Au cone, these protons originate predominantly from hydrocarbon contaminants on the surface of the target (either from the cone or shell) that may or may not have been blown off during the implosion of the shell. In any case, the implications are that the protons do not significantly interact or scatter with the compressed shell. The cone-in-shell target conditions at the time when the short-pulse laser interacts with the cone are schematically illustrated in Fig. 5.4. Shown are the cone, the compressed D$_2$ core ($\sim$ 50-μm diameter), the blowoff plasma surrounding the target, the generated hot electrons and the accelerated protons. The relative timing between the short-pulse laser and the start of the long-pulse compression lasers was varied from shot-to-shot, but was typically 3 ns. At this point in time, the blowoff plasma from the ablated shell has expanded with the ion sound speed ($c_s \equiv \sqrt{T/m_i}$), resulting in a characteristic scale length of about 400 μm-1 mm for typical coronal temperatures of $\sim$ 2 keV. The blowoff plasma from the implosion of the shell thus surrounds the target, and it is expected that the fast protons are accelerated in the presence of this long-scale-length plasma.

The maximum proton energy is of interest since it scales directly with the temperature of short-pulse-generated hot electrons. Direct measurements of the maximum energy can therefore be used to qualitatively infer how the hot-electron temperature varies with experimental parameters. The maximum proton energy was measured at various locations around the implosion using the compact WRF spectrometers on several shots (Fig. 5.5). These data incorporate gold cones with 5-μm, 10-μm and 15-μm-thick tips and 10-μm tip diameters. The data obtained in the direction transverse to the short-pulse beam ($80^\circ$) scale with intensity. A $\chi^2$-analysis indicates that these data fit a normalized ponderomotive scaling ($\propto I^{1/2}$) at $80^\circ$ (reduced $\chi^2 = 0.96$). This further confirms that these protons are accelerated by short-pulse generated hot electrons. Since the maximum energies scale with intensity as expected from theory, these protons can also be used together with models to estimate a hot-electron temperature, albeit with some caveats (see Sec. 5.5).

In contrast to the transverse direction, the maximum energies of forward going protons ($0^\circ$) neither show such scaling nor a dependence on cone-tip thickness. In addition, the maximum energies of forward-going protons are lower compared to the transverse going protons. This is consistent with simulations, which indicate that for these experiments, the angular distribution of hot electrons is azimuthally symmetric and has a maximum at $57^\circ$ relative to the forward short-pulse beam direction and falls off at lower and higher angles. As a result, it is expected that fewer and less energetic protons would be observed in the forward direction even when the cone tip is intact, which is consistent with these measurements.

Forward-going protons are accelerated from the surfaces of the compressed shell or the surrounding blowoff plasma by hot electrons that have interacted with the compressed core and lost a significant amount of energy (Fig. 5.4). Some of the slower electrons are even ranged out in the core. The electron temperatures and previously measured $\rho R$ of the compressed shell ($\sim$ 150 mg-cm$^{-2}$) are consistent with this notion, as discussed further in Sec. 5.5. As a result, the velocity distribution of forward-going electrons have a lower maximum energy and empty regions in velocity-space, thereby reducing the energies of forward-going protons relative to transverse protons.

Several WRFs were used to obtain the average maximum proton energies at each location.
Figure 5.5. Maximum proton energies measured by WRF spectrometers at 80° and 0° (where 0° corresponds to the forward short-pulse beam direction). The data points shown are averages over many WRF measurements at one location. The error bars were computed from the standard deviation of these multiple measurements. At 80° the data show reasonable agreement with the ponderimotive hot-electron scaling. The max. proton energies for the forward beam direction (0°) neither show scaling with intensity nor dependence on cone-tip thickness.

The standard deviations of these measurements were used to compute the error bars shown in Fig. 5.5. Since the spatial separation between adjacent WRFs is of order several centimeters, the observed uncertainties in the data arise from the absolute measurement uncertainty of each WRF (±200 keV) and possible spatial variations in the maximum energy of the emitted protons. For the case of forward-going protons, the uncertainties are as large as ±2 MeV, which is larger than the absolute measurement uncertainty of the WRF spectrometers. Thus, we conclude that there are real spatial variations of the maximum proton energy for forward-going protons. These observed larger spatial variations could be the reason why the scaling with intensity is not readily apparent. Furthermore, these variations are consistent with (though not an indicator of) the presence of a stochastic process, such as electrons scattering in the compressed shell. For these reasons, it is difficult to estimate a hot-electron temperature from forward-going protons, as additional physics about the electron transport must be unfolded. We defer to only transverse-going protons when estimating hot-electron temperatures in Sec. 5.5.

5.4 Source of the protons

There is evidence that the observed protons are accelerated from the entire cone surface rather than the tip alone. The data presented throughout this paper were primarily obtained with 1.2-mm-long cones, with 10-μm or 40-μm tip diameters and variable tip thickness. On a few shots, cones with a length of 1.8 mm were also irradiated. Full spectral measurements of the fast protons were not available for the 1.8-mm-long cones, but the maximum proton energies for these cones were almost twice as high (∼8-9 MeV) as those of the 1.2-mm-long cones under the same laser conditions. In contrast, differences in the maximum proton energies were not observed between cones with 10-μm and 40-μm tip diameters (lengths of 1.2 mm). The cone length thus has an effect on proton acceleration, while the cone tip diameter does not.

The effect of the cone length on proton acceleration was also observed in the yield of (p,n) reactions, which come from fast protons reacting with the aluminum target chamber. The resulting neutron time-of-flight (nTOF) data shown are in Fig. 5.6. For these two shots the nTOF settings, laser drive and target parameters were identical with the exception of the cone length. The x-ray flash, which comes from the short-pulse beam hitting the cone, and the 2.45-MeV DD-neutron signals are characteristic of these implosions. In between these signals (shown in Fig. 5.6), there
5.4 Source of the protons

exist a number of smaller peaks associated with neutrons from (p,n) reactions. The occurrence of the first (p,n) events are consistent with the time-of-flight of the fast protons across the target chamber. The integral of these signals between the proton arrival time (e.g., ~ 300 ns for 7.5 MeV protons) through 900 ns (excluding the DD-n peak) was computed for three shots: 2 with 1.2-mm-long cones 1 with a 1.8-mm-long cone. The ratio of the integrals between the 1.2- and 1.8-mm-long cone data were found to be 2.0 ± 0.5 and 3.0 ± 0.3. These ratios are comparable to the increase in surface area of between the two cones (a factor 2.25). The cross section for (p,n) in aluminum is not available for the relevant proton energies (< 10 MeV), but the (p,*) cross section increases linearly from 1-5 MeV incident proton energy, and flattens out above that through 10 MeV. The observed increase in (p,n) reactions for 1.8 mm cones is thus attributed to a combination of higher fast-proton yields and variation of the (p,n) cross section with incident proton energy.

Throughout the course of these experiments, the timing between the long-pulse OMEGA and short-pulse OMEGA EP beams was varied to find the optimal timing of the EP beam for maximum core heating and yield. Optimal timing corresponds to core heating at peak compression of the cold dense core. For effective coupling of the short-pulse laser energy to the dense core, the cone

Figure 5.6. Neutron time-of-flight signal, showing the x-ray flash, 2.45 MeV DD-n signal, and neutrons from (p,n) reactions. For these two shots, all laser and target parameters were identical with the exception of the cone length, which was 50% greater, corresponding to 2.25 times more surface area. The ratio of the total (p,n) signal of these two cone lengths is ∼ 2−3, roughly proportional to the ratio of the cone surface areas.

Figure 5.7. Maximum proton energies measured by CPS1 and WRFs. The different CPS1 and WRF measurements (at 80°) show good agreement with one another, despite the fact that they sample different azimuthal angles. The solid line is a fit to the data (∝ I^{1/2}), The max. energies of the transverse protons depend on whether the cone tip is intact when the OMEGA EP short-pulse laser arrives at the tip. When not intact (open circles), the maximum energy of the transverse protons (and hence the fields that accelerate them) are smaller.
Figure 5.8. Maximum energies of forward-going protons as a function of the OMEGA EP arrival time at the cone tip relative to the start of the long-pulse drive. Forward-going protons show no significant dependence on whether the cone tip is intact when the OMEGA EP short-pulse laser arrives at the cone tip. The shock break-out time at the cone tip, which depends on cone-tip thickness, occurs between 3.65 ns and 3.7 ns, as indicated.

The cone tip must be intact when the short-pulse laser is fired. Shock waves launched into the fuel during compression by the long-pulse OMEGA beams lasers will eventually reach the cone tip, break through, and destroy it. In this scenario, we expect poor hot electron production and hence less energetic protons. The cone tip was intact for data shown in Fig. 5.5. For two shots, the timing between OMEGA and OMEGA EP was such that the tip was broken by the shock waves when the short-pulse arrived at the tip. Shown in Fig. 5.7 are data taken at 80° using CPS1, alongside with data from WRFs (80°). The CPS1 data are generally in excellent agreement with the WRF data. This is expected since these instruments are at the same polar angle. The two shots where the tip was broken are indicated by the open circles. The maximum proton energies were significantly lower (∼40%) when the tip was not intact.

The drastic effect of the cone tip’s destruction on electron production and subsequent proton acceleration was not observed in the forward-direction, as shown in Fig. 5.8. For two shots, the 10-μm-thick cone tip was shocked before the short-pulse arrived at the cone tip. The previously measured shock break-out time, which varies with tip thickness is indicated in Fig. 5.8. Thus, neither the presence of the cone tip nor the thickness of the tip (per Fig. 5.5) affect the acceleration of protons emitted in the forward direction.

5.5 Estimates of Hot-Electron Temperature

It has been suggested that the presence of a significant preformed plasma inside the cone can lead to a filamentation instability and hence self-focusing of the short-pulse laser, leading to higher hot-electron temperatures. The relativistic filamentation instability is seeded by hot spots in the short-pulse laser intensity profile. Initially, regions of higher laser intensity expel plasma electrons due to ponderomotive pressure. This causes a modulation of the plasma dielectric function ($\epsilon = 1 - \omega^2_{pe}/\omega^2$) and hence the refractive index ($N = \epsilon^{1/2}$) transverse to the direction of laser propagation, where more intense regions of light (less electrons) have a larger index of refraction. As the laser propagates through the pre-plasma, it is refracted towards regions of higher intensity, thereby reinforcing the instability. Filamentation and self-focusing can thus occur under these conditions if the laser is transported through an underdense plasma.

In particular, 2D simulations using the fluid code HYDRA show that for these experiments, a preformed plasma with a scale length of 100 μm is present within the cone at the arrival time of the short-pulse OMEGA EP laser (Fig. 5.9). The large preformed plasma, if present, is due the laser prepulse that arises from amplified spontaneous emission (ASE). The prepulse is characterized by the laser contrast, defined as the amplitude ratio of the main drive to the prepulse. For these experiments, the energy and power contrast were $10^4$ and $10^6$, respectively.

Hotter electron temperatures, due either to self-focusing or to another physical mechanism, result in more energetic electrons that would not stop in the core as intended, thereby lowering
the overall CE. Using the proton data presented in this work, we can place a lower bound on the initial hot-electron temperature to see whether the electrons are hotter than expected from the ponderomotive scaling.

The hot-electron temperature is estimated using a plasma expansion model, which links the temperature of an initial hot-electron distribution to the proton maximum energies. In particular, \( E_M = \alpha T_H \),\(^{21}\) where \( \alpha \) depends on the expansion model.\(^{8-10,27}\) In general, \( \alpha \) has a logarithmic dependence on the hot-electron density \( n_0 \) and the laser pulse duration. The choice of appropriate model depends only on the relative timescales of the laser pulse duration \( (\tau_l) \) and the transit time of electrons through the cone wall \( (\tau_e) \).\(^9\) For these experiments, \( \tau_l \sim 20 \tau_e \). Thus, during the first part of the laser pulse, the cone tip is completely populated with hot electrons generated from the preformed plasma on the inside of the cone. For the remainder of the pulse duration, the laser maintains the temperature of these electrons. After the pulse turns off, the electrons expand adiabatically, giving their energy to the ions. A 1D fluid model has been previously used to describe this process. This two-phase fluid model\(^9,27\) treats the laser as a source term that acts to maintain a steady temperature during the pulse (isothermal expansion), and then conserves energy between electrons, ions and the accelerating field thereafter (adiabatic expansion).

The two-phase model relates the hot-electron temperature to the maximum proton energy by the relation:

\[
T_H = E_M \times [2.5 + 0.92 \ln(\omega_{pi} \tau_l)]^{-1}
\]

where \( T_H \), \( E_M \), \( \omega_{pi} \), and \( \tau_l \) are the hot-electron temperature, max. proton energy, ion plasma frequency \( (\omega_{pi} \equiv [n_e e^2/m_p \epsilon_0]^{1/2}) \) and the laser pulse duration, respectively. This formula was interpolated from numerical simulations\(^{27}\) and applies for \( \omega_{pi} \tau_l \) in the range of 5-100. The maximum energies and laser pulse duration were measured for each shot, while \( n_{e0} \) and hence \( \omega_{pi} \) were estimated using a variation of a known method.\(^{11}\) First, we determined the number of hot electrons generated by the short-pulse laser. Recent experiments on OMEGA EP showed that the laser energy conversion efficiency to hot electrons is 20\% for such kilojoule-class short-pulse lasers,\(^{19}\) and that it is independent of the laser intensity. The number of hot electrons \( (N_e) \) is then found by dividing the laser energy converted to hot electrons by the average energy of the electrons, defined by the hot-electron temperature. For the experiments presented in this work, we estimate (self-consistently, from the results of this calculation) that \( N_e \) is about \( 10^{14} - 10^{15} \). Next we obtained the volume by taking the product of the surface area of the cone and the characteristic

Figure 5.9. Simulated electron densities of the preformed plasma inside the cone. This preplasma is formed by the low-intensity pedestal \((\sim 10^{12} \text{ W-cm}^{-2})\) of the EP laser (due to amplified spontaneous light emission) before the high-intensity pulse is generated. These simulations were performed by T. Ma [Source: W. Theobald et al., “Initial cone-in-shell fast-ignition experiments on OMEGA,” Physics of Plasmas, (2011)].
scale length along the expansion dimension, given by \( \approx c \times t_l \). The hot-electron density is then just the ratio of the number of hot electrons to volume. Since the plasma frequency ultimately depends on the hot-electron temperature through the density, Eq. 5.1 is transcendental and must be solved numerically.

It is important to recognize that the density computed here \((n_e0 \sim 10^{17} \text{ cm}^{-3})\) is an overestimate. As discussed in Sec. 5.4, the protons are predominantly accelerated from either the surfaces of the cone or within the blowoff plasma surrounding it. In this calculation, we assumed that the hot-electron density is uniform, which is generally not the case. We expect the hot-electron densities to be lower upstream of the cone tip, where the ions are accelerated. From Eq. 5.1, it is evident that for a given maximum proton energy, an upper bound on \(n_e0\), and hence on \(\omega_{pi}\), corresponds to a lower-limit on the estimated hot-electron temperature.

Even though these ion expansion models primarily apply to thin-foil experiments,\(^{11,12}\) they can be used in the context of this work. However, a major distinction between thin-foil experiments and those presented here must be considered to allow for a correct interpretation of the data taken in this work. The density scale length of the ion front where the protons are accelerated is very different in these experiments. The two-phase model used here assumes that the initial density scale length of this front \((n/\nabla n)\) is small in comparison to the hot-electron Debye length (see Sec. 1.4.3 and Fig. 1.11a). While this is true for typical thin-foil experiments with short-pulse lasers, in our case the scale length of the blowoff plasma in front of the cone is \(\sim 400 \mu\text{m-1 mm}\) due to the implosion of the shell. The effective density scale length seen by the accelerating protons is roughly of this order, whereas the hot-electron Debye length is \(\sim 20 \mu\text{m}\). In this case, the maximum proton energies are lower since they scale inversely with the initial density scale length at the ion front.\(^{28}\) To quantify this difference to some extent, it has been shown that the addition of a 100-\(\mu\text{m}\)-scale-length plasma at the ion expansion front (in a scale length otherwise dominated by the much smaller hot-electron Debye length), reduced the observed maximum proton energies by about 4 times.\(^{28}\) Thus, for a given hot-electron temperature, the proton energies from these experiments are much smaller than expected by the model because of the longer density scale length at the ion front. Hence, in applying the expansion model to these experiments, it is expected that the actual temperatures are much higher than the temperatures estimated using the model. While this may seem uncertain, the aim is not to pinpoint the exact temperature, but to show that it is significantly hotter than expected from the ponderomotive scaling.

We used protons at 80° to estimate the hot-electron temperature, as these protons demonstrated the expected scaling with intensity \((\propto I^{1/2})\). For each shot, Eq. 5.1 was solved numerically to determine the lower-limit of the hot-electron temperature, and the results are shown in Fig. 5.10. The error bars correspond to the uncertainty of the maximum proton energy measurement. Shown alongside these data is the ponderomotive vacuum scaling (for the case of negligible preformed plasma inside the cone). The temperatures determined here are factors of 2-3 higher than the vacuum scaling. If this increase in temperature is due entirely to laser self-focusing in the preformed plasma, this result corresponds to a factor of 3-10 enhancement of the incident laser intensity.

It is worth noting that OMEGA EP is known to produce maximum proton energies that are higher in comparison to those of other laser systems.\(^{29}\) In particular, it has been shown that for a fixed laser intensity \((\sim 2 - 8 \times 10^{18} \text{ W-cm}^{-2})\), the maximum proton energy increases as the pulse duration is increased from 1 ps to 10 ps.\(^{30}\) Observations indicate that the maximum proton energy on OMEGA EP increases faster with the laser pulse duration than models (for example, Eq. 5.1) predict. At present, there is no explanation for this observation. We speculate that the effect itself could be due to hotter electron temperatures (for instance, due to enhanced absorption or to hot-electron refluxing) for longer pulses (10 ps), or due to additional physics of the ion acceleration process that is not incorporated into the models at this point.
5.6 Summary

In this work, we have for the first time characterized the energy loss to fast protons in cone-in-shell fast-igniton experiments. We estimate that of order 10 J, or 1% of the short-pulse laser energy is lost to fast protons. It was shown that these protons are accelerated from the surface of the cone, rather than the cone tip alone.

Finally, we have used these protons to estimate a lower bound on the initial hot-electron temperature. These estimated hot-electron temperatures (500 – 900 keV) are hotter than predicted from the ponderomotive scaling by factors of 2-3. If the enhancement of the hot-electron temperature is due entirely to laser self-focusing, this result corresponds to a factor of 3-10 enhancement of the incident laser intensity.

If laser filamentation is indeed the underlying mechanism for the observed enhancement of the hot-electron temperature, the laser intensity cannot simply be reduced to lower the temperature to improve coupling. This is because the filamentation instability that causes self-focusing and hotter temperatures can also result in a large hot-electron emission angle that can increase the ignition energy requirements of the short-pulse laser (to provide the same energy density to the core, a greater amount of short-pulse laser energy is required). Instead, the laser contrast can be improved to avoid the formation of a significant pre-plasma within the cone and hence to mitigate the instability.

5.7 References

24. ENDF/B-VII.0 Library, see http://www.nndc.bnl.gov/.
Fast-Ion Generation using Ultra-Intense Lasers

6.1 Introduction

Acceleration of fast ions generated by ultrahigh-intensity ($\geq 10^{18}$ W-cm$^{-2}$) LPI is an important topic due to its applications in medicine, HEDP and table-top particle accelerators. Medical applications include generation of protons and energetic carbon ions for cancer therapy. HEDP applications involve the development of high-energy proton backlighters for high-resolution imaging studies (e.g., field structures in ICF implosions). A better and fundamental understanding of the underlying physics in generating and accelerating ions is also important for optimizing advanced fusion concepts, such as ion fast ignition.

Over the last decade, a number of experimental studies have been conducted on proton acceleration from thin foils. These protons are sourced from hydrocarbon contaminants that reside on the surface of the foil. The effects of target parameters (geometry and materials) and laser parameters (pulse duration, contrast and intensity) on proton acceleration have been extensively studied and compared to theory. The acceleration of heavy ions has been observed and enhanced by removing surface hydrocarbon contaminants from the rear of the target, which was accomplished by ohmic heating the foil to 1,000 K. In addition, monoenergetic heavy ions were generated ($C^{5+}$) by controlling the thickness of a carbon layer (several angstroms) on the surface of the foil. A rigorous explanation for the monoenergetic nature of the observed heavy-ion spectrum is lacking, and discrepancies between the maximum energy of ions with the highest charge-to-mass ratio (either protons for unheated foils or heavy ions for heated foils) and those of other ions has not been addressed until now.

The physics of heavy fast-ion acceleration was studied in this thesis to better understand previously published results. Unheated thin-foil targets (20-$\mu$m-thick) were irradiated at intensities of $\geq 10^{19}$ W-cm$^{-2}$ and the resulting proton and heavy-ion spectra were measured. It is demonstrated here that the mean and maximum energies of the heavy ions scale with only the charge state for several ion species. Consistent with past work, it is observed that the maximum heavy-ion energies are lower than that of contaminant protons. In addition, the signature of a two-temperature electron distribution is observed, and we discuss for the first time how this distribution explains the measurements in this work and results of others, including monoenergetic heavy-ion production using thin layers of a specific material.

The heavy-ion measurements were taken at the MTW laser facility at LLE. The MTW laser is capable of delivering on-target energies as high as 10 J of IR light (1053 nm) in 1 ps. On-target intensities as high as $5 \times 10^{19}$ W-cm$^{-2}$ with a contrast of $10^8$ are possible at best focus, where 50% of the laser energy is focused to a 5 $\mu$m diameter spot (see Sec. 1.3.3 for more information). In
addition to being a platform for developing diagnostics for OMEGA EP\textsuperscript{16} and damage testing of optics, the MTW is being used for studies of LPI in the context of backlighter development and basic science. A Thomson parabola spectrometer (TPIS, see Sec. 3.3.2),\textsuperscript{17} previously designed, constructed and implemented on the MTW, was used for measurements of charged particle energy spectra presented in this chapter.

This chapter is organized as follows: In Sec. 6.2, heavy-ion experiments and results are presented, followed by a discussion of the ion-inferred hot-electron temperatures and the implications of these results (in Sec. 6.3). Sec. 6.4 concludes by discussing future measurements and some features of the detection scheme which are still being developed.

### 6.2 Heavy Ions Accelerated from Thin-Foil Targets

Initial experiments on the MTW laser were conducted using 500 \(\mu\)m \(\times\) 500 \(\mu\)m square aluminum and copper foil targets with a thickness of 20 \(\mu\)m, as shown in Fig. 6.1. The foils were irradiated with a laser energy of about 8.5 J, delivered in 1 ps. The full width at half maximum (FWHM) focal spot size was 5 \(\mu\)m in diameter, resulting in on-target intensities between 4-5 \(\times\) \(10^{19}\) W-cm\(^{-2}\). The foils were oriented in such a way that the TPIS measured energetic protons accelerated normal to the rear of the foils.

Resulting data from the CR-39, after a 20 minute etch in a 6 N solution of NaOH held at 80\(^\circ\) C, is shown in Fig. 6.2a. Data from the IP detector is shown in Fig. 6.2b. These data were acquired on different shots where either the CR-39 detector or the IP detector was fielded. Several charge states of carbon and aluminum ions are present in the CR-39 data, while protons are not observed due to a short etch time. Longer etch times could not be applied to the CR-39, as that would have caused significant overlap (saturation) of the heavy-ion tracks. From the parabolas shown in Figs. 6.2a-b, the particle distribution as a function of the parabola slope (\(\alpha\)) can be determined by establishing the parabola boundaries and then binning the number of track counts along that parabola. A typical N(\(\alpha\)) distribution for these shots is shown in Fig. 6.3, which spans the entire range of possible \(\alpha\)-values for the aluminum target. Several strong lines are present as well as many weaker ones; the expanded inset of Fig. 6.3 shows a view of the \(\alpha\)-spectrum and highlights the weaker lines. These background-subtracted N(\(\alpha\)) distributions were essential to identify the weakest ion species as they are not visible by eye in the detector images.

Fig. 6.3 shows that the TPIS, operating at an electric field of 15 kV-cm\(^{-1}\) with an aperture of 400 \(\mu\)m, was not able to differentiate some of the high-Z ions (e.g., C\(^{5+}\) and N\(^{6+}\)). To resolve this a smaller aperture should be used to decrease the parabolic line broadening. Alternatively, a stronger electric field (e.g., 40 kV-cm\(^{-1}\)) can be used if the signal is weak. The observed smaller line widths for decreasing \(\alpha\) (increasing q/m ratio of the ions) are a result of magnetic focusing, which arises from spatial non-uniformities in the field strength.

Since each parabola is divided into constant energy bins with each bin spanning a small but finite distance along the spatial direction of magnetic displacement, the background level is determined...
from just above and below the parabola. The background is interpolated through the parabola in the direction of electric displacement (see Sec. 3.3.4 and Appendix E for details). Figures 6.4 and 6.5 show the results from this background subtraction analysis for several ion species from aluminum and copper flat foils, respectively. Figure 6.6 shows proton energy spectra measured by image plates for four MTW shots of comparable laser intensities and targets (3280, 3281, 3285, 3286). Shown in these background-subtracted spectra are the error bars arising from the statistical uncertainty (95% confidence limit) of the measurement. The highest proton energies measured by the IP agree with theoretical predictions$^9,11$ for these targets and laser intensities. Predictions for aluminum or carbon ions at different charge states are not available at this time; empirical scalings are presented in the following section.

Figure 6.2. (a) Data recorded by a CR-39 detector in the TPIS from MTW shot 3117 showing various heavy-ion parabolas. In this image, lighter regions indicate higher track density. The lookup table was artificially saturated to enhance the visual appearance of parabolas. Note the localized areas of noise on the detector (e.g., bottom center of image). The TPIS was operated at 30 kV for this shot. (b) Data recorded by the IP detector for MTW shot 3285, showing carbon and proton parabolas. The TPIS was operated at voltage of 20 kV for this shot.

Figure 6.3. Histogram of the number of tracks recorded on the CR-39, binned in $\alpha$-space, for a $500 \times 500 \times 20 \, \mu m$ Al flat foil irradiated with a laser intensity of $5 \times 10^{19} \, W/cm^2$. An expanded view of the boxed region is shown as an inset, with several more weaker lines and a few overlapping parabolas. All labeled ions were separable with the exception of Al$^{9+}$ and C$^{4+}$. Some localized noise has also been identified and removed (by correlating the N($\alpha$) histogram with N(x,y) images).
Figure 6.4. Sample heavy-ion energy spectra for MTW shot 3117, obtained using CR-39. Target was a $500 \times 500 \times 20 \ \mu m$ aluminum foil, irradiated with a laser intensity of $5 \times 10^{19} \text{W-cm}^{-2}$. The dashed red lines are fits to the data (see Sec. 6.3 for discussion).

The spectra shown in Figs. 6.4 and 6.5 from aluminum and copper targets, respectively, illustrate a strong signal of certain ions and charge states over others. Note that the spectra from the aluminum target exhibits strong $C^{1+}$ and $C^{2+}$ lines while the data from the copper target does not. As mentioned above, it has been previously shown that ultra-intense laser beams incident on

Figure 6.5. Sample heavy-ion energy spectra for MTW shot 3118, obtained using CR-39. Target was a $500 \times 500 \times 20 \ \mu m$ copper foil, irradiated with a laser intensity of $4 \times 10^{19} \text{W-cm}^{-2}$. The dashed red lines are fits to the data (see Sec. 6.3 for discussion).
flat foils preferentially accelerate protons and hydrocarbons. To increase the heavy-ion energy for medical and ion fast-ignition applications, techniques have been developed to mitigate energy coupling to contaminant ions (e.g., protons) using cleaning techniques and specific target materials. In these experiments, no effort was made to remove these hydrocarbon contaminants. The difference illustrated here in the acceleration of carbon ions between aluminum and copper is likely due to the variation of contaminant levels between targets.

Despite this difference between targets, there are trends in the mean and maximum ion energy of these ions, as illustrated in Fig. 6.7. The mean energies were determined by integrating (and normalizing) the spectra of Figs. 6.4 and 6.5 along with the proton spectra of Fig. 6.6. The precision of these measurements is limited by three factors. First, the energy resolution of the TPIS, when run with a 400-μm-diameter aperture, has a worst-case uncertainty of approximately 6% for 20 MeV proton equivalent energies; this can be as low as 1% for 300 keV protons. Second, proper alignment of CR-39 during a microscope scan is limited to offsets as high as 200 μm in the magnet dispersion direction, resulting in uncertainties as high as 3.8%. The third contribution is the uncertainty arising from counting statistics, which in these data was no higher than a few percent. These three uncorrelated quantities were used to compute the error bars (within 95% confidence limits) shown in Fig. 6.7. These upper-bound estimates neglect the effect of the electric field on resolution and are thus conservative estimates. The estimates of maximum energies shown in Fig. 6.7 include a systematic observational bias of about 20% (not shown in the error bars) arising from the inability to visually discern an exact cutoff; this is evident in the spectra shown in Figs. 6.4 and 6.5. For the proton energy spectra (Fig. 6.6) the cutoff is also difficult to determine. It was taken to be the knee of the spectra between 15 and 20 MeV. There are fewer than ~200 protons above this defined cutoff, which is significantly lower than the number of protons under either of the two slopes in the spectra.

It has been suggested that protons, being lighter, are preferentially accelerated over heavier ions and thus will deplete the space charge initially setup by hot escaping electrons. This would lead to a scaling of the bulk ion energy with the q/m ratio rather than the charge, since lighter ions with higher charge states would be accelerated first, lowering the space-charge field for subsequent ions. These data, which include several species of aluminum, copper, carbon, oxygen and nitrogen ions (Fig. 6.7), show excellent scaling between mean energy and charge rather than with q/m ratio. A chi-square analysis shows that the mean energy for all ions is best described by the linear fit

\[ E_{\text{Mean}} = 0.38 \times Z - 0.07 \]  

(6.1)

with \( \chi^2 = 0.98 \). The scaling suggests that the electrostatic fields driving the acceleration do not
Figure 6.7. Mean ion energy and maximum energy vs. charge for protons and heavy ions. The proton data were obtained from image plates over four shots of comparable intensity (3280, 3281, 3285, 3286). Heavy ions shown here include several charge states of aluminum, copper, carbon, nitrogen and oxygen acquired over two flat-foil shots (3117 and 3118). The maximum energies of heavy ions scale linearly with charge (shown here as the charge number, Z), while for protons the energies are significantly higher and do not fit this trend. The maximum energies are subject to a 20% systematic observational bias (the shown error bars represent only the statistical uncertainty).

deplete over the acceleration phase of these ions. The maximum energies show a similar scaling for the heavy ions. For these ions, a chi-square analysis shows that maximum energy scales as

\[ E_{\text{Max.}} = 0.99 \times Z - 0.09 \]  

with \( \chi^2 = 0.99 \). The maximum energies of heavy ions are significantly higher than the mean, reflecting the fact that the bulk of the ions experience much lower fields. The maximum energies of protons are significantly higher than those of the heavy ions and do not fit the same linear trend.

6.3 Ion-Inferred Hot-Electron Temperatures

The observed discrepancy in maximum energies between protons and heavy ions may be explained by the presence of a two-temperature electron distribution. Simulations\(^{19,20}\) and measurements of the bremsstrahlung x-ray spectrum generated by the electrons\(^{21}\) have shown the presence of a two-temperature electron distribution for comparable targets and laser intensities and the two slopes in the proton spectra of Fig. 6.6 suggest that a two-temperature distribution is prevalent in these experiments.

In the two-temperature theoretical framework,\(^{19,20}\) the two distributions are characterized by a hot (\( T_H \)) and a cold (\( T_C \)) temperature (the term “cold” is used in a relative sense here, as these electrons can actually be quite hot), with respective number densities (\( n_H \) and \( n_C \)) and pressures (\( p_H \) and \( p_C \)). As the laser propagates through the underdense pre-plasma, a small fraction of electrons are directly heated by the laser. For a short-pulse interaction these electrons do not have time to reach thermal equilibrium with the bulk electrons. Thus, a two-temperature electron distribution with \( n_H \ll n_C \) and \( \nabla p_H / n_H \gg \nabla p_C / n_C \) is possible. The hot component will generate stronger sheath fields than the cold component (because the sheath fields \( \propto \nabla p / n \)). However, since \( n_H \ll n_C \), the charge separation is weaker for the hot component than for the cold component. Given these limits, the hot component will accelerate a small fraction of the protons to higher maximum energies; these protons will then quickly shield the sheath fields associated with the hot component since the charge separation is weak by assumption. Heavier ions and the bulk of the slower protons will experience fields generated by the colder component and therefore have lower maximum energies. Furthermore, if the cold component has enough charge separation, the heavy-ion mean and maximum energies will scale as a function of the charge number alone, and not the charge-to-mass ratio.

Simulations\(^{20}\) have shown that for \( n_C / n_H \approx 200 \) and initial temperature ratios of \( T_H / T_C \approx 900 \),
the hot-electron temperature falls on a time scale much faster than the cold component\textsuperscript{20}. In addition, 3D PIC simulations\textsuperscript{19} have shown that for $n_C/n_H \approx 100$ and $T_H/T_C \approx 50$, a smaller population of energetic ions are accelerated from the rear of the target in association with the hot component while slower ions are accelerated due to the cold component. In these experiments, we estimate a number ratio ($N_C/N_H$) of $\sim 200$ by integrating the left and right sides of the two-slope proton spectra of Fig. 6.6. These spectra were then averaged and fit to a model\textsuperscript{22} to infer the hot and cold electron temperatures (see Fig. 6.8). As discussed previously in Sec. 1.4.3, the slope of the proton spectrum scales with the electron temperature as $(ZT)^{-1/2}$. The intermediate region in the spectrum ($\sim 4 - 5$ MeV) has contributions from both the cold and hot component and is not fit. From the fits, we estimate $T_H/T_C \approx 30$ ($T_C = 72 \pm 5$ keV and $T_H = 1.9 \pm 0.1$ MeV), which is of the same order as the simulations. The experimental results are consistent with a simulated hot-electron component that drives fewer energetic protons to high energies for a short period of time and then decays away quickly, leaving a cold component with strong charge separation to accelerate heavy ions for an extended duration. This notion is further reinforced by analysis of the heavy-ion spectra. These spectra were fit to a single-temperature model, as indicated by the dashed red lines in Figs. 6.4 and 6.5. The average temperature deduced from these fits ($60 \pm 9$ keV) is comparable to the temperature inferred from the low-energy component of the proton spectrum. This further indicates that the hot-electron component accelerates a few protons while the cold

**Table 6.1.** Summary of cold-electron temperatures inferred from the proton and heavy-ion spectra. The reduced $\chi^2$ ($\chi^2_{\text{red}}$) value is shown for each fit. The average cold-electron temperature inferred from the heavy ions is $60 \pm 9$ keV.

<table>
<thead>
<tr>
<th>Ion</th>
<th>Z</th>
<th>Z/A</th>
<th>$T_C$, keV</th>
<th>$\chi^2_{\text{red}}$</th>
<th>Ion</th>
<th>Z</th>
<th>Z/A</th>
<th>$T_C$, keV</th>
<th>$\chi^2_{\text{red}}$</th>
</tr>
</thead>
<tbody>
<tr>
<td>Hydrogen</td>
<td>1</td>
<td>1</td>
<td>72</td>
<td>0.99</td>
<td>Oxygen</td>
<td>1</td>
<td>0.08</td>
<td>48</td>
<td>0.96</td>
</tr>
<tr>
<td>Aluminum</td>
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<td>0.04</td>
<td>40</td>
<td>0.82</td>
<td>Copper</td>
<td>6</td>
<td>0.09</td>
<td>165</td>
<td>0.99</td>
</tr>
<tr>
<td>Aluminum</td>
<td>3</td>
<td>0.11</td>
<td>21</td>
<td>0.96</td>
<td>Copper</td>
<td>13</td>
<td>0.20</td>
<td>77</td>
<td>1.00</td>
</tr>
<tr>
<td>Aluminum</td>
<td>4</td>
<td>0.15</td>
<td>62</td>
<td>0.99</td>
<td>Copper</td>
<td>14</td>
<td>0.21</td>
<td>71</td>
<td>1.00</td>
</tr>
<tr>
<td>Aluminum</td>
<td>5</td>
<td>0.19</td>
<td>72</td>
<td>0.99</td>
<td>Copper</td>
<td>15</td>
<td>0.23</td>
<td>66</td>
<td>0.99</td>
</tr>
<tr>
<td>Aluminum</td>
<td>8</td>
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<td>28</td>
<td>0.81</td>
<td>Copper</td>
<td>17</td>
<td>0.26</td>
<td>58</td>
<td>0.99</td>
</tr>
<tr>
<td>Carbon</td>
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<td>0.17</td>
<td>30</td>
<td>0.98</td>
<td>Copper</td>
<td>22</td>
<td>0.34</td>
<td>45</td>
<td>0.99</td>
</tr>
<tr>
<td>Carbon</td>
<td>3</td>
<td>0.25</td>
<td>78</td>
<td>1.00</td>
<td>Copper</td>
<td>26</td>
<td>0.40</td>
<td>38</td>
<td>1.00</td>
</tr>
</tbody>
</table>

**Figure 6.8.** The proton spectra obtained from four MTW shots (Fig. 6.6) were averaged (solid curve) and fit to a two-temperature model (dotted curves). The dashed curves indicate the uncertainty in the proton spectrum due to counting statistics. The hot-to-cold electron temperature ratio was inferred to be about 30.
component drives the rest of the ions. The temperatures inferred from fits to the heavy-ion spectra are summarized in Table 6.1.

In calculating these fits, it is assumed that the hot-electron temperatures are constant. This is a reasonable assumption while the laser is on but is violated thereafter. Once the laser turns off, electrons cool adiabatically and transfer additional energy to the ions. If most of the ion acceleration occurs while the laser is on, then the temperatures inferred from these time-integrated measurements are a good representation of the initial temperature. Fuchs and Robson demonstrated that this is in fact the case for comparable targets and laser conditions.\textsuperscript{9,11} They conducted an extensive study of the maximum proton energy and showed that the isothermal model alone reproduced the measured maximum proton energy if the laser pulse duration in the calculation was taken to be 1.3 times the actual pulse duration. This implies that the adiabatic expansion and ion acceleration that occurs after the laser turns off is a modest contribution to the overall expansion (on the order of 30%).

The above result suggests that the two inferred temperatures are in fact real and not an artifact of the time dependence of the expansion, since most of it occurs under isothermal conditions (i.e., the colder inferred temperature is not an artifact that results from the decay of the hot-electron temperature during the adiabatic phase of the expansion). This can be also be seen from the ion energy distribution one expects from a purely adiabatic expansion. Though the latter problem is not analytically tractable except in the case of an initial Gaussian density profile (not the case for these experiments), it illustrates more generally that the slope of the distribution function is always proportional to the initial temperature of the electrons, even when the electron-temperature is time-dependent. Using the formalism of Ref. 23, we find that the distribution function is a function of time:

\[
\frac{dN}{dE} \propto \exp \left[ -E / \left( \frac{2mc^2_{h,0}}{(x_0^2 + 2c^2_{h,0}t^2)} \right) \right]
\]  

(6.3)

where $E$ is the ion energy, $x_0$ is the thickness of the foil and $c_{h,0}$ is the initial hot-electron sound speed ($\equiv [ZT_{h,0}/m_i]^{1/2}$). The slope of the spectrum is thus time-dependent during the expansion but converges for late times. In particular, when $c_{h,0}t \gg x_0$ (a condition easily satisfied in these experiments, since $x_0 = 20 \mu m$ and $c_{h,0} = 50$ cm before the ions reach the spectrometer) the slope of the energy spectrum converges to $(ZT_{h,0})^{-1/2}$. Thus, even for the case of a time-decaying hot-electron temperature, the spectrum measured at the spectrometer would have a single slope that is a function of the initial hot-electron temperature. We therefore conclude that two temperatures are inferred and that the impact of the temperature time-dependence is small (though present).

It has been well-established that the hotter electron distribution is produced by the ponderomotive force of the laser acting on electrons (see Sec. 1.4.2). It is not clear how the colder electron distribution is produced and no analytical models exist to date. The colder temperature predicted by PIC simulations (for comparable laser and target parameters) is about 50 keV,\textsuperscript{19} in rough agreement with our measurement. We rule out vacuum heating and resonance absorption as possible mechanisms for the observed colder temperature since they require p-polarized light that was not present in these experiments.

These data have important implications for heavy-ion fast-ignition and for cancer therapy because they suggest that acceleration of a specific ion species (and charge state) is not easily controlled. Even if the acceleration of contaminants or undesired elements are suppressed by novel methods, coupling of energy into a pure beam of a pre-specified $q/m$ ratio is not trivial. As previously mentioned, no effort was made in these experiments to suppress the acceleration of protons or hydrocarbons. It has been demonstrated in the past that resistive heating of targets reduces the
number of unwanted accelerated hydrocarbons and that by coating the rear of the target one may preferentially accelerate ions of a particular element (but not charge state).\textsuperscript{24,14} Even for the case of heated targets, observations showed preferential acceleration and energy coupling to ion species with the highest $q/m$ ratio. Furthermore, two electric fields were inferred in those experiments: the first was a strong, short-lived field associated with the acceleration of ions with the highest $q/m$ ratio, and the second was a weaker field with a longer timescale that drove all other ions.\textsuperscript{24} We make the connection between the observations of those experiments and the results presented in this work. The two inferred electric fields observed in those experiments with heated targets can be explained by a two-temperature electron distribution that follows the density and temperature ordering described above: the stronger electric fields associated with the hotter temperature component accelerate ions with the highest charge-to-mass ratios, whether they be proton contaminants or heavy ions (see Fig. 6.9). These ions deplete the space charge associated with the hotter component and all other heavy ions are accelerated to lower energies. If the number of protons or heavy ions is small in comparison to the number of hot electrons, then those ions are accelerated at once resulting in a monoenergetic spectrum.

The results have an adverse impact in the context of the aforementioned applications. For cancer therapy, an ion with a given energy will have less penetrating power (and damage more surface tissue) when its charge state is higher. For ion fast-ignition applications, a higher charge state beams are more prone to scattering and result in spatial divergence of the generated ion beam, less localized power deposition, and higher ignition requirements. It is therefore desirable to increase the energy and number of ions without increasing the charge number.

6.4 Summary

Heavy-ion acceleration from the rear of thin-foils has been studied at the MTW laser facility. In initial experiments utilizing flat-foil targets at on-target intensities of $4-5 \times 10^{19}$ W-cm$^{-2}$, energetic protons (15-18 MeV max.) as well as several energetic heavy ions, including species of carbon, aluminum, copper, nitrogen and oxygen were observed. Empirical scalings of the mean and maximum energies of heavy ions were determined.

The discrepancy in maximum energies between protons and heavy ions, and the scaling of the heavy-ion energies, is explained by a two-temperature electron distribution where the hot component has weak charge separation. It has also been shown here that the mean energies of heavy ions scale linearly with only the charge of the ion implying that the space charge associated
with the colder electrons is not readily depleted as ions expand. Future experiments can be used to study how electron temperature and number density ratios affect ion acceleration, and how these ratios scale with laser parameters (e.g., contrast, intensity, etc.).

6.5 References


Conclusion

In this thesis, fast-ion measurements have been used to study aspects of direct-drive hot-spot ignition and fast-ignition implosions and laser-foil experiments. Spectral measurements of these ions were conducted using existing and new ion diagnostics that were in part developed and calibrated at MIT’s LEIA Facility.

LEIA has undergone many hardware upgrades that include the implementation of a new ion source, a multi-channel analyzer with pre-amplifiers and surface-barrier detectors and imaging diagnostics. These hardware upgrades have allowed for better control and characterization of the ion beam and hence the fusion products source. Software upgrades, including a beam-target physics simulation capability, have allowed us to verify the multi-channel analyzer calibration and to measure the energies fusion products to within a few percent. These improvements have allowed the development and calibration of charged-particle and neutron diagnostics, excluding the fast-ion diagnostics developed for this thesis which are discussed below.

Existing fast-ion diagnostics were used to measure fast protons while new instruments were developed to measure different charge states of heavy ions. In particular, fast protons were measured using the compact WRFs (≳ 3–4 MeV) and CPS (≳ 100 keV). The CR-39 used in these diagnostics was calibrated on LEIA, together with individual wedge-shaped zirconia filters used within the WRFs. The limitations of these spectrometers lead to the development of two Thomson Parabola spectrometers, which were developed in collaboration with LANL (TPIE) and SUNY Geneseo (TPIS) to measure heavy fast ions at the Omega and MTW laser facilities. These spectrometers use both magnetic and electric fields to momentum and energy analyze particles, allowing them to resolve ions with different charge-to-mass ratios. The energy resolution depends on the operating parameters but is on the order of a few percent, while the energy calibration is robust to within 3% for TPIS and 10% for TPIE, as verified by accelerator experiments and particle-tracking simulations. The electric fields produced by the electrodes (≲ 20 kV-cm⁻¹ for TPIE and ≲ 40 kV-cm⁻¹ for TPIS) were simulated using finite element analysis and detailed CAD models of the spectrometers, while the magnetic fields were provided by Dexter Magnetics, Inc. These field maps were incorporated into MIT and LANL particle-tracking simulations that were used to verify the calibration of the instruments. TPIE was originally developed for use on the OMEGA EP laser and was ported to the OMEGA laser to measure fast heavy ions from ICF implosions. On OMEGA, it was found that a significant number of low-energy ions (≲ 200 keV) scatter within TPIE and contribute a uniform background on the CR-39 detector, which was removed with the implementation of a new filtering capability to TPIE. Data reduction techniques were developed for TPIE and TPIS and a software package was written to implement analysis techniques for these instruments.

The new diagnostic capabilities were used to measure fast ions from direct-drive implosions on OMEGA. These ions are accelerated by TPD-generated hot electrons and indicate that ICF targets
charge to megavolt potentials. The large potential traps energetic electrons that can preheat the fuel, raise the adiabat and degrade compression. An estimate of the amount of preheat in OMEGA targets was obtained using previous measurements of the total hot-electron energy and the fast-ion energy measured in this thesis. For on-target laser intensities less than $6 \times 10^{14} \text{ W-cm}^{-2}$, the amount of preheat is negligible (though hot electrons are still produced and the target can charge to $\sim 200 \text{ keV}$). At an intensity of $7 \times 10^{14} \text{ W-cm}^{-2}$, target preheat is less than 20 J, corresponding to at most a 30-40% $\rho R$ reduction. Refined preheat estimates will require simulation of the hot-electron dynamics, which will require knowledge of the temporal evolution of the target potential. The fast-ion measurements were used in conjunction with radiographic images of the target support stalk to build an empirical model of how the target potential evolves in time. It was shown that the decay time predicted by the circuit model is in reasonable agreement with (limited) measurements of the potential by Hicks et al. (laser intensity of $\sim 9 \times 10^{14} \text{ W-cm}^{-2}$ and energy of 26 kJ). At a laser intensity of $4 \times 10^{14} \text{ W-cm}^{-2}$, the predicted stalk current and total fast-ion charge were also inline with measurements. At higher intensities, the predicted fast-ion charge escaping the target is larger than observed by $\gtrsim 75\%$ and radiographic measurements of the stalk are not available for comparison. The lack of stalk measurements at higher intensities could play a role in this discrepancy, since the stalk inductance at higher intensities is scaled from measurements taken at an intensity of $4 \times 10^{14} \text{ W-cm}^{-2}$. The circuit model was used to calculate the discharge time for OMEGA targets and shows that the voltage decay time increases with laser intensity (assuming a fixed target diameter of 860 $\mu\text{m}$) and has a value of about 1 ns for an intensity of $9 \times 10^{14} \text{ W-cm}^{-2}$. The decay time is predicted to be about twice as long for OMEGA targets without stalks and weakly dependent on laser parameters. This result suggests that preheat is more severe for IFE targets, which do not utilize target support stalks since they are shot with high repetition rates ($\sim 10 \text{ Hz}$). It was shown that the decay time for targets without stalks scales with the target radius. Since IFE-scale targets are a factor of 2-3 larger than those used at OMEGA, we expect decay times of 4-6 ns. This circuit model can be incorporated into hydrodynamic simulations or PIC calculations of hot-electron dynamics to obtain refined preheat estimates.

Measurements of fast protons were also used as a diagnostic tool to study fast-ignition experiments on Omega. In these experiments, 54 OMEGA beams were used to compress a CD shell using a low-adiabat laser pulse ($\rho R \sim 150 \text{ mg-cm}^{-2}$). Just before peak compression, a short-pulse ultra-intense laser was used to generate hot electrons which were intended to couple to the compressed CD shell, raise the ion temperature and increase the DD-neutron yield. Previous measurements indicated that about 20% of the short-pulse laser energy is coupled to hot electrons and that about 20% of this energy (4% of the total short-pulse energy) is coupled to the pre-assembled dense core. The poor coupling efficiency was diagnosed using measurements of fast protons that were accelerated by escaping hot electrons. Existing magnetic and compact proton spectrometers (CPS and WRFs) were used to measure the fast protons at several locations around these implosion. These protons are sourced from surface contaminants found on the target and have maximum energies of $\lesssim 7.5 \text{ MeV}$ and yields of $10^{13} - 10^{14}$. Spectral measurements of the protons were used to determine the amount of hot-electron energy coupled to protons, and it was shown that about 20% of the hot-electron energy (5% of the total short-pulse laser energy) is lost to the acceleration of fast protons. In addition, measurements of the maximum proton energies were used to infer a lower-limit on the hot-electron temperature. It was shown that the hot-electron temperature is greater than $\sim 500-900 \text{ keV}$ (a factor of 2-3 times hotter than expected) leading to poor coupling of hot electrons to the dense core. Electrons with energies of $\gtrsim 500 \text{ keV}$ can readily penetrate through the dense core ($\rho L \sim 300 \text{ mg-cm}^{-2}$), escape and accelerate ions consistent with measurements of the fast protons.

A majority of the thesis focused on using fast ions to diagnose ICF implosions. The generation
of fast ions using short-pulse lasers is also of interest due to its applications in many fields of science (e.g., ion fast-ignition, backlighters for radiography and medicine). The acceleration of protons has been previously studied and there exists an extensive literature on the influence of target and laser parameters. The physics of heavy fast-ion acceleration was studied in this thesis. Thin planar targets (20-µm-thick foils composed of Cu or Al) were irradiated at intensities of $\gtrsim 10^{19}$ W-cm$^2$ and the resulting proton and heavy-ion spectra were measured using the TPIS. It was shown that a wide range of ion species are accelerated from these targets, and that the mean and maximum energies of the ions scale with only the charge state. In particular, a two-slope proton spectrum was measured, consistent with a two-temperature hot-electron distribution. It was shown that the hot component has weak charge separation while the colder component has much stronger charge separation. The ratio of the temperatures was estimated to be about $\sim 30$ and the total number of hot electrons was inferred to be about 200 times lower than the bulk electrons. Such a distribution was used to explain measurements of the maximum and mean energies of protons and heavy ions. The maximum proton energy was observed to be much higher than that of the heavy ions, while the mean energies of protons and heavy ions were found to scale with the charge of the ion. The hot-electron component accelerates a small number of protons preferentially (since they are lighter) to very high energies ($\sim 20$ MeV) before the fields associated with this component are screened by these protons. The bulk of the electrons, with much stronger charge separation, are then capable of accelerating the remaining protons and heavy ions to mean energies that scale with the charge state of the ion, rather than the charge-to-mass ratio. These results have implications for ion fast-ignition and medical applications, where energetic singly-ionized heavy ions are more desirable since they scatter less and result in more localized power deposition.

Measurements of fast ions were demonstrated to be a valuable tool in diagnosing ICF implosions and laser-foil experiments. In future experiments, measurements of the fast ions can be used to estimate preheat, diagnose aspects of short-pulse coupling efficiency in fast ignition, or to measure fast ions from laser-foil interactions.
Appendix A

LEIA Specifications and Benchmarks

Specifications and performance benchmarks of the LEIA facility are presented in this Appendix. Various operating limits and specifications for the ion beam, vacuum system, fusion products source and charged-particle diagnostics suite are summarized in table A.1. The minimum and maximum values shown (where applicable) should never be exceeded for normal operation.

The typical values shown in table A.1 are obtained from averages over many different shots between 2009-2012. These values were obtained from the LEIA database which holds performance data acquired from each shot. These data are used to aid in the maintenance of the facility and for proper documentation of experiments. Each shot is assigned a number based on the shot date and shot number of that day (e.g., shot 2012101702 corresponds to the second shot of the day on October 17th, 2012). For each shot, the beam time and total run time are logged, alongside beam parameters and setup (distance, angle, etc) of the SBDs and other detectors. A sample shot report is shown in Fig. A.1. The first page of the report contains basic information pertaining to run time, beam parameters and detector setup. The second page holds setup information and data acquired by the charged-particle diagnostics suite (comprised of the SBD, preamplifier and MCA), including the SBD(s) used, the position, angle, filtering, aperture and counts on each detector. The charged-particle spectrum acquired from each SBD is also recorded.

These shots reports can be queried using any combination of the fields shown in the report. In addition, the database interface can generate cumulative statistics using these data. These statistics include fusion count rates (Fig. A.2), beam-on-target time (Fig. A.3), LEIA utilization for various tasks (Fig. A.4). In addition, the database contains separate tables for the target and SBDs that are used on LEIA. Each entry holds nominal specifications of the SBD or target and also holds real-time information about usage (e.g., hours of beam time on a particular target or SBD).

The information contained in the LEIA database has proven to be useful for diagnosing and resolving beam or other hardware or software failures.
Table A.1. Specifications and operating limits for various LEIA subsystems, including the ion beam, vacuum system, charged-particle diagnostics suite and fusion products source.

<table>
<thead>
<tr>
<th>Parameter</th>
<th>Min.</th>
<th>Typical</th>
<th>Max.</th>
<th>Units</th>
</tr>
</thead>
<tbody>
<tr>
<td><strong>Ion beam</strong></td>
<td></td>
<td></td>
<td></td>
<td></td>
</tr>
<tr>
<td>Ion species</td>
<td>-</td>
<td>D$_2$, $^3$He</td>
<td>-</td>
<td></td>
</tr>
<tr>
<td>Beam energy</td>
<td>70</td>
<td>140</td>
<td>150</td>
<td>kV</td>
</tr>
<tr>
<td>Beam current</td>
<td>10</td>
<td>20</td>
<td>60</td>
<td>μA</td>
</tr>
<tr>
<td>Beam size (on target)</td>
<td>3</td>
<td>5</td>
<td>10</td>
<td>mm</td>
</tr>
<tr>
<td>Probe voltage</td>
<td>-</td>
<td>5</td>
<td>6</td>
<td>kV</td>
</tr>
<tr>
<td>Extractor voltage</td>
<td>-</td>
<td>4</td>
<td>5</td>
<td>kV</td>
</tr>
<tr>
<td>Focus voltage</td>
<td>-</td>
<td>4</td>
<td>5</td>
<td>kV</td>
</tr>
<tr>
<td>Source pressure</td>
<td>10</td>
<td>50</td>
<td>100</td>
<td>mTorr</td>
</tr>
<tr>
<td>Pressure control resolution</td>
<td>-</td>
<td>&lt; 1</td>
<td>100</td>
<td>mTorr</td>
</tr>
<tr>
<td><strong>Vacuum</strong></td>
<td></td>
<td></td>
<td></td>
<td></td>
</tr>
<tr>
<td>Acc. base pressure (no gas flow)</td>
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<td>$5 \times 10^{-7}$</td>
<td>Torr</td>
<td></td>
</tr>
<tr>
<td>Acc. base pressure (nominal gas flow)</td>
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<td>$&lt; 1 \times 10^{-5}$</td>
<td>Torr</td>
<td></td>
</tr>
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<td>Chamber base pressure*</td>
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<td>$5 \times 10^{-8}$</td>
<td>Torr</td>
<td></td>
</tr>
<tr>
<td>Chamber pump down time (&lt; 10$^{-3}$ Torr)*</td>
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<td>5</td>
<td>min</td>
<td></td>
</tr>
<tr>
<td>Chamber pump down time (&lt; 10$^{-5}$ Torr)*</td>
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<td>15</td>
<td>min</td>
<td></td>
</tr>
<tr>
<td><strong>Fusion products</strong></td>
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<td></td>
</tr>
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<td>DD reaction rate</td>
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<td>$10^7$</td>
<td>s$^{-1}$</td>
</tr>
<tr>
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<td>$10^6$</td>
<td>s$^{-1}$</td>
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<td>Count rate stability§</td>
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<td>1</td>
<td>-</td>
<td>%</td>
</tr>
<tr>
<td>DD-p linewidth†</td>
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<td>150</td>
<td>-</td>
<td>keV</td>
</tr>
<tr>
<td>D$_3$He-p linewidth†</td>
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<td>140</td>
<td>-</td>
<td>keV</td>
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<td>Source position uncertainty</td>
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<td><strong>Charged-particle detection</strong></td>
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</tr>
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<td>Channels</td>
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</tr>
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<td>MeV</td>
</tr>
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<td>Energy resolution</td>
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<td>-</td>
<td>keV</td>
</tr>
<tr>
<td>Absolute energy uncertainty</td>
<td>-</td>
<td>± 75</td>
<td>-</td>
<td>keV</td>
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<tr>
<td>SBD line broadening†</td>
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<td>Pre-amp/MCA line broadening†</td>
<td>-</td>
<td>67</td>
<td>-</td>
<td>keV</td>
</tr>
<tr>
<td>Total line broadening‡</td>
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<td>70</td>
<td>-</td>
<td>keV</td>
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<td>Pre-amp gain</td>
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<td>MCA A/D resolution</td>
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<tr>
<td>MCA A/D sampling rate</td>
<td>-</td>
<td>100</td>
<td>-</td>
<td>MHz</td>
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</tbody>
</table>

* Applies for empty target chamber only
† Full width at half maximum (FWHM)
‡ Measured using an $^{241}$Am source
§ Over a 20-minute interval with feedback gas control
Figure A.1. Sample shot summary from the LEIA database. The first page (top) contains shot header information, beam and target information, detector setup and miscellaneous notes. The second page (bottom) contains information about the SBD setup, including the SBD(s) used, the position, angle, filtering, aperture and counts on each detector. The charged-particle spectrum acquired from each SBD is also recorded.
Figure A.2. Shot-averaged DD and D$^3$He count rates from 2009-2012. The mean and maximum DD count rates are $1.4 \times 10^6$ s$^{-1}$ and $5 \times 10^7$ s$^{-1}$, while the mean and maximum D$^3$He count rates are $8.9 \times 10^4$ s$^{-1}$ and $1.4 \times 10^6$ s$^{-1}$.

Figure A.3. The beam-on-target time for each LEIA shot from 2009-2012.
Figure A.4. LEIA utilization for various diagnostic development activities from 2009-2012. Total utilization was approximately 668 hours and includes beam-on-target time and shot setup.
Appendix B

A Multithreaded Modular Toolkit for Control of Complex Scientific Experiments

B.1 Introduction

In the course of building a new experimental apparatus, one is typically faced with the challenge of designing and building software for control and data acquisition. Depending on the scale of the experiment, several options are available to the experimenter. One such option is to create a control scheme from scratch, using a programming language and operating system of choice, with the use of helpful guides\(^1\) and various libraries. Another option is to use a development toolkit; such toolkits, which are available for various platforms, simplify the design process to a variable extent. They come in a number of flavors:

1. Device-driver development toolkits (DDKs),\(^2,3\) which simplify hardware communications code, allowing the experimenter to focus on the software framework for the control application itself.

2. Simplified integrated development environments (IDE), with simplified proprietary programming languages and proprietary user-interface controls.\(^4–6\)

3. Application-specific IDEs, that utilize standardized programming languages (such as C) and typically furnish the developer with generic libraries for hardware control.\(^7\)

4. Complete control solutions, which are open-source and commercial software packages specifically designed for industrial control.\(^8,9\)

Several options exist for the experimenter both for control and data acquisition; here we focus solely on the former, although there is no reason why the toolkit described here cannot be used for the latter. Note that while this toolkit does not perform functions in real-time, it is intended to interface to and monitor such controllers; a number of software packages are readily available for real-time applications.\(^10,11\)

These options each have strengths and weaknesses. DDKs ease driver development but do not provide any framework for the application which must be written from scratch. Such DDK’s for UNIX-like systems are also difficult to find.

Simplified IDEs facilitate development by introducing a simplified programming language and toolkit which encapsulates threading and low-level implementation details, such as opening communication ports and sockets. Though this eases the development of a control system, making it attractive to, if not ideal, for novice users, encapsulation of implementation details means that
Table B.1. Various toolkit types that may be used to build a full control solution. Comparison is made between the license type, operating system (OS), whether the toolkit is intended for distributed control, whether it is modular, whether it has supervisory components (e.g., interlocks) and if it has an infrastructure which eases multithreading.

<table>
<thead>
<tr>
<th>Toolkit Type</th>
<th>DDK</th>
<th>Simplified IDE</th>
<th>Application-Specific IDE</th>
<th>Complete Control Solution</th>
<th>MCT</th>
</tr>
</thead>
<tbody>
<tr>
<td>Example(s)</td>
<td>MS WDK</td>
<td>LabVIEW</td>
<td>NI LabWindows</td>
<td>RSView</td>
<td>EPICS</td>
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<td>License</td>
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<td>Windows</td>
<td>Cross-Platform</td>
<td>Open-source</td>
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<tr>
<td>Operating System(s)</td>
<td>Windows</td>
<td>Windows, OS X,</td>
<td>Windows</td>
<td></td>
<td>Open-source</td>
</tr>
<tr>
<td>Distributed</td>
<td>-</td>
<td>Limited (e.g.,</td>
<td>No</td>
<td>Yes</td>
<td>No</td>
</tr>
<tr>
<td></td>
<td></td>
<td>network variables)</td>
<td></td>
<td></td>
<td></td>
</tr>
<tr>
<td>Modular</td>
<td></td>
<td>Yes</td>
<td>No</td>
<td>Yes</td>
<td>Yes</td>
</tr>
<tr>
<td>Built-in Supervisory</td>
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<td></td>
<td>No</td>
<td>No</td>
<td>No</td>
</tr>
<tr>
<td>Component</td>
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<td></td>
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<td>No</td>
<td>Yes</td>
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<tr>
<td>Multithreading</td>
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<td>Yes</td>
<td>No</td>
<td>Yes</td>
<td>Yes</td>
</tr>
<tr>
<td>Infrastructure</td>
<td></td>
<td></td>
<td>No</td>
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</tr>
<tr>
<td>Language Support</td>
<td>C/C++</td>
<td>Proprietary</td>
<td>C</td>
<td>C</td>
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<tr>
<td>IDE</td>
<td>No</td>
<td>Proprietary</td>
<td>No</td>
<td></td>
<td>No</td>
</tr>
</tbody>
</table>

advanced debugging, control of thread and process execution, thread synchronization and mutual-exclusion of shared data structures is difficult*. Furthermore, since the programming languages and toolkits are proprietary and closed-source, one is locked-in to using propriety debugging and optimization tools from the vendor.

Application-specific IDEs typically provide the developer with convenient libraries (though often close-source) for hardware access but lack a re-usable framework for control applications. Thus, while hardware access may be readily achieved, significant amounts of time must be spent implementing an infrastructure for multithreading, centralized monitoring and supervisory functions (e.g., interlocks).

Finally, a number of open-source and commercial software packages are available which are designed specifically for industrial control. One such package is the Experimental Physics and Industrial Control System (EPICS). It has the advantage of being open-source with a built-in distributed infrastructure, making it ideal for large-scale systems. Though it is powerful and scalable, it is not well suited for novice users and lacks built-in supervisory components; it is also in excess of what is required for a small to medium sized facility (non-distributed control system, with tens of hardware components to control) such as the MIT Linear Electrostatic Ion Accelerator (LEIA). Another alternative, RSView32, is a commercial package with automation and control functionality similar to that of the MCT; it is however a closed-source package, making it less extensible, with limited operating-system support.

The Modular Control Toolkit (MCT) incorporates the strengths of each of the approaches discussed and is designed with control applications in mind. It is open-source, written in C++ and built using open, portable libraries such as GTK+ and GLib which are available under the Lesser General Public License (LGPL). It provides a modular framework, allowing one to isolate and re-use significant portions of the code as gradual changes are made to an experiment or for an entirely different experiment altogether; modularity readily allows for the systematic testing and debugging of complex code. Users have the option of launching and managing threads with

*The author recently discovered, using the toolkit described herein, a bug in the kernel driver of a particular Ethernet serial device server when it was accessed simultaneously by two threads; such a discovery is difficult, if not impossible, with heavily encapsulated systems where thread execution control is less transparent.
complete control over execution, or using the pre-existing threading infrastructure and allowing the framework to manage and control execution. Furthermore, features common to a medium or large-scale experimental apparatus, such as interlocks, are built-in to the framework. Finally, module development and framework customization is not restricted to an integrated development environment or compiler. Features of the different toolkit varieties, including the MCT, are compared in table B.1.

This appendix is organized as follows: Section B.2 gives the reader an overview of the various software components which make up the toolkit and discusses implementation details of the code; section B.3 presents some of the implementation details of a sample module; section B.4 discusses some of the development challenges and the initial deployment of the toolkit at MIT; section B.5 describes the future direction of the toolkit.

B.2 Overview and Implementation

The MCT was originally developed for use on the MIT LEIA but with a broader scope of application in mind. The specifications for the toolkit may be understood by considering the fact that several elements are common to most control systems software.

Among these include the need for interlocks to help ensure the safety of the apparatus and proper execution of the experimental process. At MIT, for example, the toolkit is used to ensure that vacuum gate valves are not inadvertently opened by an operator when doing so could put an unnecessarily high load on a pump and lead to potential failure of that pump. The toolkit is not redundant in its implementation of interlocks, and is not meant to be used to protect operators. Software interlocks should be used only as a level of redundancy for existing hardware interlocks in the context of protecting operators. Several high-voltage bias supplies at MIT are hardware-interlocked by direct electronic circuitry to either access panels or pressure gauges. This is to ensure that they are powered off when human contact is physically possible. The toolkit’s software interlocks are used to provide a parallel path of redundancy by monitoring the pressure readout from the gauge and independently controlling the power supply. Thus, both the hardware and software interlock need to be functioning properly to allow operation of the bias supply.

An additional feature common to control systems is a multithreaded code; this allows continuous communication with hardware to continue in the background while retaining a responsive user-interface (UI). Finally, since most machines or experiments evolve over time, parts of the control software evolve with the hardware while other parts are reused in subsequent versions of the experiment. This naturally leads to a modular design where each module presents both an interface to the hardware and to the operator via the toolkit, as shown in Fig. B.1; in this way modules may be re-used, removed, or evolved as the experiment changes. Note that while the toolkit does not operate in real-time, it is intended to interface to and monitor real-time hardware controllers. It is these real-time controllers which then drive and monitor actual hardware. The MCT is thus a means for centralized monitoring and operation of an experimental apparatus.

Given the modular design, it is also often desirable to share data between modules and potentially between threads. This requires mutual-exclusion locking of the data structures (shared data is limited to RAM at the present). The motivation for shared data may not be entirely obvious; certain tasks, such as logging of pertinent parameters (e.g., run-time performance metrics associated with the experiment) are centralized functions and may even be implemented as modules. Such tasks require access to all variables of interest, which are likely be distributed among modules. As another example, consider the implementation of an interlock system, where one would need to “lock down” a system when a parameter of interest crosses a threshold value. For example, one
Figure B.1. Simplified architecture of control software written using the MCT, illustrating how real-time controllers (connected to transducers), user-developed code (herein referred to as the module) and the toolkit interact together. The module is a shared object loaded by the toolkit at run-time.

Figure B.2. Class diagram for the Modular Control Toolkit (MCT) using standard Unified Modeling Language (UML) notation. The core classes are shown on the left side: Console, IntLockMan, PrefMan and ModMan, which implement the console, the interlock manager, the preferences manager and the module manager, respectively. All user modules inherit from ModBase, which provides an interface for interlock and shared variable functions. For completeness, a shared library one might use for communicating with hardware devices are shown (i.e., an object of class type <CommClass> ). Instances of classes required for hardware communication are contained by the user module.
might want to lock down operation of a power supply when a pressure reading elsewhere in the system crosses a threshold value. This requires access to both the variable holding the pressure, which is constantly updated by one module and to the interlock system of another module.

The toolkit was designed to meet all of these specifications. The architecture is illustrated in the class diagram shown in Fig. B.2 using standard Unified Modeling Language (UML) notation. Note that a significant number of auxiliary operations and attributes are suppressed in this model so as to focus on the core architecture and functionality. The entry point for the toolkit is the `Console::run()` function, which is called after a Console object is instantiated. As shown in the UML model, this Console object is a container for exactly one instance of each of the `PrefMan`, `IntLockMan` and `ModMan` classes; these objects are responsible for managing toolkit preferences, interlocks and modules, respectively. Upon execution, the console will create an instance of each of these classes, initialize the multithreading engine and then initialize each of the three aforementioned objects, at which point it will execute the GTK event loop and wait for user input.

All user-developed modules are required to inherit from the class `ModBase`, as shown in Fig. B.2. This base class provides an interface to each module for registering and unregistering interlocks and global variables, for setting and retrieving the value of global variables, for posting messages to a centralized console and for reading and writing configuration data associated with the module. The toolkit facilitates the storage and retrieval of preferences associated with any given module by handling reading, writing and parsing of configuration data to and from disk. It is left up to the module designer to use the toolkit’s interface functions to perform these tasks. Note that interlock system definitions and global variables are stored in the interlock manager and the console, respectively. Thus, the base class must access these two objects to perform the aforementioned tasks on behalf of each module. The inheritance and class permissions are set to allow the base class `ModBase` access while shielding the user-defined module class from both the implementation details of and access to the rest of the toolkit. Within the toolkit, the class `ModBase` is packaged as a shared library, which has two main advantages: (1) Since the code contained within the base class is shared between all modules, linking at run-time to a common library reduces the executable size and (2) Internal changes may be made to this base class, as optimization and enhancements are implemented in future versions of the toolkit, without the need to re-build any of the modules; this greatly enhances maintainability of the toolkit.

Since modules themselves are compiled into shared libraries and loaded at a user’s request, an additional requirement is that they must implement a pre-defined interface. This allows the module manager to properly load, unload and query each module. For ease of development, a template including a skeleton module is provided in the toolkit source package which implements these interface symbols; it also sets up an environment for properly compiling module source code into a shared library.

Next, note the function `ModBase::thread_body()`. This function may be overridden by the module and in that case should contain user-defined code to be executed within a dedicated thread; it is called repeatedly within the body of a loop. Alternatively, the user may wish to setup their own thread and ignore these functions; the toolkit does not preclude one from doing so. In any case, two additional functions, `ModBase::initialize()` and `ModBase::terminate()`, may be overridden to perform module (de-)initialization (after) before the dedicated loop (stops) starts. The threading framework available to toolkit users does not incorporate any kind of supervisory thread prioritization when scheduling threads. However, since the MCT is built using GLib, developers may use library functions to yield or to prioritize thread execution as deemed necessary.

Finally, consider the interlock manager, shown in the screenshots of the user-interface in Fig. B.3. The operator may define rules at run-time which are checked by the interlock engine. Shown in the figure is an example of a rule for locking down a turbopump (i.e., the interlock system). In
Figure B.3. Screen capture of the MCT, showing three windows (as labeled by window titles): the Console, the Interlock Manager and the Module Manager. Important information is presented to the operator in the Console window; modules may be loaded, unloaded and configured from within the Module Manager; interlock rules may be defined, grouped and (de-)activated from within the Interlock Manager interface.

For the interlock system to be active, the monitor variable “ACC_JG1” (which is a standard toolkit global variable registered by a module) must be greater than $1 \times 10^{-2}$ Torr. What actually happens when the interlock system is engaged is left up to the owning module, in this case, “leyboldtd20ctrl.so.” The interlock manager then does several things: (1) provides a means for user-defined modules which inherit from ModBase to register and check the status of interlocks (2) a user-interface for the operator to define rules for each interlock system (several may be defined for each system) and (3) an engine to check operator-defined rules, which runs in a dedicated thread in the background. In this way, modules do not need to know how they fit into the bigger picture of the experimental apparatus; they simply register and update measured quantities of interest and register interlock systems for the hardware device they are operating. How the various devices are integrated together is left up to the operator at run-time, making modules more generic and re-usable.

The activity diagram in Fig. B.4 best illustrates the operation of the interlock engine; the sequence of events outlined in the figure take place within a dedicated thread. After initialization, the engine traverses a std::map of rules. A mutual-exclusion lock is obtained on this data structure as it is traversed. For each rule, the engine checks to ensure that the monitor variable and interlock system for that rule are available (that the module(s) which registered them are loaded). If they are not available, the engine marks the rule as inactive and moves onto the next rule; if both the monitor variable and interlock system become available at a later time, the engine will mark the rule as valid and proceed. Note that to check for these two conditions, the engine must obtain read-only locks on both the interlock systems and monitor (global) variable data structures to prevent any of the module threads or the main thread from attempting to modify these data structures;
mutual-exclusion and read/write locks are implemented using standard GLib methods. Checking the validity of rules in this manner allows user-defined interlock rules to remain safely defined within the MCT as modules are loaded and unloaded.

Once a rule has been deemed valid, the engine checks to ensure that the interlock system associated with that rule has not already been flagged for lock by a rule that was previously tested positive. If the system is already flagged, the rule is ignored since the system will lock regardless; otherwise the rule is tested. If the rule is tested and found positive, the interlock system associated with the rule is then flagged for lock. Subsequent rules associated with the same interlock system are not tested, though each rule is nevertheless traversed and checked for validity. Finally, after all rules have been traversed, the lock on the rules data structure is released. The flags on all of the interlock systems are checked and systems are locked or unlocked accordingly; these flags are reset for the next iteration of the engine loop. Note that since the engine runs within a dedicated thread, it sleeps for some period of time within each iteration so to not saturate or overwhelm the hardware.

As shown here, the implementation of interlocks and global variables requires the locking of

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**Figure B.4.** Activity diagram of the interlock engine, illustrating how rules are validated and tested and how associated interlock systems are locked and released. The sequence of events depicted execute within a dedicated thread.
multiple toolkit resources from within different threads. The design of the interlock engine and global variable system ensures that deadlock does not occur. Two design paradigms were followed to ensure this: (a) all resource locking for these systems is performed internally by toolkit functions (which are called by a module) and (b) no two toolkit functions that lock the same resource ever wait for each other to complete before releasing that resource. Since locking is performed from within toolkit functions, modules do not have direct access to these resources and therefore cannot lockout a resource directly.

### B.3 Design of Modules

The design of data-flows for modules are largely left up the developer. That is, module developers may take an event-driven, data-driven or mixed approach depending on the application. For instance, certain types of equipment may wait for and respond to user interaction (event-driven), while in other cases one might wait for a device to periodically send data and then either perform processing on that data or alert the user as data comes in (data-driven). Event-driven and data-driven approaches are readily accomplished with the use of event-handlers and a dedicated thread within the module, respectively. Alternatively, and often, one might use a combination of both techniques. For example, a simple pressure gauge controller for a system might require continuous polling and readout of the pressure (data-driven) as well as some basic control (user or event-driven) for configuration and operation of the gauges. The toolkit facilitates the implementation of any combination of these models.

As an example of the flexibility allowed by the toolkit in conjunction with GTK and GLib in implementing a mixed data- and event-driven module, consider a practical module which interfaces to a vacuum pump controller. The vacuum pump itself is driven by a hardware controller with a remote serial interface; the manufacturer has provided a set of commands to retrieve information about the status of the pump, as well as commands to start and stop the pump. The module must present a user interface to the operator, displaying pump parameters such as load and temperature and allow for control of the pump, so the user can start and stop it. To achieve the first goal, it must interface with the pump and retrieve its status on a periodic basis; this is best accomplished in the background using a dedicated thread. Since the module must periodically communicate with the instrument in a separate thread and update the user-interface in the parent thread (a GTK requirement that GUI calls be made in the parent thread) it will have to dispatch function calls for updating the user-interface to the parent thread. Furthermore, the mixed event-driven and data-driven approach dictates that hardware access will occur from two threads: (1) the dedicated thread which will periodically poll the pump controller and retrieve parameters and (2) the parent thread, which handles the user-interface event-handlers, where function calls to start and stop the pump will occur. This naturally leads to the requirement of a mutual exclusion lock for hardware access. All of these constraints are readily handled using GTK and GLib. The module implementing this functionality would inherit from the class ModBase and re-implement functions as follows:

```cpp
class PumpControl : public ModBase
{
    public:
    PumpControl();
    ~PumpControl();

    void initialize();
    void terminate();
```
protected:
    void thread_body();
    GLib::Dispatcher update_gui();
    GLib::Mutex serial_mutex;
    void update_gui_disp();
    double pump_temperature;
    double pump_load;

serial pump_connection;
};

The GLib::Dispatcher object is used to dispatch the code contained in the function update_gui_disp() to the parent thread; the connection between this object and the dispatched function is made in the module’s initialize() function. The GLib::Mutex object is used to lock hardware access so that only one thread may communicate with the pump controller at a time. Finally, the serial object is an instance of a library class which allows communication over serial ports. The functions shown above may be re-implemented as follows:

void PumpControl::initialize()
{
    // Setup GUI
    // Connect dispatcher
    update_gui.connect(sigc::mem_func(*this,&PumpControl::update_gui_disp);
    // Add event handler for stopping pump
}

void PumpControl::thread_body()
{
    sleep(POLLING_PERIOD);
    // Acquire hardware access lock
    GLib::Mutex::Lock serial_access(serial_mutex);
    // Communicate with instrument, store temperature and load values in data members
    pump_connection.read(...);
    // Release hardware lock
    serial_access.release();
    // Call update_gui_disp() in main thread
    update_gui();
}
void PumpControl::update_gui_disp()
{
  // Use data members to update GUI here
}

void PumpControl::on_stop_pump()
{
  // Acquire hardware access lock
  GLib::Mutex::Lock serial_access(serial_mutex);
  // Communicate with instrument
  pump_connection.write(...);
  // Release hardware lock
  serial_access.release()
}

This bare-bones sample module allows for continuous polling of the pump while retaining a responsive control GUI. There are, however, some subtleties. For instance, the creation of a GLib::Mutex::Lock object on line 35 (from within on_stop_pump()) is a blocking call, which will wait for a call to release() (from within the function thread_body executing in a dedicated thread) before continuing. One must therefore be careful not to make the polling period or communication time per iteration too long for this will cause the GUI to become less responsive. This is seldom an issue in practice because the time-scales considered here not sufficiently long for typical communication schemes with modern laboratory instruments, even over slow (9600 baud) serial devices. Note that these kinds of timing considerations are not exclusive to the MCT, but are common to all the previously mentioned toolkits; these level of detailed considerations must be left up to the developer since these toolkits know nothing about the types of instruments being interfaced. Note that the sample code above does not register interlocks or global variables. These may be registered, unregistered and read from within any thread (this is properly handled by the MCT), giving the module designer much flexibility. Configuration data, such as the serial port parameters to use when connecting to the pump controller may also be saved and read using the toolkit. The sample presented here is simplified to demonstrate the flexibility of the MCT in its ability to accommodate various programming models. Complete modules, written in the course of developing this toolkit, included configuration dialogs for instrument setup, communications error-checking, advanced GUIs, checking and handling of interlocks, etc.

B.4 Development and Initial Deployment

The MCT was developed with several modules to serve the needs of the MIT LEIA Facility. LEIA is an accelerator comprised of several vacuum pumps and valves, numerous high-voltage bias supplies, and many pressure transducers. These components interface to real-time controllers that are interconnected with a control computer and a data acquisition computer using a fiber-optic network. In some cases, a single real-time controller drives several components simultaneously. The control computer drives all of these real-time controllers: pump controllers (2), ion source controller (1), valve controller (1) and pressure gauge controllers (2). Communication with these controllers is implemented using four modules, each with a dedicated thread. Together with the main thread and interlock engine, a total of six hardware threads are used for normal operation. The toolkit
B.5 Summary

A modular toolkit for control of scientific experiments and instruments has been developed. The toolkit allows novice developers to quickly build a multithreaded modular control application without simultaneously limiting advanced users. Users have the option and not the obligation of using the threading framework, shared variables and centralized storage of configuration data. Modules may rely heavily on these various aspects of the toolkit or implement these features independently. Since the framework is open-source and written in C++, advanced developers have the ability to use development tools of their choice to customize and improve the inner workings of the toolkit and to commit any improvements back to the user community.

Future versions of the toolkit will incorporate a number of improvements, including:

1. Full use of C++ namespaces to mitigate any ambiguities in user-developed module code.

2. An integrated diagnostic tool to allow users to monitor the number of running threads, reg-
istered interlock systems, global variables and to monitor system resource usage.

3. Ability for modules to register callback functions with the toolkit (event handlers) for handling toolkit events such as the registration of global variables and interlock systems. This is useful to modules which take all global variable data and log it to a database or to disk.

4. Priority scheduling of threads that lock toolkit resources (indirectly by calling toolkit functions); this will improve application performance as the number of modules accessing toolkit resources increases.

5. Execution of modules within a dedicated process (similar to a “sandbox”) for enhanced stability, robustness and recovery from localized data corruption and run-time errors.

The MCT has proved to be a robust control solution at the MIT LEIA Facility. It serves the needs of small to medium scale experiments and facilities, defined here as a system comprised of tens of hardware components driven by one or two computers. Relative to commercial control software or other open-source alternatives, it is cost-effective and ideal for both novice and advanced users. In addition to being a toolkit, it provides a re-usable application framework, allowing novice users to combine it with existing modules and use it out-of-the-box; advanced users are free to develop modules with little restriction.

B.6 References

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5. Data Acquisition Systems Laboratory, see http://www.dasylab.com.
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Appendix C

The Response of CR-39 Nuclear Track Detector to 1-9 MeV Protons

C.1 Introduction

The detection of charged particles (and neutrons) using solid-state nuclear track detectors is of interest to many scientific disciplines, including the field of Inertial Confinement Fusion (ICF) research; this includes both laser-plasma experiments\(^1\)\(^-\)\(^4\) as well as Z-pinch experiments. In this context, CR-39 offers several advantages over alternative charged-particle detectors. As a passive, plastic detector, it is immune to electromagnetic pulse (EMP) and largely immune to x-rays which are commonly produced in such environments. For particles with normal incidence to the detector surface, the material has 100% detection efficiency for protons up to approximately 6-8 MeV. As will be described, this is accomplished using standard optical microscope techniques to image and analyze the detector.\(^5\)

CR-39 has been used in wedge-range-filter and magnetic charged-particle spectrometers,\(^5\) as well as neutron spectrometers using the recoil technique\(^6\) at the OMEGA laser facility, and recently at the NIF for diagnosing neutrons and charged-particles from ICF implosions. CR-39 has also been fielded in a diagnostic for imaging ICF implosions,\(^7\)\(^,\)\(^8\) which has allowed the detailed study of field structures in those experiments.\(^9\)\(^,\)\(^10\) In these diagnostics, the track diameters are often used to discern different charged-particle species, and in several cases are used to determine the energy of the incident particle. A thorough understanding of the response of CR-39 to charged-particles is often important and sometimes essential for the success of such diagnostic instrumentation.

Previous work, utilizing various types of CR-39, have studied the effectiveness of various track etchants,\(^11\) and have looked at the CR-39 response to alpha particles,\(^12\)\(^-\)\(^14\) electrons,\(^15\) gammas,\(^16\) fission fragments,\(^17\) and neutrons.\(^18\) Studies of the CR-39 efficiency for detecting protons as a function of incident angle in the energy range of 1-10 MeV have also been undertaken, as well as preliminary studies which characterized the track diameter as a function of incident mean proton energy.\(^17\) In this work, we present the TasTrak CR-39 track diameter response curves to protons as a function of energy, including a study of piece-to-piece variability in addition to the effects of etch time and etchant temperature. We also present the results of a 5-year study of the effects of aging of CR-39 using 5.5 MeV alpha particles.

A charged particle leaves a track of broken molecular chains and free radicals along its trajectory as it passes through the plastic. Particle tracks are then made visible by chemical etching in hot, concentrated alkali solutions. The average diameter of the tracks left on the surface of the plastic will depend on the incident mean energy of the particles as well as how the CR-39 is processed. If the response of CR-39 to protons is understood for a range of etchant properties, we can obtain
spectral information about the particles incident on the surface of the plastic.

This appendix is structured as follows: Section C.2 describes the principal experimental setup for these measurements; Sec. C.2 presents the results of three different studies on CR-39’s track diameter response, including piece-to-piece variability, effects of different etching conditions and age of CR-39; Sec. C.3.5 illustrates how signal tracks are discerned from intrinsic noise for the different studies presented in this paper, and Sec. C.4 summarizes the study presented in this appendix.

### C.2 Experiments

The studies presented in the following section involved exposing pieces of CR-39 to protons in the energy range of 0.92 MeV to 9.28 MeV, which were produced using a linear accelerator-based fusion products generator. A deuterium (D) ion beam was focused onto an erbium (Er) target embedded with either D or Helium-3 (³He) atoms to generate DD or DHe fusion reactions. The 3 MeV DD-protons and 14.7-MeV DHe-protons were used for the CR-39 exposure. Protons of the energy range indicated above were generated by overlaying the CR-39 with aluminum filters of various thicknesses. The remaining charged fusion products (alphas and tritons) were readily ranged out with aluminum filters; the CR-39 response to the 2.45 MeV DD-neutron products is not of concern in this study since it has been demonstrated previously that the detection efficiency for these neutrons is approximately $(1.1 \pm 0.2) \times 10^{-4}$ on the front side of the CR-39 for a 6 hour etch, and lower for shorter etch times.

Measurements of the proton energy spectrum at birth as well as behind each individual aluminum filter were acquired using a silicon surface barrier detector (SBD) with a nominal depletion depth of 2000 $\mu$m, which is thick enough to completely stop 14.7-MeV protons. Two filter arrangements were constructed to range down the protons to the desired energies. The energy distribution of the protons leaving each of these filter packs was measured using an SBD (Fig. C.1). Once the CR-39 was exposed to protons, it was etched using a heated solution of NaOH and then scanned with an automated optical microscope system, which recorded the location, diameter, eccentricity, and optical contrast of each track. Following the scan, track attributes were analyzed using a custom software package developed at MIT.

### C.3 Results

#### C.3.1 Track diameter as a function of proton energy

The proton track diameter response of 1500-$\mu$m-thick CR-39 has been studied. Following exposure to protons, a 6 N solution of NaOH at a temperature of 80°C was used to etch the CR-39 for a duration of 6 hours. A 6 N solution of NaOH is typically used as it has been shown that the sensitivity of CR-39, defined here as the ratio between the etch rate of a track and that of the bulk $(V_T/V_B)$, is maximized at this molarity; higher sensitivity improves track contrast and the detection of charged particles with an optical microscope.

The diameter of proton tracks as a function of incident mean proton energy is shown in Fig. C.2. The curves were acquired using two separate accelerator exposures, utilizing the two filter packs characterized in Figs. C.1a and C.1b, on a single piece of CR-39, simply to avoid any piece-to-piece variation. The corresponding diameter distributions of protons behind each of the filters as recorded by the CR-39 are included for reference in Figs. C.1c and C.1d. During exposure to 14.7 MeV protons, the upper portion of the CR-39, containing the thinner D-D aluminum filter
Figure C.1. Energy distribution function of protons used in this study. Protons were generated from 3 MeV DD or 14.7 MeV D$_3$He fusion products with the use of stepped aluminum filters, which were overlaid on the CR-39 sample; filter thicknesses are indicated above respective peaks. Energy distributions were measured with a detector system utilizing a 2000 μm thick silicon surface barrier detector (SBD); the energy accuracy of the system was ±75 keV.
Appendix C The Response of CR-39 Nuclear Track Detector to 1-9 MeV Protons

Figure C.2. (a) Proton track diameter as a function of incident mean energy for TasTrak CR-39 from 11 separate exposures; each piece was etched for 6 hours in an 80°C solution of 6 N NaOH. Variability in the shape of the curves from one piece to the next is observed. (b) Average (solid line) and standard deviation (dashed line) of proton track diameters as a function of incident mean energy computed using results from the 11 separate exposures of Fig.C.2a.

pack, was blocked from the more energetic D-3He protons using a vacuum shutter; this eliminated the need to vent the CR-39 to atmosphere and reconfigure filters, ensuring that the experiment was not perturbed between the two exposures. Eleven such exposures (Fig.C.2a) were then made to characterize the piece-to-piece variability in the response. The average and standard deviation (dashed envelope) of these 11 exposures is shown in figure Fig. C.2b. It is important to note that the shape of each individual curve varies (see Fig. C.2a); to use these curves quantitatively one must have knowledge of the shape of the curve for a given piece. It is thought that this piece-to-piece variation is a result of the variable oxygen profile of CR-39 upon irradiation.

As illustrated in Fig. C.2, the track diameters decrease with increasing incident mean proton energy, as expected from the stopping power scaling (dE/dx ∼ 1/E) of charged particles in cold matter. The features of these curves are in qualitative agreement with a previous study utilizing 1-10 MeV protons, although that study used a different type of CR-39 (Homalite) as well as a different etchant (6 N at 70°C). In addition to these observed proton tracks, noise caused by defects and debris in the CR-39 are often observed. The diameters of these features, however, for the standard etch described above, are less than 2 μm and will not be counted as proton tracks for energies below 8 MeV.

C.3.2 Effects of Etch Time on Track Diameter

The effects of etch time were studied using four pieces of CR-39, which were exposed to protons of various energies using the aforementioned procedure. Each piece was then etched using a 6 N solution of NaOH at a temperature of 80°C and scanned several times, starting with a 30 minute etch and increasing in increments up to a total 6 hours. This procedure was repeated for a total of four pieces and the results averaged to eliminate piece-to-piece variation; the resulting diameter response curves are shown in Fig. C.3.

In several charged-particle diagnostic ICF applications, the CR-39 in use often exhibits significant variations in the proton track density. In such cases, it is desirable to apply a staged etch and scan process to the CR-39, using scans from short and long etch times for regions of high and low fluence, respectively; for regions of high fluence this avoids track overlap and in severe
Figure C.3. Proton track diameters as a function of incident mean energy for various etch times. Four samples of CR-39 were stage-etched in steps from 30 minutes up to 6 hours in a 6 N NaOH solution at 80°C; the results between the four pieces have been averaged to eliminate any piece-to-piece variation. A linear response between etch time and track diameter is demonstrated. Note that for a 30 minute etch, protons with energies greater than 3 MeV are not detected and hence not shown.

C.3.3 Effects of Etch Temperature on Track Diameter

Previous work involved studies of the effects of etchant temperature variation on the CR-39 track diameter response to 1 MeV protons as well as fission fragments and alpha particles. These studies have been extended to include the track diameter response curves of protons in the energy range 1-9 MeV for various etchant temperatures (Fig. C.4). The CR-39 for this study was etched for 6 hours in a 6 N solution of NaOH at temperatures of 75°C, 80°C, and 85°C.

Figure C.4 shows that proton track diameters grow nearly linearly for these range of temperatures. Although higher etchant temperatures are clearly desirable from a processing standpoint (due to the reduced processing time) one must consider the integrity of the charged-particle tracks. Typical tracks of 1 MeV protons and 5.5 MeV alphas for these etchant temperatures, taken with an optical microscope (Fig. C.5), show the loss of uniformity in track contrast over the surface of the track. It is unclear whether the loss of uniformity is due to temperature effects alone; this is also observed when etching CR-39 for longer times, due to etching beyond the end of the track.

Figure C.4. Proton track diameter as a function of incident mean energy for various etchant temperatures. Three separate pieces of CR-39 were exposed to protons and etched in a 6 N NaOH solution for 6 hours at different etchant temperatures. Two exposures per piece were required to obtain each curve, as indicated by the accelerator shot numbers in the legend.
Figure C.5. Microscope images (at 40× magnification) of (a) 1 MeV proton and (b) 5.5 MeV alpha particle tracks on CR-39, illustrating the non-uniformity of track contrast with increasing etchant temperature. Images were taken from three samples, each containing protons and alphas.

C.3.4 Effects of aging and time of exposure

It has been shown previously that the CR-39 response to charged-particles drops sharply within the first 30 days of manufacture, and that it is the age of the plastic itself, not of the tracks, which results in the degradation of sensitivity (lower $V_T/V_B$) with time. The effects of prolonged aging on the CR-39 were studied using 5.5 MeV alpha particles from a $^{241}$Am source and samples over a 5-year period. After exposure, samples were stored in a low-humidity, low-light environment at room temperature for a variable amount of time, after which they were etched using the standard 6 N solution of NaOH at 80°C for 6 hours. The diameters of the alpha tracks were found to decrease as a function of time between exposure and etch, as shown in Fig. C.6a. To determine whether it is the age of the plastic or the age of the track which causes a reduction in track diameter, another set of samples which were previously exposed to alpha particles were exposed again immediately before etch. The ratios of diameters of the old alpha tracks to the new alpha tracks is unity to within 5%, suggesting that it is the age of the CR-39, not the age of the track, which is responsible for the reduction in diameter, in agreement with previous claims utilizing different materials. The cause for the observed change in response may be annealing or oxidation. Given the experimental setup, we rule out long-term exposure to UV and humidity. It has been shown that storing CR-39 at or below freezing temperatures will inhibit these aging effects.

Figure C.6. (a) Track diameter of alpha particles as a function of time between exposure of CR-39 pieces to 5.5 MeV alphas and etching. Each of these pieces was kept at room temperature with minimal exposure to light for the time periods shown. (b) The effects of aging on alpha-track diameter are independent on the time of exposure. CR-39, which had been previously exposed to alpha particles, was exposed again to alpha particles just prior to etching. The differences between new and old track diameters are $\sim 5\%$. 
C.3.5 Variation of Track Contrast

The analysis of the CR-39 depends on careful discrimination of tracks based on a number of properties, including diameter, contrast, and eccentricity. The selection of signal tracks based on these properties is important in preferentially discriminating against noise. This section illustrates how the contours of constant numbers of signal tracks as a function of both diameter and contrast vary for the studies presented in this appendix. Also included is a summary table of the accelerator shot numbers corresponding to the experiments in this study.

Figure C.7 illustrates how track diameters of protons behave for proton energies near the Bragg peak of energy deposition in CR-39. The 920 keV protons in this case occupy the top-right corner of contrast vs. diameter space. Since protons with these energies are right near the Bragg peak for energy loss in CR-39, the diameters are no longer monotonic with energy and hence wrap around back to lower diameters. This results in the non-Gaussian diameter distribution shown in Fig. C.1c.

Figure C.8 shows the variation of track contrast and diameter distributions with etch time for 3.6 MeV protons; these contrast and diameter distributions correspond to the data shown in

Figure C.7. Contours of constant number of tracks as a function of track contrast and diameter for DD-p fusion products ranged down to 920 keV. The proton tracks are in the top right corner of the plot with intrinsic noise occupying regions of lower contrast and smaller diameter. Note how the diameter distribution turns around for these protons, resulting in the non-Gaussian diameter distribution of Fig. C.1c.

Figure C.8. Contours of constant number of proton tracks as a function of track contrast and diameter for a single piece of CR-39 etched in stages. Shown are data for 3.6 MeV protons generated by ranging down D-3He fusion protons with the use of aluminum filters. As the etch time is increased, the proton tracks shift towards higher contrast and larger diameters, making them more distinct from intrinsic noise.
Appendix C The Response of CR-39 Nuclear Track Detector to 1-9 MeV Protons

Figure C.9. Contours of constant number of tracks as a function of track contrast and diameter for a 6 hour etch at 75°C, 80°C and 85°C of 6 N NaOH. Two distinct islands of contours are visible in each plot: noise which occupies the bottom-left corner of contrast-diameter space and proton tracks which occupy the upper region and shift to the right for increasing etch temperatures.

Fig. C.3. The low contrast, small diameter contours in the lower left of each plot show the intrinsic noise on the sample. As the etch time grows, the contours representing protons move towards larger diameters and higher contrast, making them easier to discern from the intrinsic background.

Shown in Figure C.9 are contours illustrating the diameter and contrast distributions for 3.6 MeV protons as a function of etchant temperature. Increasing the etchant temperature shows effects similar to that of increasing etch time; proton tracks tend to become darker and larger since the etch rate is accelerated. There is some variation of intrinsic noise between the data shown for three etchant temperatures. This illustrates the variation of noise from one piece to the next since three distinct samples were used to obtain the data.

C.4 Summary

The response of TasTrak CR-39 to 1-9 MeV protons has been studied, including the piece-to-piece variability of the response and the effects of different etch times and etchant temperatures. It has been shown that the shape of the proton track diameter vs. energy response curve varies from one piece to the next, and that quantitative use of these curves requires one to have knowledge of the shape of the curve for a given piece. Furthermore, a linear relationship between etch time and track diameter is observed, allowing for staged-etch processing of CR-39. Effects due to the age of CR-39 have also been studied over a 5-year period and from these experiments it was found that 5.5 MeV alpha particle track diameters decrease as a function of age of the plastic itself and not the age of the track. These characterizations of the response of CR-39 to protons are essential for the calibration of existing diagnostics and for the development of new diagnostic capabilities for the OMEGA and NIF laser facilities.

Future work will include the study of various environmental effects important to the ICF community. The effect on track diameters of prolonged exposure of CR-39 to vacuum both before and after exposure to protons has proven to be important at large facilities, such as the NIF where CR-39 may sit in vacuum for several hours if not days before the experiment is begun. The efficiency of track detection for shorter etch times and higher count rates (where track overlap becomes significant) is also of interest for improved processing turnaround time and increased flexibility in experiment designs, respectively. Finally, a detailed study of the x-ray response of CR-39 and in particular the effect of x-rays on charged-particle tracks will become important in the harsh x-ray environments of todays and tomorrows increasingly more intense laser systems.
C.5 References


Appendix D

Derivation of Thomson Parabola Scaling Laws

D.1 Introduction

In this appendix, scaling laws for Thomson parabola diagnostics are presented, including the deflection of charged-particles in uniform magnetic and electric fields and energy resolution. These equations are useful for the conceptual design and operation of the Thomson Parabola Ion Spectrometer (TPIS) and the Thomson Parabola Ion Energy (TPIE) Analyzer. If one seeks to make a measurement with a certain resolution in a certain energy range, the following equations may be used to estimate the required field strengths and aperture required to successfully achieve the desired measurement.

D.2 Magnetic Deflection

The derivation begins by considering parallel and uniform magnetic and electric fields, a detector plane and an aperture as shown in Fig. D.1. Consider the deflection of a particle with a positive charge, mass, and energy given by $q$, $m$ and $E_p$, respectively. As the particle enters the magnetic field, it will gyrate due to $v \times B$ motion:

$$x(t) = r_{gyro} \sin(\omega_c t), \quad (D.1)$$

where the gyroradius is $r_{gyro} \equiv v_0/\omega_c$, the cyclotron frequency is $\omega_c \equiv qB/m$ and the initial particle velocity is $v_0 \equiv \sqrt{2E_p/m}$. The Larmor motion ceases when the particle leaves the magnetic field at $x = l_b$. The time it takes to traverse the face of the magnet is given by:

$$t_{gyro} = \frac{1}{\omega_c} \arcsin\left(\frac{l_b \omega_c}{v_0}\right). \quad (D.2)$$

The magnetic deflection at this point is:

$$x_1 = \frac{v_0}{\omega_c} \left[1 - \cos(\omega_c t_{gyro})\right] \approx \frac{v_0}{\omega_c} \left[1 - \sqrt{1 - \frac{l_b^2 \omega_c^2}{v_0^2}}\right] \approx \frac{v_0}{\omega_c} \left[\frac{l_b^2 \omega_c^2}{2v_0^2}\right], \quad (D.3)$$

where the last equality is an expansion in the limit of $\delta^2 \equiv l_b^2 \omega_c^2/v_0^2 \ll 1$. Physically, this corresponds to a deflection in the uniform field region which is small compared to the Larmor radius (e.g., for a 500 keV proton in the TPIS, $\delta = 0.08$). Next, we must find the deflection as the particles freely drift towards the detector. The particle velocities upon exit of the uniform $B$ region are:
Figure D.1. Schematic of the Thomson parabola geometry as applicable to both the TPIS and the TPIE Analyzer. Shown are the aperture, the permanent magnet, the electrodes and the CR-39/Image plate detector. The relevant dimensions used in the derivations throughout this appendix are labeled for reference.

\[ v_{x1} = v_0 \sin(\omega_c t_{gyro}) = l_b \omega_c \]  
\[ v_{y1} = v_0 \cos(\omega_c t_{gyro}) \approx v_0 \left( 1 - \frac{l_b^2 \omega_c^2}{2v_0^2} \right), \]  

where the same limits (\( \delta \ll 1 \)) have been once again applied. The drift time in the \( y \)-direction is:

\[ t_{bdrift} = \frac{l_{bdrift}}{v_{y1}} \]  

and the amount of deflection in the \( x \)-direction during this time is given by:

\[ x_2 = v_{x1} t_{bdrift} = \frac{l_b l_{bdrift} \omega_c}{v_0} \left[ 1 + \frac{l_b^3 \omega_c^3}{2v_0^2} \right]. \]  

The total magnetic deflection is then expressed as

\[ x = x_1 + x_2 = \frac{l_b \omega_c}{v_0} \left[ \frac{l_b}{2} + l_{bdrift} \right] + \frac{l_b^3 \omega_c^3}{v_0^2} \left[ \frac{l_{bdrift}}{2} \right]. \]  

Note that the second term in the above equation is the first correction in \( \delta \). It may be often neglected since it is of order \( \delta^3 \). The correction signifies the finite deflection of the particle as it exits the uniform field region. In typical Thomson parabola designs the majority of magnetic deflection comes from the small but finite \( v_{x1} \) imparted on the particle as it exits the field and subsequently drifts for some time.

D.3 Electric Deflection

The amount of electric deflection may be analytically determined independent of the preceding magnetic deflection. Depending on the position and angle of a particle entering in the electrodes (as determined by the preceding B-field) the detailed fringe-fields seen by particles may differ. This effect is however negligible, as confirmed by simulations (see Secs. 3.3.2 and 3.3.3). The electric
deflection is derived as follows. Consider a uniform electric field with strength \( E \) and electrode length \( l_e \). The amount of time spent in the electrode region is

\[
t_e = \frac{l_e}{v_0} = \frac{l_e}{v_0} \left[1 - \frac{l_b^2 \omega_c^2}{v_0^2}\right]^{-1/2} \approx \frac{l_e}{v_0} \left[1 + \frac{l_b^2 \omega_c^2}{2v_0^2}\right], \tag{D.9}
\]

where the same expansion in \( \delta \) has been made. The amount of electric deflection during this time is just

\[
z_1 = qE \frac{t_e^2}{2m} \approx qE \frac{l_e^2}{2m v_0^2} \left[1 + \frac{l_b^2 \omega_c^2}{2v_0^2}\right] \tag{D.10}
\]

and the velocity of the particle in the \( z \)-direction as it exits the electrodes is

\[v_{z1} = qE \frac{t_e}{m}. \tag{D.11}\]

The particle then drifts freely over a region \( l_{edrift} \) until it reaches the detector plane. The drift time is

\[
t_{edrift} = \frac{l_{edrift}}{v_0} = \frac{l_{edrift}}{v_0} \left[1 - \frac{l_b^2 \omega_c^2}{v_0^2}\right]^{-1/2} \approx \frac{l_{edrift}}{v_0} \left[1 + \frac{l_b^2 \omega_c^2}{2v_0^2}\right]. \tag{D.12}
\]

The drift velocity in the \( y \)-direction is assumed to be constant because the fringe fields will have a minimal impact on the drift velocity in this direction. The amount of electric deflection in the \( z \)-direction during the drift region is

\[
z_2 = v_{z1} t_{edrift} = \frac{qE l_e l_{edrift}}{mv_0^2} \left[1 + \frac{l_b^2 \omega_c^2}{2v_0^2}\right] \approx \frac{qE l_e l_{edrift}}{m v_0^2} \left[1 + \frac{l_b^2 \omega_c^2}{2v_0^2}\right]. \tag{D.13}
\]

The total deflection due to the electric field is thus given by

\[
z = z_1 + z_2 = \frac{qE l_e}{mv_0^2} \left[\frac{l_e}{2} + l_{edrift}\right] \left[1 + \frac{l_b^2 \omega_c^2}{2v_0^2}\right]. \tag{D.14}
\]

Like the magnetic deflection equations, the last factor may be taken to be unity for typical Thomson parabola designs. Physically, this approximation signifies that the velocity reduction in the \( y \)-direction due to deflection in the magnetic field is negligible. Significant deflection in the magnetic fields converts \( v_x \) to \( v_y \) and causes particles to spend more time in the subsequent electric fields, enhancing the electric deflection.
D.4 Parabola Slopes and Energy Resolution

The slope of a given parabola ($\alpha$) is found by combining the equations for the magnetic and electric deflection. Combining Eqs. D.8 and D.14 and eliminating $E_p$ gives the following

$$ z = \alpha x^2 $$

with

$$ \alpha \equiv \frac{m}{q} \frac{2E}{l_e} \left[ \frac{l_e}{2} + l_{edrift} \right] \frac{B^2 l_b^2}{l_b^2 + l_{bdrift}^2}. $$

The slope of a parabola, a measured quantity, is thus inversely proportional to the charge-to-mass ratio of the ion. Thus, larger parabola slopes correspond to a lower charge-to-mass ratio. Furthermore, if the materials used in a given experiment are known (e.g., CH spherical shell), then that information together with the measured $\alpha$’s specifies the types and charge states of the measured particles.

Note that each position along a fixed parabola corresponds to a unique incident particle velocity (or energy, since the materials used in an experiment are known). The most energetic particles undergo less deflection and reach the detector near the origin while less energetic particles spend more time deflecting and thus reach the detector farther away from the origin, as shown in Fig. D.2. The relationship between position (along a parabola) and energy is a non-linear function. The energy resolution of a Thomson parabola is thus a non-linear function of energy. It is largely determined by both the optics of the diagnostic and by the dispersion due to the fields. In addition, focusing (defocusing) arising from field non-uniformities can improve (degrade) the resolution. This effect is not considered in these derivations for uniform fields. At any point along a parabola with a specified slope, the energy resolution is given by the product

$$ \Delta E_p = \frac{dE_p}{dl} \Delta b, $$

where $dE/dl$ is the variation of incident particle energy along the parabola and $\Delta b$ is the projection of the aperture onto the detector plane (see Fig. D.3). We first derive an expression for $\delta b$, and then follow with a derivation of the variation of incident particle energy along a parabola ($dE/dl$).

The physical picture for aperture projection onto the detector is shown in Fig. D.3. From the

![Figure D.3](image)

Figure D.3. Geometry showing the projection of the aperture onto the detector. The dashed lines indicate particle trajectories. The extended source size, aperture diameter, source to aperture distance and aperture to detector distance all play a role on the final projection. It is assumed here that the aperture thickness is negligible in comparison to the source to aperture distance ($D$).
geometry, we can immediately write down the following:

\[
\frac{a}{s} = \frac{D_2}{D_1} \quad \text{(D.18)}
\]

\[
\frac{b}{a} = \frac{d + D_2}{D_2} \quad \text{(D.19)}
\]

\[D_1 = D - D_2 \quad \text{(D.20)}
\]

We can combine these three equations, simultaneously eliminating \(D_1\) and \(D_2\), to arrive at an expression for the projection on the detector:

\[
\Delta b = a \left[ \left( 1 + \frac{d}{D} \left( 1 + \frac{s}{a} \right) \right) \right]. \quad \text{(D.21)}
\]

The projection depends on the aperture size \((a)\), the distance from the aperture to the detector \((d)\), the distance from the source to the aperture \((D)\), and the size of the source \((s)\). Note that finite source-size effects are generally small and can be neglected. Furthermore, the equation ignores the thickness of the aperture \((t)\). Finite thickness effectively increases the distance from the source and thus enhances the inferred yield by a factor \([\left( D + t \right)/D]^2\); for both the TPIS and TPIE, this effect is of order 0.1\% and therefore negligible.

Next, we derive an expression for the variation of incident particle energy along a parabola \((dE/dl)\). The variation of energy along a given parabola may be written as

\[
\frac{dE_p}{dl} = \frac{dE_p}{dx} \frac{dx}{dl}, \quad \text{(D.22)}
\]

A chain rule expansion of \(dE_p/dl\) in \(dE_p/dz\) instead of \(dE_p/dx\) is equally valid, since only one variable is required to specify the position along a given parabola. Next, note that

\[
\frac{dl}{dx} = \left[ 1 + \left( \frac{dz}{dx} \right)^2 \right]^{1/2}, \quad \text{(D.23)}
\]

where \(dz/dx\) is found by combining Eqs. D.8, D.15 and D.16:

\[
\frac{dz}{dx} = 2\alpha x = \frac{4E l_e}{(2E_p/m)^{1/2} B l_b} \left[ \frac{l}{2} + l_{edrift} \right]. \quad \text{(D.24)}
\]

Substitution of this equation into Eq. D.23 gives

\[
\frac{dl}{dx} = \left[ 1 + \left( \frac{4E l_e}{(2E_p/m)^{1/2} B l_b} \left[ \frac{l}{2} + l_{edrift} \right] \right)^2 \right]^{1/2}. \quad \text{(D.25)}
\]

Next, we invert Eq. D.8 to get energy as a function of magnetic deflection:

\[
E_p = \frac{1}{x^2} \left[ \frac{l}{2} + l_{bdrift} \right]^2 \quad \text{(D.26)}
\]

and take the variation along \(x\) to get \(dE/dx\):
\[ \frac{dE_p}{dx} = -\frac{2}{x^3} \left[ \frac{l_b q B}{(2m)^{1/2}} \left[ \frac{l_b}{2} + l_{bdrift} \right] \right]^2 . \] (D.27)

Substitution of this equation and Eq. D.23 into Eq. D.22 gives the total variation of energy along a parabola:

\[ \frac{dE}{dl} = -\frac{2}{x^3} \left[ \frac{l_b q B}{(2m)^{1/2}} \left[ \frac{l_b}{2} + l_{bdrift} \right] \right]^2 \left[ 1 + \left( \frac{4E l_e}{(2E_p/m)^{1/2} B l_b \left[ \frac{l_b}{2} + l_{bdrift} \right]} \right)^2 \right]^{-1/2} . \] (D.28)

The magnetic deflection \( x \) may be eliminated in the above equation by backsubstitution of Eq. D.8. The result is

\[ \frac{dE_p}{dl} = -\frac{4E_p^2}{q} \left[ \left( \frac{E_p}{2m} \right)^{1/2} l_b B \left[ l_b + 2 l_{bdrift} \right] \right]^2 + \left( 2E l_e \left[ l_e + 2 l_{edrift} \right] \right)^2 \right]^{-1/2} . \] (D.29)

Finally, by combining this equation with Eqs. D.17 and D.21, the energy resolution can be expressed as:

\[ \Delta E_p = -\frac{4a(1 + d/D)}{q} \left[ \left( \frac{E_p}{2m} \right)^{1/2} l_b B \left[ l_b + 2 l_{bdrift} \right] \right]^2 + \left( 2E l_e \left[ l_e + 2 l_{edrift} \right] \right)^2 \right]^{-1/2} . \] (D.30)

To first order, the energy resolution scales linearly with the square of the particle energy and inversely with the particle charge. The resolution also depends on the magnetic and electric field geometries and strengths. Depending on the design and of the Thomson parabola, the resolution may be dominated by either the electric or magnetic term.

For typical operating values of the TPIS and the TPIE Analyzer (e.g., \( B \sim 5\text{kG}, E \sim 5\text{kV~cm}^{-1} \)), the magnetic dispersion is significantly larger than the electric dispersion. Thus, both instruments operate in the regime:

\[ \left( \frac{E_p}{2m} \right)^{1/2} l_b B \left[ l_b + 2 l_{bdrift} \right] \gg \left( 2E l_e \left[ l_e + 2 l_{edrift} \right] \right)^2 . \] (D.31)

In this limit, the expression for energy resolution may be simplified by expanding Eq.D.30. The resulting expression for the magnetic-dominated regime is:

\[ \Delta E_p \approx -\frac{4a(1 + d/D)\sqrt{2}}{q} \frac{mE_p^{3/2}}{l_b B \left[ l_b + 2 l_{bdrift} \right]} \left[ 1 - \frac{m}{E_p} \left( \frac{2E l_e \left[ l_e + 2 l_{edrift} \right]}{l_b B \left[ l_b + 2 l_{bdrift} \right]} \right)^2 \right] . \] (D.32)

The energy resolution thus scales linearly with the aperture, scales with the energy as \( \propto E_p^{3/2} \) and scales with the charge and mass as \( \propto \sqrt{m/q} \).
D.5 Summary

In summary, the magnetic deflection, electric deflection, parabola slope and energy resolution derived in this appendix are given by the following equations:

\[ \Delta M = \frac{l_b \omega_c}{v_0} \left[ \frac{l_b}{2} + l_{bdrift} \right] + \frac{l_b^3 \omega_c^3}{v_0^3} \left[ \frac{l_{bdrift}}{2} \right] \]  
(D.33a)

\[ \Delta E = \frac{qE l_e}{mv_0^2} \left[ \frac{l_e}{2} + l_{edrift} \right] \left[ 1 + \frac{l_b^2 \omega_c^2}{v_0^2} \right] \]  
(D.33b)

\[ \alpha \equiv \frac{m}{q} \frac{2E l_e}{B^2 l_b^2} \left[ \frac{l_e}{2} + l_{edrift} \right]^2 \]  
(D.33c)

\[ \Delta E_p \approx -\frac{4a(1 + d/D)\sqrt{2} mE_p^{3/2}}{q l_b B \left[ l_b + 2 l_{bdrift} \right]} \left[ 1 - \frac{m}{E_p} \left( \frac{2 E l_e}{l_b B \left[ l_b + 2 l_{edrift} \right]} \right)^2 \right] \]  
(D.33d)

These equations are valid for uniform magnetic and electric fields. They neglect the presence of fringe fields, and hence are accurate to within just a few percent. The design parameters for each of the Thomson parabolas, including relevant scale lengths and field strengths for use in the above equations, are summarized in table D.1.
### Table D.1. Specifications of the TPIE and TPIS Thomson parabolas on OMEGA, OMEGA EP and on the MTW, respectively.

<table>
<thead>
<tr>
<th>Parameter</th>
<th>Symbol</th>
<th>TPIS</th>
<th>TPIE Analyzer</th>
</tr>
</thead>
<tbody>
<tr>
<td>Laser system</td>
<td></td>
<td>MTW</td>
<td>OMEGA, OMEGA EP</td>
</tr>
<tr>
<td>Detector size ($\Delta M \times \Delta E$)</td>
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<td>10 cm $\times$ 10 cm</td>
<td>10 cm $\times$ 5 cm</td>
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<tr>
<td>Energy range</td>
<td>$E_p$</td>
<td>0.8-20 MeV ($H^+$)</td>
<td>2-100 MeV ($H^+$)</td>
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<tr>
<td>Energy resolution</td>
<td>$\Delta E_p/E_p$</td>
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<td>1.5% at 2 MeV*</td>
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<td>$a$</td>
<td>400 $\mu$m (round)</td>
<td>100, 250, 1000 $\mu$m (square)</td>
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<tr>
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<td>$B$</td>
<td>5.3 kG</td>
<td>5.6 kG, 8.4 kG</td>
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<tr>
<td>Magnet pole gap</td>
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<td>Magnetic pole length</td>
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<td>5 cm</td>
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</tr>
<tr>
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<td>Aperture-detector distance</td>
<td>$D$</td>
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* Applies for 5.4 kG magnet
† Applies for 5.6 kG magnet, 250 $\mu$m aperture and 100 cm source to aperture distance
Appendix E

Thomson Parabola Analysis Code (TPAC)

E.1 Overview

The Thomson Parabola Analysis Code (TPAC) was developed to facilitate analysis of TPIS and TPIE data. An overview of the algorithms implemented in the code and a typical analysis procedure are discussed in this appendix.

The CR-39 detector used in TPIS and TPIE Analyzer is etched in an NaOH solution to reveal latent charged-particle tracks. The detectors are subsequently scanned using an automated microscope with a high-resolution CCD camera. Information about each particle track, including diameter, contrast, eccentricity and location are stored on disk using a proprietary file format. For more information on CR-39 response to charged-particles and processing techniques, see Appendix C.

The TPAC reads these data files in an interactive manner. Absolute information about each track is read into the code, together with file trailer information that contains operational parameters (e.g., the voltage on the electrodes, aperture size used, etc.). The tracks are analyzed and discriminated based on diameter, contrast, eccentricity and location on the detector.

The TPAC user interface is shown in Fig. E.1. It is a 32-bit Windows-based application that is written in C++. The underlying analysis routines make use of the standard template library (STL) and are structured in a way that facilitates the porting of these routines to different platforms. The TPAC has a number of features that ultimately allow the user to extract energy spectra for each ion species from the CR-39 data, as described below.

E.2 Analysis Procedure

The analysis procedure described here applies to both TPIS and the TPIE data. First, the data file containing track information is loaded into the software program. At this point, the software will scan the data file header information, which indicates whether the data were taken using TPIS or TPIE, to determine which instrument calibration data to use for subsequent analysis. It will furthermore extract information about the diagnostic setup for the shot, such as the aperture size, total voltage and the distance from the aperture to the source. These parameters can also be specified by the user, as shown in region 1 of Fig. E.1.

A fiducial check is then performed by generating an $N(x,y)$ histogram (see Fig. E.2). During an experiment, the fiducial on the CR-39 is generated by fast neutrals and x-rays that are not deflected by the electric and magnetic fields within the Thomson parabola. These particles leave tracks on the CR-39 detector within a circular region that is comparable in size and location to the projection of the aperture onto the detector. The fiducial is therefore a reference point for the
Figure E.1. Screen capture of the Thomson Parabola Analysis Code (Windows), indicating the various user options available for analysis of data acquired using the TPIS and TPIE spectrometers.

amount of magnetic and electric deflection. The center of the fiducial is set to the origin (0,0) of the 2D coordinate system, which is inferred by measuring the location of the top, bottom, left and right edges of the square fiducial. Ideally, the fiducial is found by visual inspection of the CR-39 detector before it is scanned. This is sometimes not feasible as the fiducial can be faint and not detectable by eye. In this case, the coordinate system origin is typically set to the corner of the piece nearest to where the fiducial is expected and must be found manually as follows. Using the TPAC, one can generate a plot showing contours of constant track density (N(x,y)) that will reveal the location of the fiducial (see 7). The coordinate system can then be adjusted to set the fiducial at the origin using the scan offset feature of the software (3 of Fig. E.1). Sample N(x,y) plots generated by the TPAC are shown in Fig. E.3 for an aligned and a mis-aligned fiducial.

Background and data regions are selected after the location of the fiducial has been verified and corrected. The data region limits the data analysis to a particular region of the piece. These limits apply to plots (e.g., N(α), etc.) and well as ion energy-spectra, Y(E). This feature is particularly
useful if the CR-39 detector contains streaks of localized noise in regions where no signal is present.
The presence of such noise can complicate the analysis and is readily disregarded by using data region limits (\(\text{\textsuperscript{4}}\)). The background region specifies a region of the piece for the TPAC to use as background (\(\text{\textsuperscript{5}}\)). This area is sampled and the number of charged-particle tracks per unit area are subtracted from the entire piece of CR-39. This simple background subtraction is applied when generating various plots of interest (e.g., \(N(x,y)\), \(N(\alpha)\), etc.).

From here on, and at any point during the analysis, tracks can be discriminated on the basis of contrast, diameter and eccentricity (\(\text{\textsuperscript{3}}\)). Experience has shown that tracks can have eccentricities as high as 40%, due to oblique angle of incidence of tracks on the CR-39 detector. It has also been observed that tracks with moderate \(q/m\) ratios (i.e., 0.3-0.5) have low contrast values (\(\sim 10\)), while those with low \(q/m\) ratios (< 0.3) have higher contrast values (\(\lesssim 70\)). In general, all tracks with contrast less than 60-70 should be selected for the analysis until individual parabolas are analyzed, as discussed further below.

In the next step of the analysis, an \(N(\alpha)\) histogram is generated (see Fig. E.4). As discussed in Sec 3.3 and Appendix D, \(\alpha\) is the slope of a parabola; it is directly proportional to the mass-to-charge ratio of ions, and it is also a function of instrument parameters (electric and magnetic field strengths, geometry, etc.). This histogram allows one to immediately see the parabolas and the relative signal strengths of each of the particle species that were detected. For each parabola, represented in the \(N(\alpha)\) histogram as a single peak, the user must select the mean, minimum and maximum \(\alpha\)-values that characterize the peak. Since each of the parabolas typically exhibit a near-Gaussian distribution in \(\alpha\)-space, Gaussian fits may be used to determine the mean slope. The minimum and maximum \(\alpha\)-values should be chosen based on where the peak falls into the background. It is therefore important that the background region be carefully selected before the \(N(\alpha)\) histogram is generated. The three \(\alpha\)-values that characterize a peak are then entered into the parabola table (\(\text{\textsuperscript{4}}\)). The TPAC will validate the these values and allow the user to select ions

Figure E.2. TPAC-generated \(N(x,y)\) histogram showing various parabolas and the fiducial.

Figure E.3. TPAC-generated \(N(x,y)\) plots, showing (a) an aligned fiducial and (b) a misaligned fiducial. The dashed lines indicate the coordinate system origin.
Figure E.4. TPAC-generated $N(\alpha)$ histogram showing definitions of $\alpha$, $\alpha_{\text{min}}$ and $\alpha_{\text{max}}$ for each parabola. The mass-to-charge ratio increases with increasing alpha. Each parabola has an associated minimum, maximum and mean value of $\alpha$ that can be determined by fitting the (approximately) Gaussian peaks.

that best fit the measured slope. Since the TPIS and TPIE calibrations have uncertainties of a few percent, several ions are available in a drop-down box for the user to select. Shown alongside each ion is the q/m ratio of the ion and the goodness-of-fit ($d\alpha$). The latter represents the normalized difference between the measured slope and the expected slope. Values (greater than) less than unity indicate that the selected ion is expected to have a slope that is (larger) smaller than the measured slope.

Energy spectra can be extracted after the parabola table has been populated. For each parabola, the user can generate a plot showing contours of tracks as a function of track contrast and diameter. This can be achieved by setting $\alpha$-limits (3 $\circ$) and then generating an N(c,d) plot. For convenience, this procedure has been implemented as an additional N(c,d) button alongside each parabola (9 $\circ$). This plot allows the user to make the appropriate contrast and diameter cuts for that parabola. After these cuts have been determined and entered into the track limits region (3 $\circ$) the Y(E) button can be used to generate energy spectra for that parabola.

Two subroutines for background subtraction of energy spectra (Y(E)) have been implemented in the TPAC (6 $\circ$). The first sub-routine performs a flat background subtraction, where the specified background region is sampled and used as a measure of the background for the entire detector. The second sub-routine is adaptive, and for each energy bin along a parabola, looks just above and below to determine the local background level (see Fig. E.5). The background is interpolated through the parabola in the direction of electric displacement. Although this adaptive algorithm is generally preferable as it samples the noise locally along the parabola, it fails when parabolas are not widely separated. In this case, the sampled background is large compared to the real background and particle tracks are under-counted. For such cases, the flat subtraction algorithm should be used instead. For the flat algorithm, care must be taken to ensure that the selected background region has sufficient statistics and little concentrated (localized) noise. Fig. E.6 shows the effect

Figure E.5. Implementation of adaptive background subtraction in the TPAC. As shown in the inset, each parabola is divided up into bins along the magnetic dispersion axis. At each bin, the code looks just above and below the parabola to determine the local background in these regions. A linear interpolation is then used to determine the background through the parabola (along the electric dispersion axis).
Figure E.6. TPAC-generated energy spectra of singly-ionized aluminum (for this shot, the background was particularly high and the maximum ion energy is not readily measured). Shown are energy spectra (a) without background subtraction (b) with an adaptive subtraction algorithm and (c) with a flat subtraction algorithm.

of background subtraction on exponential fast-ion spectra (for a spectrum with a particularly high background).

E.3 Thomson Parabola Calibration Data

The identification of ions and extraction of energy spectra are achieved by accessing TPIS and TPIE calibration data. The TPAC stores this calibration data, which includes the magnet energy scaling, the slope of an ion species per kV of voltage across the electrodes, the maximum allowable voltage and the size of the detectors.

The calibration data was obtained using a simulation code (written in C++) that tracks particle trajectories through magnetic and electric fields. For the TPIS, a series of simulations were run for ions with a range of q/m ratios and energies to determine the magnet energy scaling (see Eq. D.33a of Appendix D) and the slope of a parabola associated with an ion species (see Eq. D.33c of Appendix D). The trajectory of each ion as it traversed the fields was simulated until it reached the CR-39 detector, after which the resulting magnetic and electric deflections were measured. Similar simulations were developed and run by LANL for TPIE. In both cases, it was found that the fringe fields have a very small impact for ions within the range of energies and q/m ratios that can be measured using the TPIS and TPIE. In particular, the magnet energy scaling and parabola slopes were found to scale with $q/\sqrt{m}$ and $m/q$, as expected.

The scalings produced by the simulation code are stored in TPAC and primarily used for data analysis. They can also be used to simulate the expected signal on the CR-39 detector. Users can select a set of operating conditions (voltage, aperture, etc., as shown in (1) of Fig. E.1) and plot the energy resolution, the magnetic deflection, and the parabolas for ions with a range of energies and q/m ratios. This feature facilitates experiment planning and can be found under the calibration menu (see Fig. E.1).
Appendix F

OMEGA Fast Ion Experimental Data Summary

A summary of fast ion data from warm glass, plastic and cryogenic targets are presented in this appendix. Heavy fast-ion measurements are presented in Sec. F.1. This is followed by data from warm and cryogenic targets (Sec. F.2) that were used to investigate the scaling of maximum proton energies (peak target voltage) with laser intensity and shell material.
F.1 Heavy Fast-Ions from Warm Plastic and Glass Targets

A summary of shots used to study heavy fast-ions are shown in Table F.1. Shown are the target and laser parameters for warm glass and plastic shots, alongside a summary of the total fast-ion energy and maximum proton energy for each shot. Following the table are data from the CR-39 detector from TPIE, with selected heavy-ion spectra and a summary of TPIE operating parameters. TPIE voltages and electric fields quoted in this appendix are measured values for each shot.

Table F.1. Summary of laser and target parameters and fast ion energetics for all heavy fast-ion measurements.

| Shot  | Shell | Shell Thick. | Pulse | Laser Energy | OD | Intensity $\times 10^{14}$ | Max. Proton Energy | Total Proton Energy | Total Heavy-Ion Energy | Total Ion Energy |%
<table>
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<td>2.2</td>
<td>SG1018</td>
<td>29.7</td>
<td>860.0</td>
<td>13</td>
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<td>94.5</td>
<td>158.2</td>
<td>0.5</td>
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<td>132.9</td>
<td>61.0</td>
<td>193.8</td>
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<td>SG1018</td>
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<td>13</td>
<td>1.2</td>
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<td>77.4</td>
<td>154.9</td>
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</tr>
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<td>SG1018</td>
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<td>858.2</td>
<td>13</td>
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<td>106.0</td>
<td>288.7</td>
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</tr>
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<td>SG1018</td>
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<td>861.0</td>
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<td>139.9</td>
<td>102.0</td>
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*Flash-coated with a 100-nm-thick layer of aluminum
Figure F.1. (a) Data recorded by a CR-39 detector in TPIE for OMEGA shot 64681. TPIE was fielded 130 cm from TCC with an aperture of 250 μm, a peak magnetic field of 5.4 kG, a total voltage of 2.9 kV (E = 2.9 kV-cm$^{-1}$) and a 1.5-μm-thick mylar filter over the CR-39. The target was irradiated at an intensity of $1.2 \times 10^{15}$ W-cm$^{-2}$. (b) Selected heavy-ion spectra measured by TPIE are shown alongside the proton spectrum measured by CPS.
Figure F.2. (a) Data recorded by a CR-39 detector in TPIE for OMEGA shot 64684. TPIE was fielded 130 cm from TCC with an aperture of 250 μm, a peak magnetic field of 5.4 kG, a total voltage of 3.0 kV (E = 3.0 kV-cm⁻¹) and a 1.5-μm-thick mylar filter over the CR-39. The target was irradiated at an intensity of 1.3 × 10¹⁵ W-cm⁻². (b) Selected heavy-ion spectra measured by TPIE are shown alongside the proton spectrum measured by CPS.

Figure F.3. (a) Data recorded by a CR-39 detector in TPIE for OMEGA shot 64685. TPIE was fielded 130 cm from TCC with an aperture of 250 μm, a peak magnetic field of 5.4 kG, a total voltage of 3.1 kV (E = 3.1 kV-cm⁻¹) and a 1.5-μm-thick mylar filter over the CR-39. The target was irradiated at an intensity of 5 × 10¹⁵ W-cm⁻². (b) Selected heavy-ion spectra measured by TPIE are shown alongside the proton spectrum measured by CPS.
Figure F.4. (a) Data recorded by a CR-39 detector in TPIE for OMEGA shot 64686. TPIE was fielded 130 cm from TCC with an aperture of 250 μm, a peak magnetic field of 5.4 kG, a total voltage of 3.8 kV (E = 3.8 kV-cm$^{-1}$) and a 1.5-μm-thick mylar filter over the CR-39. The target was irradiated at an intensity of $1.3 \times 10^{15}$ W-cm$^{-2}$. Selected heavy-ion spectra measured by TPIE are shown alongside the proton spectrum measured by CPS.

Figure F.5. (a) Data recorded by a CR-39 detector in TPIE for OMEGA shot 65266. TPIE was fielded 130 cm from TCC with an aperture of 250 μm, a peak magnetic field of 5.4 kG, a total voltage of 2.8 kV (E = 2.8 kV-cm$^{-1}$) and a 0.9-μm-thick mylar filter over the CR-39. The target was irradiated at an intensity of $1.3 \times 10^{15}$ W-cm$^{-2}$. (b) Selected heavy-ion spectra measured by TPIE.
Figure F.6. (a) Data recorded by a CR-39 detector in TPIE for OMEGA shot 65269. TPIE was fielded 130 cm from TCC with an aperture of 250 μm, a peak magnetic field of 5.4 kG, a total voltage of 3.3 kV (E = 3.3 kV-cm\(^{-1}\)) and a 0.9-μm-thick mylar filter over the CR-39. The target was irradiated at an intensity of 1.3 \times 10^{14} \text{ W-cm}^{-2}. (b) Selected heavy-ion spectra measured by TPIE.

Figure F.7. (a) Data recorded by a CR-39 detector in TPIE for OMEGA shot 65271. TPIE was fielded 130 cm from TCC with an aperture of 250 μm, a peak magnetic field of 5.4 kG, a total voltage of 2.5 kV (E = 2.5 kV-cm\(^{-1}\)) and a 0.9-μm-thick mylar filter over the CR-39. The target was irradiated at an intensity of 1.2 \times 10^{14} \text{ W-cm}^{-2}. (b) Selected heavy-ion spectra measured by TPIE.
Figure F.8. (a) Data recorded by a CR-39 detector in TPIE for OMEGA shot 65273. TPIE was fielded 130 cm from TCC with an aperture of 250 μm, a peak magnetic field of 5.4 kG, a total voltage of 2.9 kV (E = 2.9 kV-cm$^{-1}$) and a 0.9-μm-thick mylar filter over the CR-39. The target was irradiated at an intensity of 1.2 × 10$^{14}$ W-cm$^{-2}$. (b) Selected heavy-ion spectra measured by TPIE are shown alongside the proton spectrum measured by CPS.
F.2 Scaling of Maximum Ion Energies and Total Ion Energies

The scaling of the maximum proton energy with laser intensity is shown in Fig. F.9. The maximum proton energies are comparable to the results obtained by Hicks et al, although those results have been extended here by including cryogenic and thin-plastic shells. The scaling of the maximum proton and heavy-ion energies with the ion charge state are shown in Figs. F.10 and F.11 for thin glass and plastic shells, respectively. These data confirm that the maximum energies of all the fast ions scale linearly with the ion charge state, as indicated by the linear fits (dashed lines) in Figs. F.10 and F.11. Finally, the scaling of the total ion energy with maximum proton energy is summarized in Fig. F.12. The numerical values used to produce the figures in this section are summarized in several tables. Data for cryogenic implosions are shown in table F.2, while glass and plastic implosions are shown in tables F.3 and F.4.

![Figure F.9.](image)

Figure F.9. Peak target voltage as a function of the on-target laser intensity. Cryogenic and warm thin plastic shells (< 6 μm) exhibit a linear dependence with laser intensity. Thick plastic shells (> 10 μm) and thin glass shells (2 – 3 μm) also follow a linear scaling with intensity, but result in peak voltages that are systematically lower. These discrepancies are due to differences in shell laser absorption and electron stopping power.
Figure F.10. Maximum ion energies of protons and heavy ions (primarily silicon and oxygen) as a function of the ion charge state for thin-glass shells (2 – 3 μm). For each shot, the dashed line is a linear fit to the data.

Figure F.11. Maximum ion energies of protons and heavy ions (primarily aluminum) as a function of the ion charge state for Al-coated plastic shells (< 6 μm). For each shot, the dashed line is a linear fit to the data.
Figure F.12. Percentage of incident laser energy carried by fast ions, including protons from CH and SiO$_2$ shells ($\lesssim 0.3\%$), protons and deuterons from CD shells ($\lesssim 0.6\%$), and heavier fast ions CD and SiO$_2$ ($\lesssim 0.3\%$). The total incident laser energy coupled to fast ions (dashed curve) is estimated by scaling the CD dataset (solid blue circles) to match the total ion measurements (open blue circles). Where not shown, the error-bars are comparable to the symbol size.
Table F.2. Summary of laser and target parameters and fast ion energetics for cryogenic targets.

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<th>Shot</th>
<th>Shell</th>
<th>Shell Thick. μm</th>
<th>Pulse</th>
<th>Laser Energy kJ</th>
<th>OD Intensity ×10¹⁴ W-cm⁻²</th>
<th>Max. Proton Energy MeV</th>
<th>Total Proton Energy J</th>
<th>Total Ion Energy J (%)</th>
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**Table F.3.** Summary of laser and target parameters and fast ion energetics for thin-glass targets.

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*Flash-coated with a 70-nm-thick layer of aluminum
## Table F.4. Summary of laser and target parameters and fast ion energetics for plastic targets.

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Appendix G

OMEGA Fast Ignition Experimental Data Summary

A summary of fast-ignition shots on OMEGA are presented in Table G.1. Shown are the shot number, target parameters and laser parameters for both the OMEGA and OMEGA EP lasers. The focal radii quoted for each shot are known as “R80” or the radius that encircles 80% of the laser energy. The arrival time is the time at which the OMEGA EP laser arrives at the cone-tip relative to the start of the OMEGA lasers (the compression drive). All of the parameter values shown here were measured. The targets used in these experiments were all hollow spherical shells with a nominal outer diameter of 860 μm; they were not gas-filled. Other target parameters were varied from shot-to-shot, as indicated by the shorthand notation (for example “CD[40]Au[5]d10_1.8mm” refers to a 40 μm-thick spherical plastic (CD) shell embedded with a gold cone with a tip thickness of 5 μm, a tip diameter of 10 μm, and a cone length of 1.8 mm). All cones had a opening half-angle of 17° and cone-wall thickness of 50 μm.

Following the shot summary table are raw data from the OMEGA nTOF detector (Fig. G.1) and fast-proton energy spectra measured by CPS1 and CPS2 (Fig. G.2) for several shots. This appendix concludes with a table for each fast-ignition shot. Each table lists detailed laser and target parameters and also includes the maximum fast-proton energies as measured by the WRF at various angles around the implosion.
Table G.1. Summary of pertinent experimental parameters for all of the fast-ignition shots presented in this work.

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<td>56247</td>
<td>CD[40.3]Au[10]d10_1.2mm</td>
<td>19.789</td>
<td>1006</td>
<td>10</td>
<td>24.0</td>
<td>4.45</td>
<td>3.95</td>
</tr>
<tr>
<td>56971</td>
<td>CD[40]Au[16]d42_1.2mm</td>
<td>18.986</td>
<td>1055</td>
<td>11</td>
<td>24.0</td>
<td>4.24</td>
<td>3.60</td>
</tr>
<tr>
<td>56976</td>
<td>CD[38.6]Au[14]d10_1.2mm</td>
<td>18.83</td>
<td>-</td>
<td>-</td>
<td>-</td>
<td>-</td>
<td>-</td>
</tr>
<tr>
<td>59124</td>
<td>CD[38.0]Au[5]d10_1.2mm</td>
<td>18.728</td>
<td>933</td>
<td>9.2</td>
<td>27.9</td>
<td>3.32</td>
<td>-</td>
</tr>
<tr>
<td>59126</td>
<td>CD[38.6]Au[5]d10_1.2mm</td>
<td>18.535</td>
<td>942</td>
<td>8.7</td>
<td>27.1</td>
<td>3.76</td>
<td>-</td>
</tr>
</tbody>
</table>
Figure G.1. Neutron time-of-flight signal (Ch. 1 and 2) from fast-ignition experiments at Omega showing the x-ray flash, 2.45 MeV DD-n signal, and neutrons from (p,n). For these shots, all laser and target parameters were identical with the exception of the cone-length, which was 50% greater for shot 56972 (middle row), corresponding to 2.25 times more surface area. The ratio of the total (p,n) signal for 1.8 mm cones is a factor of $\sim 2 - 3$ larger than for 1.2 mm cones, which is roughly proportional to the ratio of the cone surface areas.
Figure G.2. Fast-proton spectra for fast-ignition shots, measured by CPS1 and CPS2.
### Table G.2. Summary of target, laser and proton data for OMEGA shot 53043.

#### Shot 53043

<table>
<thead>
<tr>
<th>Target Metrology</th>
<th>Laser Parameters</th>
</tr>
</thead>
<tbody>
<tr>
<td>Shell Diameter</td>
<td>860 μm</td>
</tr>
<tr>
<td>Shell Thick.</td>
<td>40 (CD) μm</td>
</tr>
<tr>
<td>Cone Material</td>
<td>Au</td>
</tr>
<tr>
<td>Cone Tip Thick.</td>
<td>15 μm</td>
</tr>
<tr>
<td>Cone Tip Dia.</td>
<td>10 μm</td>
</tr>
<tr>
<td>Cone Length</td>
<td>1.2 mm</td>
</tr>
<tr>
<td>OMEGA Energy</td>
<td>17.85 kJ</td>
</tr>
<tr>
<td>OMEGA Pulse Shape</td>
<td>LA241701P</td>
</tr>
<tr>
<td>EP Energy</td>
<td>512 J</td>
</tr>
<tr>
<td>EP Focal Spot</td>
<td>24.0 μm</td>
</tr>
<tr>
<td>EP Intensity</td>
<td>$2.08 \times 10^{18}$ W-cm$^{-2}$</td>
</tr>
<tr>
<td>EP Pulse Width</td>
<td>11 ps</td>
</tr>
<tr>
<td>EP Arrival Time</td>
<td>3.59 ns</td>
</tr>
</tbody>
</table>

#### WRF Maximum Proton Energy

<table>
<thead>
<tr>
<th>Location</th>
<th>Measurements (MeV)</th>
<th>Average</th>
<th>Std. Dev</th>
</tr>
</thead>
<tbody>
<tr>
<td>TIM 3 (70°)</td>
<td>4.70 5.35</td>
<td>5.03 MeV</td>
<td>0.46 MeV</td>
</tr>
<tr>
<td>TIM 4 (80°)</td>
<td>6.44 6.00</td>
<td>6.04 MeV</td>
<td>0.30 MeV</td>
</tr>
<tr>
<td>TIM 5 (0°)</td>
<td>4.80 5.10</td>
<td>5.23 MeV</td>
<td>0.5 MeV</td>
</tr>
</tbody>
</table>

### Table G.3. Summary of target, laser and proton data for OMEGA shot 53527.

#### Shot 53527

<table>
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<th>Target Metrology</th>
<th>Laser Parameters</th>
</tr>
</thead>
<tbody>
<tr>
<td>Shell Diameter</td>
<td>860 μm</td>
</tr>
<tr>
<td>Shell Thick.</td>
<td>40 (CD) μm</td>
</tr>
<tr>
<td>Cone Material</td>
<td>Au</td>
</tr>
<tr>
<td>Cone Tip Thick.</td>
<td>10 μm</td>
</tr>
<tr>
<td>Cone Tip Dia.</td>
<td>10 μm</td>
</tr>
<tr>
<td>Cone Length</td>
<td>1.2 mm</td>
</tr>
<tr>
<td>OMEGA Energy</td>
<td>18.13 kJ</td>
</tr>
<tr>
<td>OMEGA Pulse Shape</td>
<td>LA241701P</td>
</tr>
<tr>
<td>EP Energy</td>
<td>820 J</td>
</tr>
<tr>
<td>EP Focal Spot</td>
<td>24.0 μm</td>
</tr>
<tr>
<td>EP Intensity</td>
<td>$2.79 \times 10^{18}$ W-cm$^{-2}$</td>
</tr>
<tr>
<td>EP Pulse Width</td>
<td>13 ps</td>
</tr>
<tr>
<td>EP Arrival Time</td>
<td>3.55 ns</td>
</tr>
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#### WRF Maximum Proton Energy

<table>
<thead>
<tr>
<th>Location</th>
<th>Measurements (MeV)</th>
<th>Average</th>
<th>Std. Dev</th>
</tr>
</thead>
<tbody>
<tr>
<td>TIM 3 (70°)</td>
<td>4.46 5.06 4.76 4.56 5.00</td>
<td>4.77 MeV</td>
<td>0.26 MeV</td>
</tr>
<tr>
<td>TIM 4 (80°)</td>
<td>6.44 6.00 6.01 5.71</td>
<td>6.04 MeV</td>
<td>0.30 MeV</td>
</tr>
<tr>
<td>TIM 5 (0°)</td>
<td>4.80 5.10 5.78</td>
<td>5.23 MeV</td>
<td>0.5 MeV</td>
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</tbody>
</table>
### Table G.4. Summary of target, laser and proton data for OMEGA shot 53528.

<table>
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<th>Target Metrology</th>
<th>Laser Parameters</th>
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</thead>
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<td>Shell Diameter</td>
<td>860 μm</td>
<td>OMEGA Energy</td>
</tr>
<tr>
<td>Shell Thick.</td>
<td>40 (CD) μm</td>
<td>OMEGA Pulse Shape</td>
</tr>
<tr>
<td>Cone Material</td>
<td>Au</td>
<td>EP Energy</td>
</tr>
<tr>
<td>Cone Tip Thick.</td>
<td>5 μm</td>
<td>EP Focal Spot</td>
</tr>
<tr>
<td>Cone Tip Dia.</td>
<td>10 μm</td>
<td>EP Intensity</td>
</tr>
<tr>
<td>Cone Length</td>
<td>1.2 mm</td>
<td>EP Pulse Width</td>
</tr>
<tr>
<td></td>
<td></td>
<td>EP Arrival Time</td>
</tr>
</tbody>
</table>

<table>
<thead>
<tr>
<th>WRF Maximum Proton Energy</th>
<th>Location</th>
<th>Measurements (MeV)</th>
<th>Average</th>
<th>Std. Dev</th>
</tr>
</thead>
<tbody>
<tr>
<td>TIM 3 (70°)</td>
<td>5.80</td>
<td>5.82</td>
<td>5.60</td>
<td>-</td>
</tr>
<tr>
<td>TIM 4 (80°)</td>
<td>7.30</td>
<td>7.10</td>
<td>6.90</td>
<td>6.40</td>
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<tr>
<td>TIM 5 (0°)</td>
<td>4.97</td>
<td>5.50</td>
<td>6.46</td>
<td>5.70</td>
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### Table G.5. Summary of target, laser and proton data for OMEGA shot 53529.

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<th>Target Metrology</th>
<th>Laser Parameters</th>
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<tbody>
<tr>
<td>Shell Diameter</td>
<td>860 μm</td>
<td>OMEGA Energy</td>
</tr>
<tr>
<td>Shell Thick.</td>
<td>38.1 (CD), 2 (CH) μm</td>
<td>OMEGA Pulse Shape</td>
</tr>
<tr>
<td>Cone Material</td>
<td>Au</td>
<td>EP Energy</td>
</tr>
<tr>
<td>Cone Tip Thick.</td>
<td>10 μm</td>
<td>EP Focal Spot</td>
</tr>
<tr>
<td>Cone Tip Dia.</td>
<td>10 μm</td>
<td>EP Intensity</td>
</tr>
<tr>
<td>Cone Length</td>
<td>1.2 mm</td>
<td>EP Pulse Width</td>
</tr>
<tr>
<td></td>
<td></td>
<td>EP Arrival Time</td>
</tr>
</tbody>
</table>

<table>
<thead>
<tr>
<th>WRF Maximum Proton Energy</th>
<th>Location</th>
<th>Measurements (MeV)</th>
<th>Average</th>
<th>Std. Dev</th>
</tr>
</thead>
<tbody>
<tr>
<td>TIM 3 (70°)</td>
<td>7.00</td>
<td>7.50</td>
<td>-</td>
<td>-</td>
</tr>
<tr>
<td>TIM 4 (80°)</td>
<td>7.40</td>
<td>7.18</td>
<td>7.40</td>
<td>7.34</td>
</tr>
<tr>
<td>TIM 5 (0°)</td>
<td>5.30</td>
<td>5.50</td>
<td>5.80</td>
<td>-</td>
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### Table G.6. Summary of target, laser and proton data for OMEGA shot 54409.

<table>
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<th>Laser Parameters</th>
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</thead>
<tbody>
<tr>
<td>Shell Diameter</td>
<td>860 μm</td>
<td>OMEGA Energy</td>
</tr>
<tr>
<td>Shell Thick.</td>
<td>39.5 (CD) μm</td>
<td>OMEGA Pulse Shape</td>
</tr>
<tr>
<td>Cone Material</td>
<td>Au</td>
<td>EP Energy</td>
</tr>
<tr>
<td>Cone Tip Thick.</td>
<td>20 μm</td>
<td>EP Focal Spot</td>
</tr>
<tr>
<td>Cone Tip Dia.</td>
<td>10 μm</td>
<td>EP Intensity</td>
</tr>
<tr>
<td>Cone Length</td>
<td>1.2 mm</td>
<td>EP Pulse Width</td>
</tr>
<tr>
<td>EP Arrival Time</td>
<td>3.43 ns</td>
<td></td>
</tr>
</tbody>
</table>

<table>
<thead>
<tr>
<th>WRF Maximum Proton Energy</th>
</tr>
</thead>
<tbody>
<tr>
<td>Location</td>
</tr>
<tr>
<td>TIM 3 (70°)</td>
</tr>
<tr>
<td>TIM 4 (80°)</td>
</tr>
<tr>
<td>TIM 5 (0°)</td>
</tr>
</tbody>
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### Table G.7. Summary of target, laser and proton data for OMEGA shot 54409.

<table>
<thead>
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<th>Target Metrology</th>
<th>Laser Parameters</th>
</tr>
</thead>
<tbody>
<tr>
<td>Shell Diameter</td>
<td>860 μm</td>
<td>OMEGA Energy</td>
</tr>
<tr>
<td>Shell Thick.</td>
<td>39.5 (CD) μm</td>
<td>OMEGA Pulse Shape</td>
</tr>
<tr>
<td>Cone Material</td>
<td>Cu</td>
<td>EP Energy</td>
</tr>
<tr>
<td>Cone Tip Thick.</td>
<td>20 μm</td>
<td>EP Focal Spot</td>
</tr>
<tr>
<td>Cone Tip Dia.</td>
<td>10 μm</td>
<td>EP Intensity</td>
</tr>
<tr>
<td>Cone Length</td>
<td>1.2 mm</td>
<td>EP Pulse Width</td>
</tr>
<tr>
<td>EP Arrival Time</td>
<td>3.43 ns</td>
<td></td>
</tr>
</tbody>
</table>

<table>
<thead>
<tr>
<th>WRF Maximum Proton Energy</th>
</tr>
</thead>
<tbody>
<tr>
<td>Location</td>
</tr>
<tr>
<td>TIM 3 (70°)</td>
</tr>
<tr>
<td>TIM 4 (80°)</td>
</tr>
<tr>
<td>TIM 5 (0°)</td>
</tr>
</tbody>
</table>
Table G.8. Summary of target, laser and proton data for OMEGA shot 54410.

**Shot 54410**

<table>
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<th>Laser Parameters</th>
</tr>
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<td><strong>Shell Diameter</strong></td>
<td>860 μm</td>
</tr>
<tr>
<td><strong>Shell Thick.</strong></td>
<td>39.3 (CD) μm</td>
</tr>
<tr>
<td><strong>Cone Material</strong></td>
<td>Cu</td>
</tr>
<tr>
<td><strong>Cone Tip Thick.</strong></td>
<td>30 μm</td>
</tr>
<tr>
<td><strong>Cone Tip Dia.</strong></td>
<td>10 μm</td>
</tr>
<tr>
<td><strong>Cone Length</strong></td>
<td>1.2 mm</td>
</tr>
</tbody>
</table>

<table>
<thead>
<tr>
<th><strong>WRF Maximum Proton Energy</strong></th>
</tr>
</thead>
<tbody>
<tr>
<td><strong>Location</strong></td>
</tr>
<tr>
<td>TIM 3 (70°)</td>
</tr>
<tr>
<td>TIM 4 (80°)</td>
</tr>
<tr>
<td>TIM 5 (0°)</td>
</tr>
</tbody>
</table>

Table G.9. Summary of target, laser and proton data for OMEGA shot 54411.

**Shot 54411**

<table>
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<th>Laser Parameters</th>
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<tbody>
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<tr>
<td><strong>Shell Thick.</strong></td>
<td>39.3 (CD) μm</td>
</tr>
<tr>
<td><strong>Cone Material</strong></td>
<td>Au</td>
</tr>
<tr>
<td><strong>Cone Tip Thick.</strong></td>
<td>15 μm</td>
</tr>
<tr>
<td><strong>Cone Tip Dia.</strong></td>
<td>10 μm</td>
</tr>
<tr>
<td><strong>Cone Length</strong></td>
<td>1.2 mm</td>
</tr>
<tr>
<td><strong>EP Arrival Time</strong></td>
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</table>

<table>
<thead>
<tr>
<th><strong>WRF Maximum Proton Energy</strong></th>
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</thead>
<tbody>
<tr>
<td><strong>Location</strong></td>
</tr>
<tr>
<td>TIM 3 (70°)</td>
</tr>
<tr>
<td>TIM 4 (80°)</td>
</tr>
<tr>
<td>TIM 5 (0°)</td>
</tr>
</tbody>
</table>
### Table G.10. Summary of target, laser and proton data for OMEGA shot 54412.

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<td>OMEGA Pulse Shape</td>
</tr>
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<td>Cone Material</td>
<td>Au</td>
<td>EP Energy</td>
</tr>
<tr>
<td>Cone Tip Thick.</td>
<td>5 μm</td>
<td>EP Focal Spot</td>
</tr>
<tr>
<td>Cone Tip Dia.</td>
<td>10 μm</td>
<td>EP Intensity</td>
</tr>
<tr>
<td>Cone Length</td>
<td>1.2 mm</td>
<td>EP Pulse Width</td>
</tr>
<tr>
<td></td>
<td></td>
<td>EP Arrival Time</td>
</tr>
<tr>
<td></td>
<td></td>
<td>WRF Maximum Proton Energy</td>
</tr>
<tr>
<td>Location</td>
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<td>Average</td>
</tr>
<tr>
<td>TIM 3 (70°)</td>
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<td>-</td>
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<tr>
<td>TIM 4 (80°)</td>
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<td>-</td>
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<td>TIM 5 (0°)</td>
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### Table G.11. Summary of target, laser and proton data for OMEGA shot 55148.

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<th>Laser Parameters</th>
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<td>OMEGA Energy</td>
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<tr>
<td>Shell Thick.</td>
<td>40.0 (CD) μm</td>
<td>OMEGA Pulse Shape</td>
</tr>
<tr>
<td>Cone Material</td>
<td>Au</td>
<td>EP Energy</td>
</tr>
<tr>
<td>Cone Tip Thick.</td>
<td>15 μm</td>
<td>EP Focal Spot</td>
</tr>
<tr>
<td>Cone Tip Dia.</td>
<td>10 μm</td>
<td>EP Intensity</td>
</tr>
<tr>
<td>Cone Length</td>
<td>1.2 mm</td>
<td>EP Pulse Width</td>
</tr>
<tr>
<td></td>
<td></td>
<td>EP Arrival Time</td>
</tr>
<tr>
<td></td>
<td></td>
<td>WRF Maximum Proton Energy</td>
</tr>
<tr>
<td>Location</td>
<td>Measurements (MeV)</td>
<td>Average</td>
</tr>
<tr>
<td>TIM 3 (70°)</td>
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<td>-</td>
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<tr>
<td>TIM 4 (80°)</td>
<td>-</td>
<td>-</td>
</tr>
<tr>
<td>TIM 5 (0°)</td>
<td>5.20</td>
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Table G.12. Summary of target, laser and proton data for OMEGA shot 55150.

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<th>Laser Parameters</th>
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<td>OMEGA Energy</td>
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<tr>
<td>Shell Thick.</td>
<td>38.7 (CD) μm</td>
<td>OMEGA Pulse Shape</td>
</tr>
<tr>
<td>Cone Material</td>
<td>Au</td>
<td>EP Energy</td>
</tr>
<tr>
<td>Cone Tip Thick.</td>
<td>16 μm</td>
<td>EP Focal Spot</td>
</tr>
<tr>
<td>Cone Tip Dia.</td>
<td>10 μm</td>
<td>EP Intensity</td>
</tr>
<tr>
<td>Cone Length</td>
<td>1.2 mm</td>
<td>EP Pulse Width</td>
</tr>
<tr>
<td></td>
<td></td>
<td>EP Arrival Time</td>
</tr>
</tbody>
</table>

<table>
<thead>
<tr>
<th>Location</th>
<th>Measurements (MeV)</th>
<th>Average</th>
<th>Std. Dev</th>
</tr>
</thead>
<tbody>
<tr>
<td>TIM 3 (70°)</td>
<td>5.30 6.50</td>
<td>-</td>
<td>5.90 MeV 0.85 MeV</td>
</tr>
<tr>
<td>TIM 4 (80°)</td>
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<td>-</td>
<td>-</td>
</tr>
<tr>
<td>TIM 5 (0°)</td>
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</table>

Table G.13. Summary of target, laser and proton data for OMEGA shot 55154.

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<th>Laser Parameters</th>
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<td>OMEGA Energy</td>
</tr>
<tr>
<td>Shell Thick.</td>
<td>38.7 (CD) μm</td>
<td>OMEGA Pulse Shape</td>
</tr>
<tr>
<td>Cone Material</td>
<td>Au</td>
<td>EP Energy</td>
</tr>
<tr>
<td>Cone Tip Thick.</td>
<td>16 μm</td>
<td>EP Focal Spot</td>
</tr>
<tr>
<td>Cone Tip Dia.</td>
<td>10 μm</td>
<td>EP Intensity</td>
</tr>
<tr>
<td>Cone Length</td>
<td>1.2 mm</td>
<td>EP Pulse Width</td>
</tr>
<tr>
<td></td>
<td></td>
<td>EP Arrival Time</td>
</tr>
</tbody>
</table>

<table>
<thead>
<tr>
<th>Location</th>
<th>Measurements (MeV)</th>
<th>Average</th>
<th>Std. Dev</th>
</tr>
</thead>
<tbody>
<tr>
<td>TIM 3 (70°)</td>
<td>6.20</td>
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<td>6.20 MeV -</td>
</tr>
<tr>
<td>TIM 4 (80°)</td>
<td>-</td>
<td>-</td>
<td>-</td>
</tr>
<tr>
<td>TIM 5 (0°)</td>
<td>-</td>
<td>-</td>
<td>-</td>
</tr>
</tbody>
</table>
Table G.14. Summary of target, laser and proton data for OMEGA shot 56242.

Shot 56242

<table>
<thead>
<tr>
<th>Target Metrology</th>
<th>Laser Parameters</th>
</tr>
</thead>
<tbody>
<tr>
<td>Shell Diameter 860 μm</td>
<td>OMEGA Energy 19.92 kJ</td>
</tr>
<tr>
<td>Shell Thick. 40.5 (CD) μm</td>
<td>OMEGA Pulse Shape LA241701P</td>
</tr>
<tr>
<td>Cone Material Au</td>
<td>EP Energy 995 J</td>
</tr>
<tr>
<td>Cone Tip Thick. 11 μm</td>
<td>EP Focal Spot 24.0 μm</td>
</tr>
<tr>
<td>Cone Tip Dia. 10 μm</td>
<td>EP Intensity 4.40×10¹⁸ W·cm⁻²</td>
</tr>
<tr>
<td>Cone Length 1.2 mm</td>
<td>EP Pulse Width 10 ps</td>
</tr>
<tr>
<td></td>
<td>EP Arrival Time 3.43 ns</td>
</tr>
</tbody>
</table>

<table>
<thead>
<tr>
<th>WRF Maximum Proton Energy</th>
</tr>
</thead>
<tbody>
<tr>
<td>Location</td>
</tr>
<tr>
<td>TIM 3 (70°)</td>
</tr>
<tr>
<td>TIM 4 (80°)</td>
</tr>
<tr>
<td>TIM 5 (0°)</td>
</tr>
</tbody>
</table>

Table G.15. Summary of target, laser and proton data for OMEGA shot 56245.

Shot 56245

<table>
<thead>
<tr>
<th>Target Metrology</th>
<th>Laser Parameters</th>
</tr>
</thead>
<tbody>
<tr>
<td>Shell Diameter 860 μm</td>
<td>OMEGA Energy 19.37 kJ</td>
</tr>
<tr>
<td>Shell Thick. 40.4 (CD) μm</td>
<td>OMEGA Pulse Shape LA241701P</td>
</tr>
<tr>
<td>Cone Material Au</td>
<td>EP Energy 996 J</td>
</tr>
<tr>
<td>Cone Tip Thick. 11 μm</td>
<td>EP Focal Spot 24.0 μm</td>
</tr>
<tr>
<td>Cone Tip Dia. 10 μm</td>
<td>EP Intensity 3.83×10¹⁸ W·cm⁻²</td>
</tr>
<tr>
<td>Cone Length 1.2 mm</td>
<td>EP Pulse Width 12 ps</td>
</tr>
<tr>
<td></td>
<td>EP Arrival Time 3.75 ns</td>
</tr>
</tbody>
</table>

<table>
<thead>
<tr>
<th>WRF Maximum Proton Energy</th>
</tr>
</thead>
<tbody>
<tr>
<td>Location</td>
</tr>
<tr>
<td>TIM 3 (70°)</td>
</tr>
<tr>
<td>TIM 4 (80°)</td>
</tr>
<tr>
<td>TIM 5 (0°)</td>
</tr>
</tbody>
</table>
### Table G.16. Summary of target, laser and proton data for OMEGA shot 56247.

**Shot 56247**

<table>
<thead>
<tr>
<th>Target Metrology</th>
<th>Laser Parameters</th>
</tr>
</thead>
<tbody>
<tr>
<td>Shell Diameter</td>
<td>860 μm</td>
</tr>
<tr>
<td>Shell Thick.</td>
<td>40.3 (CD) μm</td>
</tr>
<tr>
<td>Cone Material</td>
<td>Au</td>
</tr>
<tr>
<td>Cone Tip Thick.</td>
<td>10 μm</td>
</tr>
<tr>
<td>Cone Tip Dia.</td>
<td>10 μm</td>
</tr>
<tr>
<td>Cone Length</td>
<td>1.2 mm</td>
</tr>
<tr>
<td></td>
<td></td>
</tr>
</tbody>
</table>

**WRF Maximum Proton Energy**

<table>
<thead>
<tr>
<th>Location</th>
<th>Measurements (MeV)</th>
<th>Average</th>
<th>Std. Dev</th>
</tr>
</thead>
<tbody>
<tr>
<td>TIM 3 (70°)</td>
<td>4.81 4.95 4.56 4.80</td>
<td>4.78 MeV</td>
<td>0.16 MeV</td>
</tr>
<tr>
<td>TIM 4 (80°)</td>
<td>- - - -</td>
<td>- -</td>
<td>- -</td>
</tr>
<tr>
<td>TIM 5 (0°)</td>
<td>4.91 4.74 4.98 5.00</td>
<td>4.91 MeV</td>
<td>0.12 MeV</td>
</tr>
</tbody>
</table>

### Table G.17. Summary of target, laser and proton data for OMEGA shot 56971.

**Shot 56971**

<table>
<thead>
<tr>
<th>Target Metrology</th>
<th>Laser Parameters</th>
</tr>
</thead>
<tbody>
<tr>
<td>Shell Diameter</td>
<td>860 μm</td>
</tr>
<tr>
<td>Shell Thick.</td>
<td>40 (CD) μm</td>
</tr>
<tr>
<td>Cone Material</td>
<td>Au</td>
</tr>
<tr>
<td>Cone Tip Thick.</td>
<td>16 μm</td>
</tr>
<tr>
<td>Cone Tip Dia.</td>
<td>42 μm</td>
</tr>
<tr>
<td>Cone Length</td>
<td>1.2 mm</td>
</tr>
<tr>
<td></td>
<td></td>
</tr>
</tbody>
</table>

**WRF Maximum Proton Energy**

<table>
<thead>
<tr>
<th>Location</th>
<th>Measurements (MeV)</th>
<th>Average</th>
<th>Std. Dev</th>
</tr>
</thead>
<tbody>
<tr>
<td>TIM 3 (70°)</td>
<td>4.38 4.50 4.34 4.20</td>
<td>4.36 MeV</td>
<td>0.12 MeV</td>
</tr>
<tr>
<td>TIM 4 (80°)</td>
<td>- - - -</td>
<td>- -</td>
<td>- -</td>
</tr>
<tr>
<td>TIM 5 (0°)</td>
<td>4.00 4.30 4.50 - -</td>
<td>4.27 MeV</td>
<td>0.25 MeV</td>
</tr>
</tbody>
</table>
Table G.18. Summary of target, laser and proton data for OMEGA shot 56972.

Shot 56972

<table>
<thead>
<tr>
<th>Target Metrology</th>
<th>Laser Parameters</th>
</tr>
</thead>
<tbody>
<tr>
<td>Shell Diameter</td>
<td>OMEGA Energy</td>
</tr>
<tr>
<td>860 μm</td>
<td>18.32 kJ</td>
</tr>
<tr>
<td>Shell Thick.</td>
<td>OMEGA Pulse Shape</td>
</tr>
<tr>
<td>40 (CD) μm</td>
<td>LA241701P</td>
</tr>
<tr>
<td>Cone Material</td>
<td>EP Energy</td>
</tr>
<tr>
<td>Au</td>
<td>1060 J</td>
</tr>
<tr>
<td>Cone Tip Thick.</td>
<td>EP Focal Spot</td>
</tr>
<tr>
<td>15 μm</td>
<td>24.0 μm</td>
</tr>
<tr>
<td>Cone Tip Dia.</td>
<td>EP Intensity</td>
</tr>
<tr>
<td>42 μm</td>
<td>4.69×10^{18} W-cm^{-2}</td>
</tr>
<tr>
<td>Cone Length</td>
<td>EP Pulse Width</td>
</tr>
<tr>
<td>1.8 mm</td>
<td>10 ps</td>
</tr>
<tr>
<td></td>
<td>EP Arrival Time</td>
</tr>
<tr>
<td></td>
<td>3.66 ns</td>
</tr>
</tbody>
</table>

<table>
<thead>
<tr>
<th>WRF Maximum Proton Energy</th>
</tr>
</thead>
<tbody>
<tr>
<td>Location</td>
</tr>
<tr>
<td>TIM 3 (70°)</td>
</tr>
<tr>
<td>TIM 4 (80°)</td>
</tr>
<tr>
<td>TIM 5 (0°)</td>
</tr>
</tbody>
</table>

Table G.19. Summary of target, laser and proton data for OMEGA shot 56973.

Shot 56973

<table>
<thead>
<tr>
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<th>Laser Parameters</th>
</tr>
</thead>
<tbody>
<tr>
<td>Shell Diameter</td>
<td>OMEGA Energy</td>
</tr>
<tr>
<td>860 μm</td>
<td>18.47 kJ</td>
</tr>
<tr>
<td>Shell Thick.</td>
<td>OMEGA Pulse Shape</td>
</tr>
<tr>
<td>39.7 (CD) μm</td>
<td>LA241701P</td>
</tr>
<tr>
<td>Cone Material</td>
<td>EP Energy</td>
</tr>
<tr>
<td>Au</td>
<td>1065 J</td>
</tr>
<tr>
<td>Cone Tip Thick.</td>
<td>EP Focal Spot</td>
</tr>
<tr>
<td>17 μm</td>
<td>24.0 μm</td>
</tr>
<tr>
<td>Cone Tip Dia.</td>
<td>EP Intensity</td>
</tr>
<tr>
<td>42 μm</td>
<td>4.28×10^{18} W-cm^{-2}</td>
</tr>
<tr>
<td>Cone Length</td>
<td>EP Pulse Width</td>
</tr>
<tr>
<td>1.2 mm</td>
<td>11 ps</td>
</tr>
<tr>
<td></td>
<td>EP Arrival Time</td>
</tr>
<tr>
<td></td>
<td>3.65 ns</td>
</tr>
</tbody>
</table>

<table>
<thead>
<tr>
<th>WRF Maximum Proton Energy</th>
</tr>
</thead>
<tbody>
<tr>
<td>Location</td>
</tr>
<tr>
<td>TIM 3 (70°)</td>
</tr>
<tr>
<td>TIM 4 (80°)</td>
</tr>
<tr>
<td>TIM 5 (0°)</td>
</tr>
</tbody>
</table>
### Table G.20. Summary of target, laser and proton data for OMEGA shot 56976.

<table>
<thead>
<tr>
<th>Shot 56976</th>
<th>Target Metrology</th>
<th>Laser Parameters</th>
</tr>
</thead>
<tbody>
<tr>
<td>Shell Diameter</td>
<td>860 μm</td>
<td>OMEGA Energy</td>
</tr>
<tr>
<td>Shell Thick.</td>
<td>38.6 (CD), 1 (CH) μm</td>
<td>OMEGA Pulse Shape</td>
</tr>
<tr>
<td>Cone Material</td>
<td>Au</td>
<td>EP Energy</td>
</tr>
<tr>
<td>Cone Tip Thick.</td>
<td>14 μm</td>
<td>EP Focal Spot</td>
</tr>
<tr>
<td>Cone Tip Dia.</td>
<td>10 μm</td>
<td>EP Intensity</td>
</tr>
<tr>
<td>Cone Length</td>
<td>1.2 mm</td>
<td>EP Pulse Width</td>
</tr>
<tr>
<td></td>
<td></td>
<td>EP Arrival Time</td>
</tr>
</tbody>
</table>

#### WRF Maximum Proton Energy

<table>
<thead>
<tr>
<th>Location</th>
<th>Measurements (MeV)</th>
<th>Average</th>
<th>Std. Dev</th>
</tr>
</thead>
<tbody>
<tr>
<td>TIM 3 (70°)</td>
<td>-</td>
<td>-</td>
<td>-</td>
</tr>
<tr>
<td>TIM 4 (80°)</td>
<td>-</td>
<td>-</td>
<td>-</td>
</tr>
<tr>
<td>TIM 5 (0°)</td>
<td>-</td>
<td>-</td>
<td>-</td>
</tr>
</tbody>
</table>

### Table G.21. Summary of target, laser and proton data for OMEGA shot 59124.

<table>
<thead>
<tr>
<th>Shot 59124</th>
<th>Target Metrology</th>
<th>Laser Parameters</th>
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</thead>
<tbody>
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<td>Shell Diameter</td>
<td>860 μm</td>
<td>OMEGA Energy</td>
</tr>
<tr>
<td>Shell Thick.</td>
<td>38.0 (CD) μm</td>
<td>OMEGA Pulse Shape</td>
</tr>
<tr>
<td>Cone Material</td>
<td>Au</td>
<td>EP Energy</td>
</tr>
<tr>
<td>Cone Tip Thick.</td>
<td>5 μm</td>
<td>EP Focal Spot</td>
</tr>
<tr>
<td>Cone Tip Dia.</td>
<td>10 μm</td>
<td>EP Intensity</td>
</tr>
<tr>
<td>Cone Length</td>
<td>1.2 mm</td>
<td>EP Pulse Width</td>
</tr>
<tr>
<td></td>
<td></td>
<td>EP Arrival Time</td>
</tr>
</tbody>
</table>

#### WRF Maximum Proton Energy

<table>
<thead>
<tr>
<th>Location</th>
<th>Measurements (MeV)</th>
<th>Average</th>
<th>Std. Dev</th>
</tr>
</thead>
<tbody>
<tr>
<td>TIM 3 (70°)</td>
<td>-</td>
<td>-</td>
<td>-</td>
</tr>
<tr>
<td>TIM 4 (80°)</td>
<td>-</td>
<td>-</td>
<td>-</td>
</tr>
<tr>
<td>TIM 5 (0°)</td>
<td>-</td>
<td>-</td>
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</tr>
</tbody>
</table>
Table G.22. Summary of target, laser and proton data for OMEGA shot 59126.

**Shot 59126**

<table>
<thead>
<tr>
<th>Target Metrology</th>
<th>Laser Parameters</th>
</tr>
</thead>
<tbody>
<tr>
<td>Shell Diameter</td>
<td>OMEGA Energy</td>
</tr>
<tr>
<td>860 μm</td>
<td>18.54 kJ</td>
</tr>
<tr>
<td>Shell Thick.</td>
<td>OMEGA Pulse Shape</td>
</tr>
<tr>
<td>38.6 (CD) μm</td>
<td>LA241701P</td>
</tr>
<tr>
<td>Cone Material</td>
<td>EP Energy</td>
</tr>
<tr>
<td>Au</td>
<td>942 J</td>
</tr>
<tr>
<td>Cone Tip Thick.</td>
<td>EP Focal Spot</td>
</tr>
<tr>
<td>5 μm</td>
<td>27.1 μm</td>
</tr>
<tr>
<td>Cone Tip Dia.</td>
<td>EP Intensity</td>
</tr>
<tr>
<td>10 μm</td>
<td>3.76×10^{18} W-cm^{-2}</td>
</tr>
<tr>
<td>Cone Length</td>
<td>EP Pulse Width</td>
</tr>
<tr>
<td>1.2 mm</td>
<td>8.7 ps</td>
</tr>
<tr>
<td></td>
<td>EP Arrival Time</td>
</tr>
<tr>
<td></td>
<td>- ns</td>
</tr>
</tbody>
</table>

<table>
<thead>
<tr>
<th>WRF Maximum Proton Energy</th>
</tr>
</thead>
<tbody>
<tr>
<td>Location</td>
</tr>
<tr>
<td>TIM 3 (70°)</td>
</tr>
<tr>
<td>TIM 4 (80°)</td>
</tr>
<tr>
<td>TIM 5 (90°)</td>
</tr>
</tbody>
</table>