CHARGED-PARTICLE SPECTROSCOPY:
A NEW WINDOW ON INERTIAL CONFINEMENT FUSION

by

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Abstract
A new charged-particle diagnostic for inertial confinement fusion experiments has been designed,
constructed and fielded on the OMEGA laser system. The spectrometer consists of a 7.6 kilogauss
permanent magnet and CR-39 nuclear track detectors which can probe energies between 0.1 – 30 MeV
and particle yields between $10^8$ – $10^{15}$. A rapid track scanning system that has been developed is used to
analyze between $10^3$ – $10^6$ particles per shot. Highly resolved spectra have been obtained for a wide
range of primary fusion products, such as D-D protons and tritons, D-$^3$He protons and alphas, D-T
alphas, as well as neutron-scattered fuel ions, or “knock-ons”. Benchmark comparisons of the data with
available neutron measurements show reasonable agreement: fusion yields are similar, on the average, as
are fuel ion temperatures measured by Doppler broadening and line ratio techniques in certain regimes.
The spectrum of triton and deuteron knock-ons have enabled fuel $\rho R$’s to be determined, while the
energy downshift of the fusion products has been used to determine ablator $\rho R$. In addition, copious
fluxes of ablator protons have been detected with energies up to 1.4 MeV when laser intensities exceed a
threshold of $1 \times 10^{14}$ W/cm², while fusion products are sometimes measured with accelerations of several
hundred keV. The symmetry of all measurements is assessed by using two nearly identical spectrometers
mounted with orthogonal views of the target.

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This work is the result of a close collaboration between groups at three laboratories: the Laboratory for Laser Energetics (LLE) at the University of Rochester; Lawrence Livermore National Laboratory (LLNL); and the Plasma Science and Fusion Center at the Massachusetts Institute of Technology (MIT). Hopefully the pages to follow will do justice to the efforts of all the individuals who have worked together to ensure success of this project.

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1 Introduction

Based on the title of this work, possible objectives of an introductory chapter might be to describe the broad discipline of charged-particle spectroscopy and its application in various scientific fields; or perhaps to outline the world’s energy needs, future projections of a crisis, and the potential role of controlled fusion reactors; or maybe to establish the historical context of nuclear fusion and the drama played out by personalities such as Eddington, Gamow and Teller. No doubt all three aspects, and many more, are worthy of discussion and important to gaining an understanding of the broader context of this specialty. But choices must be made. What follows is a brief historical description of the scientific origins of nuclear fusion research, from problems concerning the source of stellar energy, to development of the hydrogen bomb, and finally to the beginnings of controlled fusion research. This leads to an overview of the key aspects of the inertial confinement approach to nuclear fusion, ending with an outline of the increasingly important role of nuclear diagnostics in current and future experiments. It is hoped that this will set the stage for the introduction of charged-particle spectroscopy and its function in elucidating the physics of inertial fusion conditions.

1.1 The origins of fusion research

In 1926, Sir Arthur Eddington, Plumian Professor of Astronomy at Cambridge University, published his classic work, *The Internal Constitution of the Stars*, a brilliant exposition on the physics of stellar interiors. Using astrophysical principles still relevant today, he determined the temperature, pressure and density profiles through the center of the sun without any assumptions about the source of solar energy. Many energy sources had been proposed, including chemical burning, gravitational collapse, and meteor showers but none could maintain the sun’s power output for its required several billion year lifetime. Eddington himself made no secret of his guess. Using Einstein’s Special Theory of Relativity that showed the equivalence of mass and energy, he speculated that the 0.8 % mass difference between a helium nucleus and its four component hydrogen nuclei was the origin of the sun’s power. Somehow hydrogen was being converted to helium, releasing the immense energies locked up in nuclear matter. The only problem was that, at the temperatures of 40 million degrees that he himself had calculated for the center of the sun¹, classical physics claimed that the velocities of hydrogen nuclei were not nearly enough for them to overcome their mutual repulsion. In defiant tone Eddington stood his ground: “We do not argue with the critic who urges that the stars are not hot enough for this process; we tell him to go and find a hotter place.”¹[1]. References to Hades aside, it was ultimately not necessary to find a hotter place. While it was the Theory of Relativity that set Eddington on the right trail, it would be the other great revolution in 20th century physics, Quantum Mechanics, that provided the final answer.

¹ Compare this to the modern accepted value of 15 million degrees.
In a chronological coincidence typical in scientific discovery, the year 1926 marked the birth of Quantum Mechanics. With these powerful new weapons of theory, one of the problems tackled was the transmutation of atoms, in particular radioactivity. George Gamow established that there was a finite probability for alpha particles in heavy nuclei to tunnel through the potential well created by the strong nuclear force and thereby escape the confines of the nucleus. In 1929, by applying the tunneling process in reverse, Atkinson and Houtermans [2] showed that hydrogen nuclei could react with heavier nuclei at much lower temperatures than predicted by classical physics. In their paper, which may be regarded as the beginning of nuclear fusion research, they laid the groundwork for the theory of thermonuclear reactions. Atkinson and Houtermans had proved what Eddington could only surmise – nuclear fusion was responsible for energy production in the stars.

Experimental demonstration of fusion on earth was achieved in 1932 at the Cavendish laboratory when Cockcroft and Walton, utilizing the particle accelerator that bears their name, achieved the first artificial nuclear reactions by fusing hydrogen and lithium-7. The fact that nuclear fusion could now not only be achieved inside the sun, but also inside the Cavendish, was surely not lost on them and perhaps at this time the first serious talk was made of tapping fusion as a terrestrial energy source. Whatever was said though, it was enough for the laboratory director, Ernest Rutherford, to openly dismiss the possibility as “moonshine”[3].

Thoughts about thermonuclear fusion thus remained a theoretical exercise for many more years, confined to the study of stellar interiors. Although Atkinson and Houtermans had showed nuclear fusion to be theoretically possible, the precise sequence of reactions inside the sun had yet to be determined. However, a flurry of activity in 1938, characteristically inspired by Gamow, led to a description of the carbon cycle (Bethe and Von Weizsäcker) and the proton-proton chain (Bethe and Critchfield) – the two most common reaction sequences. Then in 1939 came the discovery of fission and the onset of World War II and it was not long before serious speculation began on a terrestrial application of nuclear fusion – the hydrogen bomb.

To make any use of fusion as an energy source, huge amounts of energy must first be consumed to heat the nuclear fuel to temperatures high enough for a substantial number of fusion reactions to occur. Under the right conditions, this ignites a self-propagating nuclear burn that can release huge amounts of energy

---

2 The notion that a particle accelerator can be used as a fusion energy device is a common misconception. The fundamental difficulty is that the probability for a fusion reaction between accelerated ions is so small that most of the ions are lost or brought to rest without having reacted. These lost ions carry away much greater energy than can be replaced by the occasional fusion reaction.
extremely rapidly – the necessary requirements for a bomb. However, before the advent of the fission bomb, it was unclear how to generate the temperatures and densities necessary to spark thermonuclear ignition. Although it would take until 1945 to completely solve the problems surrounding a practical fission bomb, as early as 1941, Enrico Fermi was already speculating with Edward Teller about the possibility of igniting deuterium using heat from a fission chain reaction [4].

During the summer of 1942, Robert Oppenheimer gathered together in Berkeley a group of physicists, including Bethe and Teller, whose job was to throw light on the actual design of a fission bomb [5]. It did not take long for the group to become involved in calculations, initiated by Teller, to assess the viability of igniting deuterium. Because fusion fuel is not constrained by a critical mass, like uranium or plutonium, there is no theoretical limit to the size of the bomb, and the resulting “limitless” explosion seemed particularly attractive. When further study showed that pure deuterium fuel would not burn fast enough to reach ignition, Konopinski suggested adding tritium. The deuterium-tritium reaction was known, or perhaps established at the time, to proceed more rapidly than that of pure deuterium [4]. By the end of the summer it was apparently clear to the group that, barring unforeseen engineering obstacles, a hydrogen bomb was possible. It is interesting to note that experimental measurements of the deuterium-deuterium and deuterium-tritium cross-sections were not made until 1945, eventually appearing in the open literature in 1949 [6, 7]. Short on experimental measurements, the theoretical nuclear physicists of the day had to make educated predictions about nuclear cross-sections. The guesses were good enough though that even today the first controlled fusion reactor scenarios utilize deuterium and tritium.

Following the war, the first concerted efforts were made to harness nuclear fusion energy in a controlled manner. The most apparent method seemed to be by confining a hot plasma by magnetic fields, as was done by Sir George Thomson in London (1947) and William Allis in Oxford (1948) [3]. Although discussions about possible magnetic geometries took place between Allis and Fermi, Teller and others, the American effort appeared not to start until 1951 when the toroidal pinch was selected for confinement study in Los Alamos. Lyman Spitzer, unaware of other fusion work, built his stellarator in Princeton in 1952. The magnetic confinement of plasmas today forms the mainstay of worldwide controlled fusion research.

During this period, hydrogen bomb development continued and the first ignited thermonuclear reaction exploded over Eniwetok Island on November 1st 1952 – just at the time Teller was setting up a new weapons laboratory at Livermore. By this time, the fundamental ideas of inertial confinement fusion (ICF) were already in place. This is no surprise since the basic concept of ICF is identical to that of the hydrogen bomb. The difference between controlled ICF and uncontrolled bombs is merely one of scale:
if the explosions are small enough they are considered controlled. By 1958, John Nuckolls at Livermore had already made calculations for a fusion power plant that would use a steady succession of small-scale hydrogen bombs; however, he realized that in order for such an idea to be commercially viable, something other fission bombs had to be used to drive the implosions [8].

When the laser was invented in 1960, it was immediately recognized by the Livermore scientists as the appropriate implosion driver. Detailed computer simulations were performed on laser compressed DT fuel targets and, with the construction of a Ruby laser in Livermore in 1962, the ICF program was officially born. The first openly published description of this work and the ideas of ICF finally appeared in 1972 in a now classic paper by Nuckolls, Wood, Thiessen, and Zimmerman [9].

Such a military heritage still plays a key role in ICF today. In particular, the U.S. Government’s stockpile stewardship program, whose purpose is to maintain the safety, reliability, and performance of the nuclear arsenal in the absence of underground testing, provides the main backing for the new National Ignition Facility – a laser fusion laboratory due for completion in 2003, that, if successful, will provide the first demonstration of controlled fusion ignition. However, while defense applications are paramount, there is strong awareness of the contributions such experiments can make to other scientific disciplines, in particular the studies of astrophysical phenomena. For example, unprecedented tests of solar and stellar models as well as under-studied fusion reactions will be possible under NIF conditions. Perhaps it is only appropriate then that a discipline whose origins lie in the astrophysical questions of the past can begin to help solve the related questions of today.

1.2 Nuclear fusion by inertial confinement

1.2.1 Physical conditions

The ultimate intent of inertial confinement fusion is to generate net fusion energy by setting off a steady stream of controlled, thermonuclear explosions. To produce each of these miniature explosions, approximately 5 mg of D-T fuel must first be imploded to a density of ~ 400 g/cm³ and a temperature near 20 keV, whereupon fusion reactions can proceed at a high rate. Rapid expansion quenches these conditions in tens of picoseconds, but, before this happens, enough fusion energy has been generated to compensate for the energy used in driving the implosion. It is the inertia of the imploded assembly that “confines” the fuel for this brief burn period. To provide some justification for such conditions, this section will follow a few simple calculations that establish the necessary density and time scales as well as the energy needs. It will be shown how imploding spherical targets by laser ablation can begin to meet these requirements. Two points should be kept in mind: Firstly, these calculations are zero-dimensional
and so should only be regarded as order-of-magnitude estimates; secondly, the physical conditions derived here are those necessary for ignition and high gain which have yet to be achieved on current experiments.

The fuel conditions just described, that are required to achieve a large net energy gain, can be derived by analyzing the energy balance of an ignited, inertially-confined fuel mass [10]. In order to accomplish high gain it is crucial to achieve ignition – the point where enough of the generated fusion energy heats the surrounding fuel to initiate further fusion reactions, thereby generating a self-propagating burn wave. This reduces the amount of input energy required since now only a small part of the fuel, the "spark", needs to be heated from the outset. In deuterium-tritium (D-T) fuel, it is only the 3.5-MeV alphas that can heat the fuel directly. The amount of alpha energy generated per unit mass of fuel, \( E_a \), is given by

\[
E_a = \frac{Q_a}{m_D + m_T} f_B
\]  

(1.1)

where \( Q_a \) is the energy per alpha particle (3.52 MeV), \( m_D \) and \( m_T \) are the deuteron and triton masses respectively, and \( f_B \) is the fraction of fuel consumed by the fusion reactions. To calculate \( f_B \), consider the reaction rate equation, with equal numbers of D and T ions:

\[
\frac{dn}{dt} = -\frac{n^2}{2} \langle \sigma v \rangle
\]  

(1.2)

where \( \langle \sigma v \rangle \) is the reaction rate parameter for the D-T reaction and \( n \) is the total number density of ions. Assuming that \( \langle \sigma v \rangle \) stays constant over the duration of the burn, this equation can be integrated to give \( n \) as a function of the burn time \( \tau \) and the initial number density \( n_0 \). Since the burn fraction is defined as \( f_B = (n_0 - n)/n_0 \), this can be used to show that

\[
f_B = \frac{n_0}{n_0 + 2/\langle \sigma v \rangle \tau}
\]

(1.3)

Roughly speaking, the burn time \( \tau \) is the time it takes for a rarefaction wave, which quenches the burn, to propagate from the outside of the fuel to its center. More precisely, it is the mass-weighted average time during which the fuel ions, at the various radii, are able to react before passage of the converging rarefaction wave. Performing this mass-average over the fuel gives \( \tau = \frac{\rho R}{c_s} \) where \( R \) is the radius of the fuel region and \( c_s \) is the sound speed in the hot plasma\(^3\). Using the fact that the initial mass density, \( \rho = <m_i>n_0 \), where \( <m_i> \) is the average ion mass, Eq. (1.3) can be re-written as

\[
f_B = \frac{\rho R}{\rho R + 8c_s <m_i>/\langle \sigma v \rangle} = \frac{\rho R}{\rho R + 6(g/cm^2)}
\]

(1.4)

\(^3\) Note that some calculations use \( \tau = R/3c_s \).
The constant $6 \text{ g/cm}^2$ appears because, in the temperature range $20 - 40 \text{ keV}$, the term $c_s/\langle \sigma v \rangle$ is almost independent of temperature, where $c_s = \sqrt{2T/\langle m_i \rangle}$, assuming equal electron and ion temperatures. The term $\rho R$ is defined as the fuel areal density.

Now, at ignition, the fusion energy deposited by the alphas during the burn time – the “confinement time” of this inertial plasma – must be equal to the thermal energy of the plasma, where the plasma thermal energy per unit mass is $3T/\langle m_i \rangle$. This assumes that other energy losses, in particular through bremsstrahlung and electron conduction, are small – a fact confirmed by more detailed calculations of ignition conditions [11]. Thus, at ignition, using Eqs. (1.1) and (1.4),

$$\frac{Q_a}{2\langle m_i \rangle} = \frac{\rho R}{\rho R + 6} = \frac{3T}{\langle m_i \rangle}$$

(1.5)

Solving for $\rho R$ at a temperature of $20 \text{ keV}$ shows that ignition is achieved when $\rho R = 0.2 \text{ g/cm}^2$. Note this equation assumes that all the alpha energy is deposited in the fuel, or, more specifically, that the range of alphas is small compared to the fuel radius. It is a coincidence that, around $10 \text{ keV}$ at densities of $\sim 100 \text{ g/cm}^3$, the alpha particles have a range $\rho \lambda_a \approx 0.3 \text{ g/cm}^2$ [12]; therefore, ignition occurs when the alpha particles are just being completely ranged out in the fuel. On the other hand, it is no coincidence that a $\rho R$ of $0.2 \text{ g/cm}^2$ corresponds to an $n \tau$ of $1.7 \times 10^{14} \text{ s/cm}^3$, similar to the Lawson criterion.

Although $\rho R$’s of $0.3 \text{ g/cm}^2$ are sufficient for ignition, the burn fraction given by Eq. (1.4) at these areal densities is only 5%. ICF targets need much higher burn fractions in order to achieve the necessary gain to compensate for the inefficiencies accumulated by imploding a target. It is generally accepted [13] that $\rho R$’s of $3 \text{ g/cm}^2$ are attainable and would result in high gain targets with a burn fraction of 33%. Unlike in magnetic fusion, just reaching ignition in ICF is insufficient for generating high gain.

The above discussion of ignition and burn sets a lower limit on the fuel areal density but says nothing about the required absolute density, mass of the target, or the attainable fusion output energy and necessary input energy. Using the $\rho R$ limit though, these parameters can be estimated. To start with, consider the mass of a spherical target, which, for a fixed $\rho R$, is proportional to $1/\rho^2$.

$$M = \frac{4\pi}{3} \left( \frac{\rho R}{\rho^2} \right)^{3/2} \sim \frac{1}{\rho^2}$$

(1.6)

If the target is kept at a normal liquid density of $0.21 \text{ g/cm}^3$, then to achieve the $\rho R$ of $3 \text{ g/cm}^2$ necessary for high gain, 2.6 kg of D-T would be required. At a burn fraction of 33%, with each D-T reaction releasing 17.6 MeV of fusion energy, this much fuel would produce $3 \times 10^{14} \text{ J}$ or 72 kt of TNT equivalent – hardly a miniature explosion. Realistically then, to achieve sufficient gain with manageable energy
yields, smaller targets must be compressed to high densities. Now, in a realistic implosion, the fuel region at maximum density does not form a uniform sphere but rather a spherical shell around the high temperature, and lower density, hot spot which initiates the burn. If this high density shell has a thickness $R$ and an outer radius $2R$, then at densities of 400 g/cm$^3$, the required fuel mass is 5 mg. This mass would release 600 MJ of fusion energy. The burn time under these conditions, given by $\frac{1}{4} \frac{R}{c_s}$, is $\sim 20$ ps.

The total energy required to compress a target to 400 g/cm$^3$ and heat it to 20 keV depends heavily upon whether or not the target is heated first or compressed first. Matter at low temperatures is easiest to compress since its pressure is lower and therefore the $PdV$ work required is smaller. The most energy efficient technique, and in fact the one that is crucial for success in the near future, is thus to compress the entire target first, at the lowest possible temperature, and then to heat a small central hot spot to several keV, igniting the fusion burn. Any pre-heating of the fuel before the necessary densities are reached will reduce the achievable densities, $\rho R$’s, and subsequent energy gains. Starting with cryogenic, solid D-T, the relevant equation-of-state for much of the compression is that of a partially degenerate gas. For simplicity, the fully degenerate (non-relativistic) Fermi gas equation-of-state will be used. Here, the pressure, $P$, is given by [14]

$$P = \frac{2}{5} n_e \varepsilon_F = \frac{h^2}{20m_e} \left( \frac{3}{\pi} \right)^{\frac{2}{3}} \frac{5}{n_e^5}$$

(1.7)

where $n_e$ is the electron number density (equivalent to the ion density in the case of D-T), $\varepsilon_F$ is the Fermi energy, $h$ is Planck’s constant and $m_e$ is the electron mass. Since this is an ideal Fermi gas, the pressure is temperature independent; however, finite temperature corrections can be added which increase the pressure at any given density. Now for any type of gas, $P = 2/3 \rho \varepsilon$, where $\varepsilon$ is the energy per unit mass. Using this with Eq. (1.7) shows that to compress 5 mg of D-T fuel to 400 g/cm$^3$ takes 90 kJ. If the central hot spot is only 2% of the total mass, then to heat it to 10 keV requires 60 kJ. Thus, by heating only a small fraction of the fuel, the compressive work and heating energy are comparable and the total energy required by the target is 150 kJ. If the efficiency of delivering this energy to the target is 5%, then the energy required of any target driver is 3 MJ and the total gain achieved is 200.

In order to determine which type of drivers are suitable for imploding a D-T target to these conditions, it is useful to estimate the pressure inside the target during peak compression. At 400 g/cm$^3$, $n = 5 \times 10^{26}$ cm$^{-3}$, and, using Eq. (1.7), $P = 7 \times 10^{11}$ atm. These densities and pressures are of the same order as those inside the cores of hydrogen burning stars such as the sun. More to the point though, any driver of the target implosion has to be able to generate these kinds of pressures. Since chemical explosives are limited to $10^6$ atm, before the advent of the laser the only way to achieve this was with fission bombs. Laser-driven ablation of matter, however, can produce $\sim 10^8$ atm on the surface of a solid [15], and, using
a spherical implosion with carefully tailored convergence, the required pressures above $10^{10}$ atm can be accomplished, *inside the compressed target*. It is hoped that in the future, ablation driven by heavy ions can also perform this task as these drivers have the potential to be much more efficient than lasers.

The technique of ICF then is to implode a spherical target by laser ablation of its surface. By gradually ramping up the laser pulse to avoid propagating large shock waves through the target (which would preheat the fuel), the fuel needs to be almost isentropically compressed to several hundred times liquid densities. At the point of maximum compression a central hot spot must be ignited by the convergence of several converging shock waves. A typical ICF target, shown in Figure 1.1, consists of a spherical fuel region covered by a low Z ablator. The ablator produces the rocket blow-off and is mostly consumed by the time peak compression is reached. Future high gain targets will consist of a spherical shell of solid DT fuel, but current targets are mostly filled with D-T or D-D gas at several atmospheres.

Although $\rho R$'s of a few g/cm$^2$ and densities of several hundred g/cm$^3$ are necessary for high gain, the conditions currently attainable are far more modest. Present day fuel $\rho R$'s are generally only around a few to tens of mg/cm$^2$, with densities of a few g/cm$^3$ (achieved with the simpler gas-filled, rather than cryogenic targets) although fusion temperatures of $5 - 20$ keV are readily achieved. In order to close the gap in fuel compression though, a number of challenges must be tackled.

![Figure 1.1: Schematic of an ICF target showing the ablator and fuel regions. In the cryogenic targets proposed for high gain experiments, the solid fuel is arranged in a hollow shell surrounded a gas interior. Highest compressions are achieved with targets whose shells are thin compared with their radii. However these targets are more unstable to the Rayleigh-Taylor instability. Capsule design then must provide a balance between these competing effects.](image-url)
1.2.2 Primary challenges for ICF today

In the original 1972 paper by Nuckolls et al. [9], it was concluded that a laser energy of 1 kJ was sufficient to achieve ignition – a goal that would have made ignition possible with current laser energies of 30 – 40 kJ. Since then though, a better understanding of the Rayleigh-Taylor instability and laser-plasma instabilities has revised this energy requirement to greater than 1 MJ. These effects, whose severity was considerably underestimated by Nuckolls, are still the main concerns of ICF today.

The Rayleigh-Taylor (RT) instability occurs at the interface between fluids of different densities when the acceleration is in the same direction as the density gradient across the interface [16]. During compression of an ICF target, this situation occurs twice: first when the hot, low density plasma corona accelerates the denser ablator shell, and second when the lower density, high pressure fuel core decelerates the converging shell. The effect of the RT instability is to cause unstable growth of ripples at the interface between the high and low density regions. Ripples evolve into fingers penetrating into the rarefied fluid and bubbles rising through the dense fluid, mixing the two regions and causing the shell to breakup, compromising the achievable target compression [17, 18]. Much effort has been devoted to studying this hydrodynamic behavior, primarily using 2-D codes and planar foil experiments. It is difficult to diagnose the implosion symmetry of spherical targets with high convergence since the standard x-ray imaging techniques readily image only the outside of the capsule – the core generally being optically thick to soft x rays – although a few attempts have been made [19]. The degree of fuel mix can be estimated indirectly by measuring \( \rho R \) and comparing the result, and the measured fusion yield, to the predictions of 1-D simulations. Since 1-D compression produces the highest densities, a measured shortfall can be attributed to a certain amount of mix. In light of the importance of uniform implosions, diagnosing asymmetries in spherical geometry will provide crucial feedback on the behavior of hydrodynamic instabilities during convergence.

Growth of the RT instability originates with initial seed perturbations such as those arising from non-uniformities in laser irradiation. The main drawback with directly illuminating the ablator surface is that the focal plane of a laser beam potentially contains large and uncontrollable intensity variations. These are caused by interference of the phase distortions from numerous optical elements and, untreated, are the source for deleteriously large perturbations on the target implosion. One method to circumvent these problems is the indirect-drive approach to ICF, developed in 1975. Instead of directly irradiating the target surface, indirect drive first illuminates and ionizes the inner walls of a high Z enclosure, or hohlraum, surrounding the capsule. The resulting x rays uniformly bathe the target and drive the implosion. Despite the loss in coupling efficiency from the laser to DT fuel the improved uniformity

\[ \text{Nuckolls' calculation required that } < 0.05 \text{ mg of fuel be brought to } \rho R \sim 0.3 \text{ g/cm}^2. \]
made indirect drive approach more promising than direct drive for many years. The advent of beam smoothing techniques, in particular Induced Spatial Incoherence (ISI) and Smoothing by Spectral Dispersion (SSD), in the mid-1980’s enabled direct drive to make a comeback. The OMEGA laser system at the University of Rochester is the primary direct drive facility in the United States, while the NOVA laser at Lawrence Livermore National Laboratory is the focus of indirect drive (although both laser systems can be reconfigured for either technique). The experiments described in this thesis were all performed on direct drive experiments at OMEGA.

Although Rayleigh-Taylor mix and the issues surrounding drive symmetry are key challenges for ICF today, much effort is also being channeled towards an improved understanding of laser-plasma instabilities. The severity of these plasma collective effects has necessitated placing an upper limit on the applied laser intensity. A simple model for capsule implosion, as described by Lindl [10], predicts that the relationship between the required driver energy, $E_D$, and the laser intensity, $I$, is $E_D \sim I^{2.9}$. This means that to minimize the required driver energy, it is advantageous to use as high intensities as possible – in 1972, Nuckolls assumed that intensities of $10^{17}$ W/cm$^2$ could be used. By the late 1970’s however, it became evident that at intensities much greater than $10^{15}$ W/cm$^2$, less than half of the laser energy is absorbed by the target, much of it being lost to stimulated Raman scattering (SRS) and stimulated Brillouin scattering (SBS). More importantly, a substantial fraction of the absorbed energy is carried by suprathermal electrons, accelerated by plasma collective effects such as SRS and the two-plasmon instability. These electrons prematurely heat the fuel and prevent high compressions from being achieved. To mitigate these effects, laser intensities are generally kept in the range $10^{14} - 10^{15}$ W/cm$^2$. Additionally, the laser frequency has a profound impact on the severity of laser-plasma instabilities. Shorter wavelength light penetrates to higher densities where more of the laser energy is converted to plasma thermal energy by collisional absorption, rather than to suprathermal electrons by instabilities in the less dense corona. For this reason, the OMEGA and NOVA lasers both work with frequency-tripled light, at 0.35 µm, from Neodymium-Glass lasers.

Since the late 1980’s there has been reasonable confidence that the combination of reduced laser intensities and shorter wavelength lasers would be sufficient to provide excellent laser plasma coupling and tolerable hot electron generation [20]. However, because of the multitude of constraints on laser energy, it is highly desirable to maximize the available intensity and wavelength parameter space for ICF target design. This means that it is crucial to have an accurate understanding of the thresholds for onset of these different instabilities. For this reason, there is still considerable interest in the various plasma collective effects, each of which will be explained in more detail in Chapter 4. Of particular note is the significant reduction in instability levels due to laser beam smoothing techniques such as ISI and SSD.
Removing local hot spots in the laser intensity enables higher average laser intensities to be used. Understanding the benefits of laser beam smoothing is a major thrust of laser plasma studies.

1.2.3 The OMEGA experiment

At the present time, the OMEGA laser system at the Laboratory for Laser Energetics, University of Rochester, is the flagship of the U.S. inertial fusion program. OMEGA is a neodymium-doped phosphate glass laser capable of delivering up to 30 kJ of frequency-tripled, 0.35 µm light [22]. This energy can be distributed into 60 beams to irradiate a spherical target with a root-mean-square uniformity of ~ 5%. The uniformity of individual beams is achieved through the implementation of two-dimensional SSD. Each beam can be individually focused with a pointing accuracy of ±16 µm and a timing resolution on target of ±10 ps. A variety of pulse shapes are possible, with widths ranging from 115 ps to a few nanoseconds. The laser is typically operated at a repetition rate of about 1 shot per hour. The OMEGA target chamber is a 3.3 m diameter aluminum “soccer ball” with 60 beam ports and 32 diagnostic ports. A vacuum of about 10⁻⁵ Torr is maintained during experiments.

The other main ICF experiment in the U.S., the 40 kJ, 10 beam NOVA laser is reducing its operations as focus shifts towards construction of the new National Ignition Facility (NIF). The NIF, due for completion in the year 2003 is designed to achieve ignition by 2005. This facility will comprise a 1.8 MJ, 192-beam laser. The NIF design aims to achieve a target gain of 10 – 20, producing approximately 10¹⁹ D-T neutrons. In comparison, the maximum D-T yield achieved on OMEGA so far has been 10¹⁴.

1.3 ICF diagnostics & charged-particle spectroscopy

In order to probe the wide range of physical conditions that characterize a laser-driven ICF target, a large variety of diagnostics are utilized to detect plasma, optical, x-ray, neutron and charged-particle emissions [22]. The optical instruments, operating between the ultra-violet and near-infra red, are mostly used to examine laser plasma interactions, while calorimeters measure the energy of plasma blow-off as a gauge of laser absorption and energy transfer in the target. X-ray detectors probably form the most diverse group of diagnostics, ranging from microscopes and pinhole cameras to spectrographs and streak cameras. Together they can provide space and time resolved images of the implosion, as well as measurements of core temperatures and ablactor ρR’s. With the achievement of higher ρR’s, which make imploded targets much more optically thick to standard x-ray diagnostics, there is an increasing need to turn towards more penetrating radiations to measure core conditions. For this reason, nuclear diagnostics, in particular neutron and charged-particle (specifically charged-fusion-product) detectors, will play an important role in the current and future generation of experiments [23].
At present, neutron detectors are by far the most advanced of all ICF nuclear diagnostic technologies and many of them are part of the standard suite of routine ICF diagnostics [24]. Time-of-flight measurements using plastic scintillators have been used in both single-hit and current mode to measure ion temperatures and fuel $\rho R$'s. Placing the scintillator close enough to the target to reduce Doppler broadening enables the fusion burn history to be determined, while integrated D-T and D-D yields are extracted using the activation of copper and indium.

In contrast, charged-fusion-product diagnostic methods have been used sparingly, with the knock-on $\rho R$ measurement technique perhaps being the only one regarded as "routine" [25, 26]. Absent of charge, neutrons escape virtually untouched from the target and arrive at detectors with a spectrum that reflects purely the birth conditions in the core. On the other hand, charged particles are affected by ranging through the fuel and ablator material and may interact with fields in the surrounding corona. This means that charged-particles contain inherently more information than neutrons. The use of several different types of ions aids in extracting some of this information: the different particles are effectively multiple measurements of the various unknowns.

The purpose of this thesis is to examine the wealth of detailed information that can be extracted from charged particles using highly resolved spectral measurements. A spectrometer has been designed, built and fielded and a rich array of spectra, from both charged fusion products and accelerated coronal ions, has been observed. Some of the results can be compared directly to neutron measurements, giving baseline confidence in the instrument, while others provide the first evidence of new and unexpected physical effects.
2 The charged-particle spectrum

The wealth of information contained within the charged-particle spectrum from ICF experiments is due in large part to the variety of different reactions and particle species that contribute. Between ~1 and 30 MeV, the spectrum consists of a diverse range of discrete lines and continua made up of protons, deuterons, tritons, and alphas generated by nuclear reactions. Discrete lines are produced by fusion reactions while continuous spectra up to a maximum, endpoint energy are caused by the elastic scattering of ions by fusion neutrons. These processes are usually divided into primary reactions, secondary reactions (fusion of a primary product), and "knock-on" interactions (neutron scattering). In addition, tertiary reactions utilizing all three processes in sequence can produce particles with the highest energies. The low energy particles below ~1 MeV, on the other hand, have an entirely different origin, consisting mostly of protons accelerated by laser plasma phenomena outside the target. An overview of all these reactions important to charged-particle spectroscopy will be presented, with brief remarks about how they may be used to diagnose certain target parameters, though discussion of the latter will be the subject of the chapter to follow.

In general, the most copious particles above about 1 MeV are the products of primary fusion reactions. The primary fusion products appear as discrete lines at energies predominantly determined by the nuclear energetics of each reaction. Such lines are discrete (with a narrow width) because the thermal energy of the fuel ions, at a few keV, is small compared to the energy of the products. Depending on the type of fuel used, these products can include 3.0-MeV protons and 1.0-MeV tritons from D-D reactions, 3.7-MeV alphas and 14.7-MeV protons from D-³He reactions, and 3.5-MeV alphas from D-T reactions. Detection of these particles is an important benchmark in charged-particle spectroscopy as their discrete energies are instrumental in verifying the absolute energy calibration of any spectrometer. In addition, the absolute yield value determined by the instrument can be tested by comparing available neutron measurements, at least for the D-T and D-D reactions. Comparing neutron temperature measurements with those from the Doppler width of particle lines and the ratio of fusion products also provides a good benchmark test for the charged-particle technique.

\[
\begin{align*}
\text{Primary fusion reactions} & : \\
D + D & \rightarrow T (1.01 \text{ MeV}) + p (3.02 \text{ MeV}) \\
& \rightarrow n (2.45 \text{ MeV}) + ^{3}\text{He} (0.82 \text{ MeV}) \\
D + T & \rightarrow \alpha (3.52 \text{ MeV}) + n (14.07 \text{ MeV}) \\
D + ^{3}\text{He} & \rightarrow \alpha (3.67 \text{ MeV}) + p (14.68 \text{ MeV})
\end{align*}
\]

Knock-on particles are generated by elastic scattering of 14.1-MeV D-T neutrons with the surrounding fuel and ablator material. For instance, neutrons scattering off deuterons and tritons in the fuel generate
deuterons up to 12.5 MeV and tritons up to 10.6 MeV [25, 27, 28], while neutron collisions with hydrogen (in the ablator or fuel) can produce protons up to 14.1 MeV [29]. The endpoint energy of these particles is determined purely by reaction kinematics while the precise shape of the spectrum of a given ion is governed by the neutron differential energy cross-section. The yield of knock-ons is proportional to the fuel $\rho R$.

\[
D + T \rightarrow \alpha (3.5 \text{ MeV}) + n (14.1 \text{ MeV}) \quad \text{(step 1)}
\]

14.1-MeV neutron knock-ons

\[
n (14.1 \text{ MeV}) + p \rightarrow n' + p (\leq 14.1 \text{ MeV}) \quad \text{(step 2a)}
\]

\[
n (14.1 \text{ MeV}) + D \rightarrow n' + D (\leq 12.5 \text{ MeV}) \quad \text{(step 2b)}
\]

\[
n (14.1 \text{ MeV}) + T \rightarrow n' + T (\leq 10.6 \text{ MeV}) \quad \text{(step 2c)}
\]

Secondary fusion particles are the result of two sequential fusion reactions. The first step is the D-D reaction, producing tritons and $^3$He ions. The 1.0-MeV tritons can then fuse with thermal deuterons to produce alphas between 1.4 and 6.7 MeV, while the 0.8-MeV $^3$He ions can react with thermal deuterons to generate protons between 12.5 and 17.4 MeV. The secondary particles have a wider band of energies than the primaries as a result of substantial kinematic broadening and, for sufficiently tenuous plasmas, their yields are proportional to the fuel $\rho R$. Although secondary particles have been examined in the past [30, 33] (with the neutron secondary method being a common way of measuring $\rho R$ [31, 32, 33]) the technique is not the focus of this experiment and thus will not be discussed further.

\[
D + D \rightarrow T (1.01 \text{ MeV}) + p (3.02 \text{ MeV}) \quad \text{(step 1)}
\]

Secondary fusion reactions

\[
T (1.01 \text{ MeV}) + D \rightarrow \alpha (6.7-1.4 \text{ MeV}) + n (11.8-17.1 \text{ MeV}) \quad \text{(step 2)}
\]

\[
D + D \rightarrow n (2.45 \text{ MeV}) + ^3\text{He} (0.82 \text{ MeV}) \quad \text{(step 1)}
\]

\[
^3\text{He} (0.82 \text{ MeV}) + D \rightarrow \alpha (6.6-1.7 \text{ MeV}) + p (12.5-17.4 \text{ MeV}) \quad \text{(step 2)}
\]

Tertiary reactions in fuel containing deuterium, tritium and $^3$He can generate the highest energy particles of interest [34]. This process involves a sequence of all three mechanisms described above: initially, primary D-T reactions produce 14.1-MeV neutrons; some of these neutrons go on to generate elastically scattered deuterons up to 12.5 MeV; and then a fraction of these high energy deuterons fuse with $^3$He fuel ions, producing a spectrum of protons up to 30.8 MeV. This spectrum is peaked near the maximum energy, mirroring the endpoint peak in the knock-on deuteron spectrum. The yield of these particles is approximately proportional to $(\rho R)^2$. A neutron producing tertiary counterpart is described in reference [35]. Owing to their low yields, tertiary fusion products have yet to be detected and thus will not be discussed further; however, the necessity, within the next several years, for diagnosing the high areal-
densities of targets approaching ignition will mean that tertiary processes, with their associated highly penetrating particles, will become crucial diagnostic techniques.

\[
\begin{align*}
\text{30.8-MeV tertiary reaction chain} & \quad D + T \rightarrow \alpha (3.5 \text{ MeV}) + n (14.1 \text{ MeV}) \\
& \quad \text{(step 1)} \\
& \quad n (14.1 \text{ MeV}) + D \rightarrow n' + D (\leq 12.5 \text{ MeV}) \\
& \quad \text{(step 2)} \\
& \quad D (12.5 \text{ MeV}) + ^3\text{He} \rightarrow \alpha + p (\leq 30.8 \text{ MeV}) \\
& \quad \text{(step 3)}
\end{align*}
\]

The spectra described here are those produced purely by the nuclear reactions themselves, all of which occur near the center of the imploded target. The path of any escaping particle must of course involve passage out through the dense fuel and ablator material, a process that may alter the nascent particle spectrum. In particular, the energy losses of an emerging particle can be used to determine the total \(\rho R\) (ablator and fuel) of the target. At high enough \(\rho R\)'s though, these particles can be completely ranged out. For example, in an ablator consisting of carbon and hydrogen at 1 keV, the range of 3-MeV D-D protons, 14.7-MeV D-\(^3\)He protons and 30.8-MeV tertiary protons is 50 mg/cm\(^2\), 300 mg/cm\(^2\) and 1000 mg/cm\(^2\) respectively (Figure 2.1). For this reason it is important to utilize a diverse array of particles that have a wide variety of ranges spanning and exceeding the different target conditions.

![Figure 2.1: The range of protons at important energies in a CH ablator plasma for a variety of ablator temperatures. The 3.0 MeV and 14.7 MeV protons are useful for probing current experiments, while the highly penetrating 30.8 MeV proton will be useful in future, high-density conditions. Ion densities of \(10^{24} \text{ cm}^{-3}\) were assumed in these calculations.](image)

The charged-particle spectrum below \(\sim 1 \text{ MeV}\) is dominated by ions, from the ablator plasma, that have been accelerated by electric fields in the target corona. Under certain conditions, the laser can drive plasma instabilities which cause ejection of suprathermal electrons from the target. This generates radial
electric fields that accelerate the ions up to several hundred keV. If nuclear reaction products are produced during the presence of such fields, the spectra of these products can be upshifted slightly.

The division of the charged-particle spectrum into two regimes – that above and below ~ 1 MeV – is roughly a division between information about the core and ablator regions comprising the target, and about the laser-plasma interaction region outside the target. Above ~ 1 MeV, the spectrum is comprised almost exclusively of particles born of nuclear reactions within the target, such as primary fusion reactions and knock-on interactions. This spectrum bears information predominantly about core and ablator conditions. Chapter 3 will describe quantitatively how important target parameters such as fuel \( \rho R \), ion temperature, and ablator \( \rho R \) may be determined using this spectrum. On the other hand, below ~ 1 MeV, the particles are comprised mostly of ions accelerated from the ablator by fields in the laser-plasma region. Chapter 4 will outline possible mechanisms responsible for these particle accelerations and how they can be used to diagnose hot electron temperatures.
3 Diagnosis of core and ablator conditions

The diagnosis of direct-drive ICF experiments can roughly be divided into two categories: the study of laser performance and the study of target performance [36]. Characterization of laser performance involves issues such as beam spot size, positioning, smoothing and synchronization, as well as the diverse array of laser-plasma issues related to the coupling of laser light to the plasma. Study of target performance involves measurement of parameters such as core $\rho R$ and temperature, ablator $\rho R$, and fusion yield. Interestingly, charged-particle spectroscopy can provide information about both categories: the fusion products are useful for diagnosing core and ablator conditions, while accelerated plasma ions characterize the behavior of certain laser-plasma interactions. This chapter will investigate the core and ablator parameters so far accessible by charged-particle spectroscopy.

3.1 Primary fusion yields

Yield is simply the total number of primary fusion reactions of a particular type produced by a shot. For instance, in modern ICF experiments, the D-T yield ranges from $10^9$ to $10^{14}$ while the D-D yield is typically between $10^8$ and $10^{12}$. Future experiments to achieve ignition on the National Ignition Facility (NIF) are projected to have maximum yields of $10^{19}$ D-T reactions [23]. For the maximum yields of $10^{14}$ achievable in current D-T experiments, the ratio of total fusion energy produced to laser incident on the target is of order $10^{-3}$; on the NIF, this is expected to reach $\sim 10$.

Currently yield measurements are all made using neutron diagnostics. The most common approach, which has the largest dynamic range, is based on neutron activation. Typically, the 14.1-MeV D-T neutron yield is determined using copper activation [37] while indium activation is used to measure the yield of 2.5-MeV D-D neutrons [38]. These reactions are shown in Table 3.1.

<table>
<thead>
<tr>
<th>Neutron Energy</th>
<th>Origin</th>
<th>Activation reaction</th>
<th>Threshold Energy</th>
</tr>
</thead>
<tbody>
<tr>
<td>14.1 MeV</td>
<td>$^3$H(d,n)$^4$He</td>
<td>$^{63}$Cu (n,2n) $^{62}$Cu</td>
<td>10.7 MeV</td>
</tr>
<tr>
<td>2.5 MeV</td>
<td>$^2$H(d,n)$^3$He</td>
<td>$^{115}$In (n,n') $^{115m}$In</td>
<td>0.38 MeV</td>
</tr>
</tbody>
</table>

Table 3.1: The neutron activation reactions that are used in determining the yield of 14.1-MeV D-T neutrons and 2.5 -MeV D-D neutrons.

The copper activity is measured by counting the 0.511 MeV gamma-rays from positron annihilation, while the indium activity is determined by measuring the number of 336 keV gamma-rays emitted. The relationship between isomer activity and incident neutron flux can be determined by calibrations [38] or by calculations using the reaction cross sections, self-absorption factors of the material (Cu or In) and detector efficiency for gamma-rays of interest. The effects of neutron scattering within the target chamber environment must also be taken into consideration.
An absolute yield measurement from charged fusion products is achieved by counting the number of fusion products. The D-T yield can be determined by counting the number of alphas generated in the $^3\text{H}(d,n)^4\text{He}$ reactions, while the D-D yield can be measured using the number of protons or tritons produced by the $^3\text{H}(d,p)^3\text{H}$ reaction. The alpha yields provide direct comparison to D-T neutron yields, while proton or triton yields must be related to neutron yields via the branching ratio for $^2\text{H}(d,p)^3\text{H}$ versus $^2\text{H}(d,n)^3\text{He}$ reactions shown in Figure 3.1.

![Graph of D(D,p)T/D(D,n)$^3\text{He}$ Yield Ratio](image)

Figure 3.1: Predicted ratio of proton to neutron yields from the D-D reaction at different temperatures [39]. The curve fit used is within 1% of numerical calculations of $<\sigma>$, integrating from the basic cross-sections; however, the cross-section measurements are only accurate to 5%.

Reference [40] reports that alpha yields measured by CR-39 plastic track detectors are in good agreement with the copper calibration results; on the other hand, reference [41] reports a 30% – 50% deficiency of D-D protons compared to D-D neutrons, with the protons being measured using a magnetic spectrometer and CR-39.

Asymmetries in the particle fluxes, and thus uncertainties in the yield measurements from any single detector, can be estimated by placing multiple detectors around the target chamber.
3.2 Fuel ion temperatures

Two ways of using primary fusion products to measure fuel ion temperature are the method of yield ratios for different fusion reactions and the method of Doppler line widths. These are independent techniques which, used concurrently, provide a valuable consistency check.

3.2.1 Fusion yield ratios

The fusion reaction rate per unit volume, \( R \), is given by

\[
R = n_A n_B \langle \sigma v \rangle
\]

(3.1)

where \( n \) is the particle density, subscripts \( A \) and \( B \) refer to the different particle species, and \( \langle \sigma v \rangle \) is the rate coefficient understood to be an average over particle distributions for both species. Since the energy distribution of particles in thermal equilibrium is fully described by the local temperature, \( \langle \sigma v \rangle \) can be considered to be purely a function of \( T_i \), the ion temperature. The variation of the rate coefficient with temperature for relevant fusion reactions is given in Figure 3.2.

![Fusion rate coefficients](image)

Figure 3.2: Rate coefficients for fusion reactions important to charged-particle spectroscopy [39]. The D-D,p curve represents the rate coefficient of the D-D proton branch only. At temperatures of \( \sim 15 \) keV, the gradients of the D-T and D-D,p coefficients are very similar, meaning that their ratio is almost constant. On the other hand, the D-D,p and D-\(^3\)He reactions always have very different gradients, causing their ratio to change rapidly with temperature; this is useful as a temperature diagnostic.

Measurement of temperature using the ratio of two reaction yields has the advantage of factoring out unknown quantities, common to both lines (density and volume), which contribute to the measured absolute yield. In charged-particle spectroscopy of fusion products, the lines are, for example, due to D-T
\( \alpha \)'s or D-D protons, or perhaps D-\( ^3 \)He protons. In a D-T plasma both D-T and D-D reactions will occur at rates given by

\[
R_a = n_p n_T \langle \sigma v \rangle_{DT} \quad R_p = \frac{n_p^2}{2} \langle \sigma v \rangle_{DD-p}
\]

(3.2)

where \( R_a \) is the production rate of D-T \( \alpha \)'s, \( R_{DD-p} \) is the production rate of D-D protons \( \langle \sigma v \rangle_{DD-P} \) is the rate coefficient for the proton branch of the D-D reaction. For a plasma containing a D:T atomic ratio of 1:1, the \( \alpha \) to proton yield ratio is:

\[
\frac{Y_\alpha}{Y_{DD-P}} = 2 \frac{\langle \sigma v \rangle_{DT}}{\langle \sigma v \rangle_{DD-P}}
\]

(3.3)

Similarly, for plasmas containing a 1:1 atomic ratio of D and \( ^3 \)He, the D-D proton to D-\( ^3 \)He proton (or \( \alpha \)) yield ratio is

\[
\frac{Y_{DD-P}}{Y_{DHe}} = \frac{1}{2} \frac{\langle \sigma v \rangle_{DD-P}}{\langle \sigma v \rangle_{DHe}}
\]

(3.4)

These ratios are shown in Figure 3.3. The rapidly varying D-D proton to D-\( ^3 \)He proton rate coefficient ratio makes it a sensitive temperature diagnostic. The highly penetrating 14.7 MeV D-\( ^3 \)He proton is crucial for probing high \( \rho R \) targets when other charged particles are ranged out. In these cases, the D-D neutron yield can be used to determine the D-D component of the yield. The ratio of D-T \( \alpha \) and D-D proton yields has been measured before by time of flight spectroscopy using a plastic scintillator [42] and been found to be consistent with other measurements.

Figure 3.3: The ratio of important charged fusion product yields. The rapidly varying D-D proton to D-\( ^3 \)He proton ratio is a particularly sensitive temperature diagnostic. In addition, the highly penetrating 14.7 MeV D-\( ^3 \)He proton enables high \( \rho R \) targets to be probed, particularly if the D-D neutron yield is used instead of the D-D proton yield.
3.2.2 Doppler broadening

The finite temperature of a plasma leaves its imprint on the energy spectrum of products generated in fusion reactions, in particular the primary reactions. This is because the kinetics of the reacting particles determine the output energy of the products. Using the results of calculations by Brysk [43] and Hutchinson [44], for a reaction \( A + B \rightarrow C + D \), the energy spectrum of reaction products, \( C \), from a thermal plasma is given by a Gaussian:

\[
f(E_C) \sim \exp \left[ -\frac{(E_C - \langle E_C \rangle)^2}{2 \Sigma^2} \right] = \exp \left[ -\left( E_C - \langle E_C \rangle \right)^2 \right] \frac{4m_cm_D}{(m_C + m_D)MQT} \]

(3.5)

where \( \Sigma^2 = \frac{2m_cm_D}{(m_C + m_D)MQT} \) \( \Sigma \)

(3.6)

where \( T \) is the temperature of the reactants, \( Q \) is the nuclear energy release of the reaction, \( m_C \) (\( m_D \)) is the mass of particle \( C \) (\( D \)) and \( M = m_A + m_B \). The average energy, \( E_C \), is given by:

\[
\langle E_C \rangle = \frac{m_D}{m_C + m_D} Q
\]

(3.7)

Note that, in these calculations, all terms of order \( T/Q \) have been neglected. For the reactions of interest, \( \Sigma^2(D - T) = 5630T_{keV} \), \( \Sigma^2(D - D, p) = 1510T_{keV} \), \( \Sigma^2(D - ^3He) = 5880T_{keV} \) \( \Sigma \)

(3.8)

where, \( \Sigma \) is in keV. The width of the spectrum, termed the Doppler width, thus provides an ion temperature estimate that is independent of absolute yield measurements or nuclear cross section calculations.

Doppler width measurements of neutron spectra using the time-of-flight technique have been undertaken in both “current-mode” [45, 46] and “single-hit” mode [47, 48, 49]. In general, current mode detectors are useful for yields of \( 10^{10} \) and above while single-hit arrays are useful at lower yields down to \( 10^6 \).

Time-of-flight methods have also been used in determining ion temperatures from the spectrum of D-T alphas [50, 42]. The instrumental response of the scintillator detector systems used in these experiments limited the energy resolution to \( \sim 130 \) keV. Reference [50] reported that the measured temperatures were only an upper limit, claiming that other effects were also responsible for the observed broadening. These effects were energy straggling, non-uniform or time-varying energy losses in the target, and gross hydrodynamic motion of the fuel. Attempts to account for these effects have produced some agreement with ion temperatures predicted from the two-dimensional hydrodynamic code LASNEX. Reference [42] reported that, after accounting for time-varying collisional energy losses in the fuel and ablator using computational modeling of the target implosion, the ion temperatures were in reasonable agreement with
code predictions and temperature measurements using D-T alpha to D-D proton yield ratios. The exception was when the alphas were emitted while laser illumination was still present. In this case, enhanced spectral broadening took place that was attributed to time-varying electric fields in the target corona. In neither of these references was mention made of the D-D proton width.

3.3 Areal density ($\rho R$)

As described in Section 1.2.1, ignition of an ICF target at temperatures of ~10 - 20 keV is achieved if fuel areal densities, or $\rho R$'s, are greater than ~ 0.3 g/cm². The fuel $\rho R$ is an indicator of how much of the alpha energy is given to the fuel and how close the target is to achieving ignition. Areal density is defined as the line-averaged density-radius product $\langle \rho R \rangle$, and is given by

$$\langle \rho R \rangle = \int_0^R \rho dr$$

(3.9)

Unless otherwise stated, $\rho R$ should be assumed to mean $\langle \rho R \rangle$.

For a test particle moving through a target material, this parameter determines the number of target particles that are "in the way" of the test particle per unit area normal to its direction of motion. Thus $\rho R$ is a measure of the likelihood for interactions between the test particle and the material. Because of the vast difference between the Coulomb and nuclear cross sections, the dominant effect of $\rho R$, at least for densities typical for ICF, is different for nuclear interaction processes (neutron scattering and fusion reactions) than it is for Coulomb interaction processes: a higher $\rho R$ tends to increase the yield of neutron-scattered ions and secondary fusion reactions, whereas it tends to downshift the average energy of charged particles traveling out of the plasma. Both of these effects can be used as diagnostics for $\rho R$.

The knock-on, secondary, and tertiary reaction methods are techniques that determine fuel $\rho R$ from the yields of particular side nuclear processes. Each process can be identified by its characteristic spectrum. The advantage of these techniques is that the fuel $\rho R$ can be determined directly, without using complicated models, computational or analytic. The main disadvantage is that they are second-order (or third-order) processes, requiring at least two sequential nuclear reactions, which means that their probabilities are low. Any detector of these particles must contend with the much larger background of primary fusion products. Another measurement approach is to detect the energy downshift of primary charged-fusion products, a technique that avoids the problem of low signal-to-noise. Energy loss is, in general, a result of ranging in both the fuel and the ablator and is affected most strongly by areal density and electron temperature. It is more complicated to unfold the contributions from each of these effects but under many relevant conditions, this can be done.
3.3.1 Fuel areal density from scattered ions

The principle of this method is to measure the number of deuterons and tritons elastically scattered by 14.1-MeV primary fusion reactions. The yield of these ions, termed knock-on particles, will be shown to be proportional to the fuel $\rho R$ and the neutron yield, and, for conditions of interest, is three to four orders of magnitude below the yield of the primary products. The characteristic spectrum of the scattered ions allows their unambiguous identification, provided that the signal-to-noise issues can be understood. This method was first described in reference [25].

Assuming that a yield, $Y_0$, of 14.1-MeV neutrons is produced at the center of a D-T target, the number of elastically scattered particles, $Q^{KO}$, generated by passage of these neutrons through a spherical fuel mass of radius $R$ is given by

$$Q^{KO} = \int_0^R Y_0 (\sigma_d n_d(r) + \sigma_t n_t(r)) dr$$

(3.10)

where $\sigma_d$ and $\sigma_t$ are the elastic scattering cross sections for deuterons and tritons respectively, and $n_d$ and $n_t$ are their respective number densities. This expression assumes that the probability for interaction of the neutrons with the fuel is small enough that depletion of the neutron beam, or multiple scattering, is negligible. Typically, $\sigma < 10^{-24}$ cm$^2$ and $R \sim 10^{-2}$ cm, and, for reasonably high-density conditions, $n \sim 10^{24}$ cm$^{-3}$, which give a suitably small neutron interaction probability of $< 10^{-2}$.

Now, using Eq. (3.9), the $\rho R$ of the fuel is given by

$$\langle \rho R \rangle = \int_0^R (m_d n_d(r) + m_t n_t(r)) dr = (2\gamma + 3)m_p \int_0^R n_t(r) dr$$

(3.11)

where $m_d$, $m_t$, $m_p$ are the deuteron, triton and proton masses respectively and, $\gamma = n_d/n_t$, which is independent of $r$. Replacing $n_d$ with $\gamma$ in Eq. (3.10) and then substituting Eq. (3.11) leaves

$$Q^{KO} = \frac{Y_0 (\gamma \sigma_d + \sigma_t)}{(2\gamma + 3)m_p} \langle \rho R \rangle$$

(3.12)

Thus the yield of knock-on particles is directly proportional to the fuel $\rho R$ and the neutron yield. Measurement of $\rho R$ using this method involves identifying the knock-on particles accepted into a given detector solid angle and simply counting their numbers. An independent measurement of the neutron yield $Y_0$ is required for this calculation. Note that $\langle \rho R \rangle$ as written here is actually temporally integrated over the burn period. This will be true for all measurements of $\langle \rho R \rangle$ discussed, unless stated otherwise. The above calculations implicitly assume that all neutrons are generated at the center of the target; should the burn be uniformly distributed over the fuel, the right-hand side of Eq. (3.12) should be multiplied by 0.75.
It is somewhat misleading, to claim that this method is purely a yield counting technique that does not require a measurement of the energy spectrum of knock-on particles. Any practical detector will have a lower energy cutoff for the spectrum of detectable knock-on energies shown in Figure 3.4. Thus, the cross sections used in Eq. (3.12) must be the effective cross sections, $\sigma^{\text{eff}}$, for generating knock-ons above the threshold or within a pre-determined energy range. For this purpose, the high energy "ramp" feature, present in both the deuteron and triton spectra, is particularly useful. This ramp, between the nascent energies 7.3 – 10.6 MeV for the tritons and 9.6 – 12.5 MeV for the deuterons, should be distinguishable if the spectrum is downshifted. By only counting knock-ons within this feature, $\sigma^{\text{eff}}$ will have been pre-defined. The effective cross-sections and the calculated knock-on yields for tritons and deuterons within these ramp regions are summarized in Table 3.2.

<table>
<thead>
<tr>
<th>Nascent energies for ramp</th>
<th>Effective cross section</th>
<th>KO Yield Expression</th>
<th>KO Yield Convenient form</th>
</tr>
</thead>
<tbody>
<tr>
<td>Deuterons 9.6 – 12.5 MeV</td>
<td>0.104 barns</td>
<td>$\frac{\sigma_d^{\text{eff}}}{5m_p} \langle \rho R \rangle Y_0$</td>
<td>0.0125$&lt;\rho R&gt;Y_0$</td>
</tr>
<tr>
<td>Tritons 7.3 – 10.6 MeV</td>
<td>0.142 barns</td>
<td>$\frac{\sigma_t^{\text{eff}}}{5m_p} \langle \rho R \rangle Y_0$</td>
<td>0.0170$&lt;\rho R&gt;Y_0$</td>
</tr>
</tbody>
</table>

Table 3.2: Knock-on yields calculated for the ramp region, assuming a 1:1 deuteron to triton number ratio. In the expression for KO Yield in convenient form, $<\rho R>$ is in g/cm$^2$. These expressions assume a central point burn model; if the target undergoes uniform burn, the expressions are multiplied by 0.75.

Figure 3.4: Differential cross-sections for elastic scattering of 14.1-MeV neutrons on protons, deuterons and tritons (adapted from [51]), where the energy is the scattered ion energy. The spectra of scattered particles emerging from an ICF target should have similar shapes to these cross-sections.
The downshift of the knock-on spectrum and the yield of knock-ons are actually two independent measurements of $\rho R$. Spectral downshift is a function of both fuel and ablator $\rho R$, while the knock-on yield is purely a function of fuel $\rho R$ (if there are no deuterons or tritons in the ablator). Clearly the total $\rho R$ should be greater than or equal to the fuel $\rho R$. A description of how energy downshifts from multiple species can be used to determined fuel and ablator conditions will be given in Section 3.3.2.

The knock-on method will work only if the scattered deuterons or tritons emerge from the target with enough energy to be registered and identified by the diagnostic. Taking this “minimum” energy to be 5 MeV for the peak knock-on energy, knock-on deuterons are useful if the total $\rho R$ is below $\sim 0.1 \, \text{g/cm}^2$, assuming a 1 keV plasma. A technique for increasing the maximum $\rho R$ accessible by knock-ons is to mix the D-T fuel with a significant fraction of hydrogen and measure elastically scattered protons. The predicted proton knock-on spectrum is shown in Figure 3.4. Generation of protons out to 14.1-MeV potentially enables measurement of fuel $\rho R$'s up to $\sim 0.25 \, \text{g/cm}^2$. Since the knock-on proton spectrum is constant at all energies, the total number of protons can be determined from the magnitude of the spectrum (in yield per unit energy) at any energy (assuming of course minimum distortion to this flat profile), without the need for a ramp feature or effective cross-section values.

Extensive studies of $\rho R$ using knock-on particles have been made on the OMEGA experiment [27, 28]. All of the work has utilized CR-39 plastic track material with different thickness range filters. An elaborate system to discriminate between signal and noise tracks was implemented that caused restrictions in the accessible energy window: Only deuterons from 4.6 – 6.8 MeV and tritons between 5.4 – 10.3 MeV could be extracted unambiguously from the noise background. Typically fewer than 100 particles were counted and the energy resolution was around 1 MeV. The uncertainty in $\rho R$ was about 30%. $\rho R$'s were from 0.1 – 1 mg/cm$^2$ for gas targets and 10 – 40 mg/cm$^2$ for cryogenic targets [52].

Experiments to detect proton knock-ons have been made by utilizing nuclear emulsions [29]. Due to detector limitations, only protons above 9.5 MeV could be distinguished from the deuteron knock-ons. Typically about 100 protons were detected above this energy with an energy resolution of 1 MeV. The uncertainty in $\rho R$ was about 20% for measured $\rho R$'s of 3 mg/cm$^2$.

The great advantage of the knock-on method is that it is a simple and direct technique for determining the fuel areal density. It is also possible to determine ablator $\rho R$ using knock-ons from the ablator. This was recently done using deuteron knock-ons from a deuterated ablator [53]. This provides a good comparison of the energy ranging methods for ablator $\rho R$ measurement to be described in Section 3.3.2.
3.3.2 Charged-particle energy loss as a measure of ablator (and fuel) $\rho R$

While the yield of knock-ons (as well as secondary and tertiary particles) is a direct measure of fuel $\rho R$, the spectral downshift of these, and in fact all charged fusion products, gives an estimate of the total target $\rho R$. In the most general case, the slowing down of charged particles in ICF targets depends on the areal density and temperature of both the fuel and ablator, a total of four unknowns. For many cases of interest though, most of the energy is lost in the ablator – reducing the number of unknowns to two. In the case of high energy protons, with velocities much greater than the plasma electron velocities, the slowing down is independent of temperature, allowing $\rho R$ to be determined from a single energy loss measurement. The advantage of using energy downshifts is that primary fusion products can be used, thus avoiding the problems of low signal-to-noise associated with knock-ons, secondaries and tertaries.

The slowing down of charged particles in a plasma has been studied using a variety of approaches, among them binary collisions [54, 55], plasma dielectric response [56], and Monte Carlo simulations [57]. The method of Li & Petrasso [12], utilizing the binary collision formulation, will be used for all the calculations in this thesis. Using this approach, the stopping powers for a range of fusion-relevant particles are shown in Figure 3.5.

![Figure 3.5](image)

Figure 3.5: Stopping powers, normalized to the square of the charge, for protons, deuterons, tritons and alphas in a 1 keV plasma. The plasma is composed of D-3He at a density of $10^{24}$ cm$^{-3}$. The peak occurs when the particle velocity is near the electron plasma velocity. The small discontinuity in the curves is due to the imposition of collective effects when particle velocities exceed electron plasma velocities.
Before discussing how energy loss measurements can be used to determine $\rho R$, a few qualitative remarks concerning the physics of charged-particle stopping in ICF plasmas will be made. In a plasma where the velocity of the charged fusion product (test particle) is much greater than the average electron velocity, the energy loss, which is due mostly to plasma electrons, becomes virtually independent of temperature. Physically this can be interpreted as the regime where the fusion product velocity is much higher than the average random velocity of electrons and thus, from the perspective of the rapidly moving particle, the electrons are stationary. Any further decrease in temperature has no effect on the energy loss. This is the regime relevant to the stopping of 14.7 MeV $^3$He protons in plasmas of a few keV and is important since now a single energy loss measurement can be used to extract a $\rho R$ measurement. When the test particle velocity is comparable to, or slightly less than, the average electron velocity, the stopping power, which is still dominated by losses to plasma electrons, becomes dependent on the electron temperature. Losses to plasma ions become important when electron velocities are considerably greater than the test particle velocity, as in the case of a 3.5 MeV alpha moving through a 20 keV plasma. For all these different regimes, large-angle scattering has a negligible effect on energy loss [58].

Strictly speaking, the energy loss of charged particles is a function of density as well as temperature and $\rho R$; however, this is only a weak logarithmic dependence. For experimental purposes, the density need only be specified to within an order of magnitude. Density effects become more significant for degenerate plasmas since, under these conditions, average electron velocities are density dependent.

To understand how energy loss techniques can be used to determine $\rho R$, consider the measurement of the energy downshift of two particle lines. Each energy loss measurement defines a particular curve in $\rho R$ vs. $T$ parameter space as calculated by the stopping power predictions for a particular plasma and plasma density. By detecting two independent, yet simultaneous, line shifts, two different curves may be drawn and the point where they intersect uniquely defines the areal density and temperature. An example of this is shown in Figure 3.6. This technique was examined in reference [59]. By using more than two particle types, the problem can be over-determined and the errors in the measurements can be ascertained. Should an unknown accelerating potential be present, a third particle loss measurement must be performed to find a unique solution to the three unknowns. By using a fourth loss measurement, the whole problem can be over-determined.

It is important to recognize that the target parameters inferred from these energy loss measurements are path-averaged. Typically there can be significant radial variation in target $\rho R$ and temperature, and measuring the particle energy loss determines the "stopping-power-weighted" average values of these parameters. This means that regions with high density and low temperature will be weighted the most
strongly. Care must taken when comparing these results with those from other instruments such as x-ray diagnostics which are often sensitive to the maximum ablator temperatures. In comparison, the fuel $\rho R$ techniques described above determine a path-averaged but unweighted $\rho R$ value.

![Figure 3.6](image_url)

Figure 3.6: Predicted relationships of $\rho R$ and $T$ if D-D protons lose 0.5 MeV and D-T alphas lose 1.5 MeV. These curves were calculated for an SiO$_2$ plasma at $10^{24}$ cm$^{-3}$. The intersection point indicates a $\rho R$ of 7.5 mg/cm$^2$ and a temperature of 1.0 keV.

An experiment to measure the downshift of 3-MeV D-D protons is described in Reference [60]. A magnetic spectrometer with CR-39 was used with a quoted energy resolution of 300 keV and typically 100 – 2000 tracks were counted from D-D yields of order $10^4$. Inferred total $\rho R$ was 1 – 10 mg/cm$^2$ with errors of about 30%. Concurrent secondary neutron measurements found the fuel $\rho R$ to be 1.6 – 1.9 times less than this total $\rho R$.

For developing and debugging this energy loss technique, it is useful to study the losses in thin-shelled targets of all the various primary fusion products – D-D protons and tritons, D-$^3$He protons and alphas, and D-T alphas – in order to check for consistency in the slowing down predictions. Once such consistency checks have been validated and confidence in the stopping formula predictions established, thicker-shelled targets can be studied. Experiments with thicker ablators are in fact the focus of experiments relevant to high-gain ICF. Under such conditions, it will be the 14.7 MeV D-$^3$He proton that will be of the greatest utility, since the other particles will be ranged out. As mentioned earlier, a given energy loss measurement of the D-$^3$He protons can be used to predict $\rho R$ directly, without cross-measurements from other particles, since the stopping power of such high velocity particles is quite insensitive to temperature. A theoretical study of the effects of different target conditions on this proton spectrum is provided in reference [61]. Since energy loss measurements can utilize primary fusion products, they have the capability for low-yield $\rho R$ measurements.
4 Laser-plasma instabilities, hot electrons and accelerated ions

The objective of laser heating in direct-drive ICF is to convert laser energy rapidly into plasma thermal energy, thus generating the pressures necessary for high-density target implosions. This conversion is achieved most efficiently via collisional absorption, or inverse bremsstrahlung. However, collective effects involving plasma waves reduce the coupling of laser light to useful thermal energy by instead scattering light and producing suprathermal electrons. Much effort has been devoted to studying, and trying to limit, these various instabilities. Studies have concentrated usually on observing the reflected laser light, x-ray emissions from hot electrons, or the hot electrons themselves. It is well-known that the presence of such electrons can generate strong electrostatic fields which accelerate ions from the ablator surface to suprathermal energies. In addition, these fields accelerate charged fusion products emerging from the core. The spectrum of fast ablator ions and the acceleration of fusion products present alternative methods for probing conditions in the corona.

4.1 Laser-plasma instabilities and production of hot electrons

In order to achieve high gain in laser fusion experiments, it is necessary for a significant fraction of the laser light to be absorbed by the target. However, under certain conditions, a portion of this absorbed energy can go into stimulating electron plasma waves which generate deleterious suprathermal electrons that preheat the target, preventing high compression from being achieved. The generation of hot electrons is reduced to acceptable levels by keeping laser intensities below $\sim 10^{15}$ W/cm² and using sub-micron wavelength light. Producing sub-micron light requires frequency doubling or tripling of the 1.05 µm wavelength light produced by Neodymium-Glass lasers. Under these conditions, the fraction of absorbed laser energy is $> 80\%$ [20], with only about 0.01% – 0.1% given to hot electrons [62].

A schematic of the different laser-plasma interaction processes relevant to ICF are shown in Figure 4.1. Collisional absorption, which occurs throughout the underdense plasma up to the critical density, is the most desirable form of energy deposition since it heats the bulk of the electron distribution function, rather than just the most energetic particles. In fact, slower electrons are preferentially heated since they are more collisional. This form of absorption is most effective at higher densities; thus, shorter wavelength light, which has a higher critical density, deposits more of its energy this way.

On the other hand, if the laser energy is absorbed in the form of electron plasma waves, nonlinear damping, through particle-trapping or wave-breaking, results in energy being given to a small population of energetic electrons. There are three main instabilities that lead to formation of electron plasma waves: resonant absorption, stimulated Raman scattering (SRS), and two-plasmon decay (TPD). Resonant
absorption [63] occurs at the critical surface and results in excitation of a Langmuir wave at the plasma frequency. Stimulated Raman scattering [64] takes place at densities below “quarter-critical” – or a quarter of the critical density – and refers to decay of the incident electromagnetic wave to an electron plasma wave and a scattered electromagnetic wave. Two-plasmon decay [64] describes the decay of an incident electromagnetic wave into two electron plasma waves, and is only possible near the quarter-critical surface. Both SRS and TPD are parametric instabilities and thus have a threshold intensity for excitation below which damping of the daughter waves and plasma inhomogeneities do not allow the instability to grow. Expressions for these thresholds and their expected values for 0.35 µm light are shown in Table 4.1. One important experimental observation [65] concerning SRS and TPD is that, whereas Raman scattering is strongly reduced by application of beam smoothing (SSD in this case), two-plasmon decay is affected only slightly.

Figure 4.1: Schematic electron density profile through the underdense corona region showing the various laser-plasma interactions important in laser fusion experiments and where they occur. Laser energy is deposited up to the critical density and electron conduction transfers energy to higher densities. Resonant absorption, two-plasmon decay and stimulated Raman scattering generate electron plasma waves which are responsible for suprathermal electrons.

There are two other mechanisms important in laser-plasma interactions, though these do not (at least directly) lead to generation of electron plasma waves and suprathermal electrons. Stimulated Brillouin scattering, another resonant instability, involves decay of the incident wave into an ion acoustic wave and a scattered electromagnetic wave. Filamentation is a non-resonant instability, and describes the non-
linear self-focusing of the incident laser by ponderomotive or thermal forces. This can cause break-up of the beam into intense filaments which may exacerbate many of the processes listed above. Both these processes can occur throughout the underdense plasma.

<table>
<thead>
<tr>
<th></th>
<th>TPD</th>
<th>Absolute SRS</th>
<th>Convective SRS</th>
</tr>
</thead>
<tbody>
<tr>
<td>Threshold (W/cm²)</td>
<td>$0.54 \times 10^{14} \left( \frac{T_e}{L_N \lambda} \right)$</td>
<td>$4 \times 10^{14} \frac{L_N^{4/3} \lambda^{2/3}}{L_N \lambda}$</td>
<td>$4 \times 10^{15} \frac{L_N \lambda}{L_N}$</td>
</tr>
</tbody>
</table>

For $\lambda = 0.35 \mu m$, $8 \times 10^{14}$ W/cm², $9 \times 10^{14}$ W/cm², $1 \times 10^{16}$ W/cm²

Table 4.1: Threshold intensities for two-plasmon decay (TPD) [18] and stimulated Raman scattering (SRS) [62]. The thresholds are different for the convective SRS instability (which occurs below quarter-critical) and the absolute SRS instability (which occurs near quarter-critical). $\lambda$, the laser wavelength, is in units of $\mu m$; $L_N$, the density scale length, is in units of 100 $\mu m$; and $T_e$, the electron temperature, is in units of keV. Sample intensity values are calculated for $\lambda = 0.35 \mu m$, $L_N = 100 \mu m$, and $T_e = 0.5$ keV.

Because of the detrimental effects of hot electrons on target compression, many studies have been undertaken to examine the temperature and fraction of laser energy absorbed by such electrons, particularly for long (> 1 $\mu m$) wavelength lasers where the effects are most severe. For $\lambda < 1 \mu m$ at fusion relevant intensities though, only a few experiments designed to directly measure the electron spectrum have been undertaken [66, 67, 68]. These indicated that for intensities of $10^{15}$ W/cm² and $\lambda = 0.35 \mu m$, hot electrons had temperatures of approximately 50 keV and carried a total energy of about $10^4 - 10^5$ of the incident laser energy. Disagreement exists over whether two-plasmon decay [66] or stimulated Raman scattering [68] is the source instability.

### 4.2 Accelerated ions

The presence of hot electrons in laser fusion experiments has been observed to cause acceleration of ions up to suprathermal energies [69, 70, 71]. Theoretical attempts to interpret the resulting spectra have utilized analytical techniques [72, 73, 74, 75, 76] as well as numerical simulations [69, 77]. Studies of the acceleration mechanism have generally focused on the process of plasma expansion into vacuum: space charge is produced by the more mobile electrons which move rapidly into the vacuum region generating electrostatic fields which eventually draw the ions out to follow them. This will be the point of view taken throughout this thesis.

Two simple analytical models of this ion acceleration process have been presented in detail in Appendix A with the purpose of deriving a hot electron temperature from the measured ion spectrum. The first model is a global view of the target as a charged sphere, which acquires a positive potential when a small fraction of suprathermal electrons escape from its surroundings. This approach gives an estimate of the
maximum energy achieved by the accelerated ions (given by the acquired target potential) and how it is related to electron temperatures and total energies. Results from this calculation are shown in Figure 4.2.

![Figure 4.2: Predicted variation of target potential for hot electron temperatures up to 150 keV. This model calculates the potential by determining the number of electrons with enough energy to escape the target potential they generate. The maximum ion energy acquired is assumed to be equal to the target potential multiplied by the charge of the ion.](image)

The second model is a closer examination of the ion flow process by determining a self-similar solution to the one-dimensional fluid equations. This predicts that the slope of the ion velocity spectrum is a function of the electron temperature, as given by

\[
\frac{dN}{dv} = K \exp\left(-\frac{v}{\sqrt{T/m}}\right) \quad (4.1)
\]

where \(dN/dv\) is the number of ions per unit velocity, \(K\) is a normalizing constant, \(v\) is the ion velocity, \(m\) is the ion mass, and \(T\) is the hot electron temperature. The slope of the \(\ln(dN/dv)\) versus \(v\) curve thus provides a measurement of \(T\). With these two models, a single measurement of the accelerated ion spectrum out to its maximum energy provides two different ways to extract the hot electron temperature.

Energetic ion emissions from laser-produced plasmas were first observed in the early 1960's [78] and have been studied over a wide variety of laser wavelengths, intensities, and target conditions (see the extensive review of charge and Faraday cup experiments given by Gitomer et al. [69]). One important discovery was that the bulk of fast ion signals seen on these experiments was due to protons. This was true regardless of the target composition: aluminum, gold, glass, plastic, or any other material. The reasoning was that only tens of angstroms of hydrogen containing contaminants needed to be present on
the surfaces of these targets in order to produce the measured signals. Such contaminants were probably water vapor or oils. Quantitative measurements of ion distribution functions from CH plasmas [79] indicated that the slowest protons reached the detector at an earlier time than the fastest carbon ions.

Compiling data from a broad range of experiments, Gitomer showed a correlation between the mean fast ion energy and the x-ray determined electron temperature, where the mean energy was determined by the peak in the fast ion current measured on a time-of-flight spectrometer. In addition, the ion energy and electron temperature both scaled with $I\lambda^2$ - a quantity that is proportional to the electron quiver velocity and is a measure of the ponderomotive force strength. Much data exists from observations of ion emissions during the 1970's and 1980's when they were observed to carry away a substantial portion of the incident laser energy. However, since the implementation of sub-micron wavelength lasers and the subsequent reduction of hot electron producing plasma instabilities as described above, the literature appears to contain much less on measurements of accelerated ions from current laser fusion experiments. Therefore, in order to obtain an estimate for OMEGA conditions, Gitomer's scaling between ion energy and $I\lambda^2$ will be used. At $I\lambda^2 = 10^{14}$ W/cm$^2$.\mu m$², corresponding to a laser intensity of $10^{15}$ W/cm$^2$ and wavelength of 0.35 \mu m, this scaling predicts a mean ion energy between 1 - 20 keV.

It is difficult, however, to compare this mean ion energy obtained from time-of-flight measurements, with results from a magnetic spectrometer, such as the one described in this thesis, since the magnet has a low energy cutoff, which might obscure the bulk of the lower energy ions that contribute to the mean. One characteristic of such a spectrometer though, is its ability to finely resolve the spectrum of fast ions out to and including the maximum energy. However, in order to compare results from this diagnostic with previous results in the literature, it is necessary to find studies that have located the endpoint energy of ions and determined electron temperatures from the slope of the ion profile.

Pearlman and Morse [80] investigated the maximum energy of accelerated protons and found that for $I = 10^{14}$ W/cm$^2$, $\lambda = 1 \mu m$ - an equivalent $I\lambda^2$ to the experiments described herein - using a 50 ps pulse delivering 1 J, protons were detected out to 38 keV, which is not too much higher than the mean energies reported by Gitomer. These spectra closely followed the exponential velocity profile given by Eq.(4.1), the slope indicating an electron temperature of 2.5 keV. This temperature is consistent with other measurements at the same $I\lambda^2$ as indicated by Gitomer's correlation between $T_e$ and $I\lambda^2$ which show a temperature range of 0.8 – 4 keV. To compare the electron temperature deduced from the maximum proton energy, Eq. (A.4) (from Appendix A) is used with a target potential of 38 kV, a potential well “radius” of 1 mm, and a suprathermal electron energy of 0.01 J. This gives a temperature of $\sim$ 6 keV, slightly above that derived from the slope of the ion spectrum, but of the same order.
However, if electron temperatures are determined primarily by the value of $I\lambda^2$, these measured temperatures are at odds with the direct electron spectral measurements of Rousseaux [68] reported above for $I = 10^{15}$ W/cm$^2$ and $\lambda = 0.26$ μm which produced electron temperatures of 50 keV. Note that the experiments reported by Rousseaux used 25 J in 450 ps pulses. Temperatures inferred from the ion spectra described in this thesis, at $I = 10^{15}$ W/cm$^2$, $\lambda = 0.35$ μm, and 1 ns pulse length should be able to distinguish between hot electron temperatures of 2 keV and 50 keV for an $I\lambda^2$ of $10^{14}$ W/cm$^2$.μm$^2$.

Apart from the studies by Pearlman, there have been other experiments which have verified that ion spectra obey the self-similar velocity profile [81, 82, 83] given by Eq. (4.1); however some showed differences. In particular, Decoste and Ripin [70] report the appearance of oscillations in their energy spectra that they speculated were due to temporal variations of the pressure gradient giving bursts of ions at decreasing energies. Also, Wagli and Donaldson [71] show proton spectra with a distinctive peak, and Fews et al. [84] describe the unpredictable appearance of both exponentially decaying spectra and flat spectra, though these later experiments were performed at much higher intensities of $10^{18}$ W/cm$^2$.

Closely related to the production of fast ions is the observation of accelerated charged fusion products. Fusion products pass through the corona region and can be expected to receive similar accelerations to those acquired by the fast ablator ions. Measurements of D-T alphas and D-D protons [42, 85] show energy upshifts of 310 keV and 260 keV respectively, though these were probably lower limits since the particles appeared to undergo some ranging. In these experiments, laser intensities were $5 \times 10^{15} - 5 \times 10^{16}$ W/cm$^2$ at 1.05 μm, with laser energies of 70 – 155 J at a pulse width of 40 – 75 ps. At these values of $I\lambda^2$, Gitomer’s correlation reports mean fast ion energies of up to 300 keV which is consistent with the observed upshifts of these fusion products. The decay of the accelerating fields has been probed using fusion products produced at different times during the implosion [85]. The strength of the accelerating fields decays over the duration of the laser pulse because of plasma expansion, since the field strength is approximately proportional to the ratio of electron temperature to density scale length. At the end of the laser pulse, this decay is even more rapid because hot electron production ceases, the remaining hot electrons cool during expansion, and currents from the stalk neutralize the target. Non-thermal broadening of the D-T alpha line width was attributed to these field variations with time.

In summary, both accelerated fusion products and fast ions (which will from here on be termed ablator ions to distinguish them from the fusion products) are thus indicators of the magnitude of the target potential and the hot electron temperature. An interesting question has been posed of what the hot electron temperature at an $I\lambda^2$ of $10^{14}$ W/cm$^2$.μm$^2$ really is – 2 keV or 50 keV. The maximum ion accelerations at these temperatures should be either tens of keV or hundreds of keV respectively – a difference that should be immediately apparent.
5 Spectrometer design

The essential challenge of time-integrated charged-particle measurements on ICF experiments is to determine the energy and identity of hundreds, perhaps thousands, of particles arriving within several nanoseconds. Two basic approaches have been used in the past. The first is to use a detector with a high degree of spatial resolution, such as plastic track material, which allows individual particles to be spatially discriminated when the detector is examined after the shot. The second is to use time-of-flight techniques which discriminate groups of particles by their arrival time. In general, spatially resolving the particles using track detectors allows individual particles to be uniquely identified, though extracting the data can be laborious and slow. TOF methods are much less labor-intensive but their inability to resolve individual particles makes them more prone to ambiguous signals. Given the wide variety of particles and energies, it is usually important to be able to discriminate individual particles. The spectrometer to be described here utilizes a dipole magnet for spectral dispersion, and CR-39 track detectors for particle identification. In future, charge-coupled devices, which have demonstrated the capability of resolving thousands of individual particles, will replace some of the track detectors, enabling efficient electronic readout. Ultimately, pin-diodes placed behind the magnet will provide the additional advantage of time discrimination of signal and noise. The magnet plus CR-39 design allows the most reliable and flexible setup for examining the large range of particles, energies and yields that comprise charged-particle spectra. Once these spectra have been well-characterized, electronic methods may be used to expedite the measurements.

Track detectors have a long history of use in ICF experiments. Two experimental setups close to the one described here used a dipole magnet and CR-39 for measuring protons from D-D fusion [60] and from an intense proton beam [86]. The latter experiment also utilized pin-diodes in conjunction with the CR-39. In another experiment performed before the discovery of CR-39, LR115 cellulose nitrate track material was used behind a magnet [81] to detect D-T alpha particles. LR115 is insensitive to protons above a few hundred keV so D-D protons could not be detected. The same paper reports the use of a different magnet, again with LR115, that was used to detect ablator ions, although individual particle species could not be identified. Numerous experiments have utilized CR-39 to measure charged fusion products without the aid of a magnet. Many of these were reported earlier (Section 3.3) in the context of \( pR \) measurements using knock-ons (and also secondary products). It is possible to achieve energy resolution from the track detectors alone since track diameters are a measure of the incident particle stopping power. However, resolution is limited and many ambiguities can result when multiple species are present, although the simplicity and convenience of this experimental setup is difficult to match. One experiment reported the use of nuclear emulsions to measure proton knock-ons [29]. The advantages of emulsions over simpler track detectors seem to be outweighed by the extra effort required to extract data.
There have been a few reported TOF measurements of charged fusion products: D-T alphas were used for ion temperature measurements [50] and for probing electric fields [42, 85]. Experiments to determine the ablator $\rho R$ have been performed by comparing the downshift of D-T alphas and D-D protons [87]. Typically these experiments use a quadrupole magnet to focus particles onto plastic scintillators which are coupled to photomultiplier tubes. Most ablator ion spectra have been measured using TOF techniques (see [69]).

5.1 Instrument design

5.1.1 Design concept

A successful charged-particle spectrometer must not only determine the energy of the ions and discriminate them from the background, it must also distinguish between the different types of light ions. Fusion products of interest generated in ICF experiments are protons, deuterons, tritons, $^3$He, and alphas. In addition to these particles, ablator ions may also consist of the carbon, oxygen and silicon nuclei that make up the plastic and glass ablator materials. For the purposes of this discussion, a particle is uniquely identified if its mass number, $A$, and its atomic number, $Z$, are determined. Along with the energy, $E$, of the particle, this makes 3 independent unknowns. Given the restricted types of particles present though, it is usually sufficient to determine just $E$ and $A$. For instance, protons, deuterons and alphas are uniquely identified by their mass since no other relevant particles have mass numbers of 1, 2 or 4 respectively. In addition, the energy of the tritons (from D-D or knock-on reactions) will distinguish them from $^3$He ions for most energy regimes of interest. The requirements for an ICF charged-particle spectrometer thus reduce to determining $E$ and $A$.

In order to determine these two unknowns, at least two measurements must be made on the same particle. This spectrometer uses a dipole magnet and CR-39 track material: the dispersion of particles through the magnet provides one measurement, the track response provides the other. A schematic of the instrument is shown in Figure 5.1. Ion trajectories through a magnetic field are governed by the particle gyroradius, $r_L = p/(ZeB)$, where $p$ is the particle momentum, $e$ is the electron charge, and $B$ is the magnetic field intensity. In terms of $A$, $Z$ and $E$ then, $r_L \sim AE/Z^2$. For a well-collimated source of particles (as is the case for ICF targets which are effectively point sources) all particles with the same $AE/Z^2$ will follow the same trajectories through the magnet and arrive at the same detector positions. The energy of an ion with a given $A/Z^2$ can then be determined by its arrival position behind the magnet. Degeneracies exist though. For instance 3-MeV protons, 1-MeV tritons, and 3-MeV alphas will all arrive

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5 The problem becomes more involved when considering ions with multiple charge states, such as $C^{6+}$, $C^{5+}$, which may exist in the ablator ion spectrum. These situations need to be resolved case by case.
at the same positions. They must then be discriminated by examining the different track diameters created by each particle type.

7.6 kG MAGNET

50 keV

200 keV

600 keV

1.0 MeV

3.0 MeV

10 MeV

Figure 5.1: Schematic of the spectrometer concept. A pentagonal dipole magnet disperses protons in the range 0.1 – 30 MeV. CR-39 nuclear track detectors are placed in the dispersion arc.

To understand how track detectors can break the degeneracy that arises for different ions with the same \( AE/Z^2 \), the variation of track diameters with particle \( E, A \) and \( Z \) must be understood. To first approximation, the track diameter of a particle incident normally on the detector surface is a function of the particle stopping power, \( dE/dx \), at the detector surface. The approximation is reasonable for many of the MeV charged fusion products of interest. The stopping power of MeV ions in solid materials is described well by the Bethe-Bloch formula:

\[
\frac{dE}{dx} = \frac{Z^2}{v^2} \left( \frac{e^4}{4\pi \varepsilon_0^2} \right) \left( \frac{nz}{m_e} \right) \ln \left( \frac{2m_e v^2}{I} \right)
\]

where \( v \) and \( Z \) are the particle velocity and atomic number, \( n \) and \( z \) are the number density and atomic number of the target material, \( m_e \) is the electron mass, and \( I \) is the mean excitation energy of the target atomic electrons. In terms of the incident particle properties, neglecting the slowly varying logarithm term means that \( dE/dx \sim Z^2/v^2 \). Therefore,
\[ \frac{dE}{dx} \sim \frac{AZ^2}{E} = \left( \frac{A^2}{AE / Z^2} \right) \] (5.2)

At a given detector position behind the magnet, \( AE/Z^2 \) is constant so the stopping power of different particles at any given position varies with mass alone, and not charge. Relating the track diameter, \( D \), to the particle stopping power by a power law, \( D \sim (dE/dx)^{1/\eta} \), where \( \eta \sim 1 \) (the conditions when this applies are discussed in Appendix B) then at any detector position the diameter is related to the particle mass by:

\[ D \sim \left( \frac{dE}{dx} \right)^{1/\eta} \sim A^{2/\eta} \] (5.3)

Although this description is an oversimplification, the qualitative result of Eq. (5.3), namely that, at any given detector position, particles with larger mass have larger track diameters, generally holds. A model that describes the general process of track growth that allows track diameters to be predicted is described in Appendix B.

In summary, the magnet dispersion determines the particle energy, while the track response determines the particle mass, and hence identity. This general concept is used in some detector systems used in particle physics experiments [88]. A similar detector system using CR-39 and PIN-diodes has been used to characterize intense proton beams [86]. Application to ICF experiments is outlined in references [89, 90] where an electronic version of this spectrometer, using charge-coupled devices (CCD) instead of CR-39, is described. The CCD response to individual charged particles is very similar to that of CR-39: the intensity of the particle signal is directly related to the particle stopping power and all the arguments above apply equally well.

### 5.1.2 Dipole magnet

The design of the magnet is driven by two competing needs. First, the field must be strong enough to disperse particles from up to 10.6-MeV tritons (particles with the largest gyroradius) as set by the knock-on triton endpoint. Second, the magnet must be compact and lightweight enough that it can be placed into a re-entrant module inside the target chamber to maximize signal (although the first spectrometer built for this experiment is mounted outside the target chamber). A good compromise is found in the 7.6 kilogauss Neodymium-Iron-Boron pentagonal permanent magnet shown in Figure 5.2. The pentagonal shape ensures that there is minimal "wasted" field, meaning that there is little field area through which particles of some energy do not pass. This fact is important to keeping the size and weight of the magnet low. The magnet and yoke structure is 28 cm long, 17 cm wide (to the apex) and 20 cm high, while the gap height is 2 cm. The entire construction weighs about 160 pounds and two virtually identical magnets were constructed by Dexter Corporation [91].

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The magnetic field between the pole faces is 7.6 kilogauss and has 2% uniformity in its central region. The magnetic material is constructed of multiple pieces of Nd-Fe-B which are epoxied together. Steel pole faces smooth out inhomogeneities in the fields that fill the gap region. Fringing fields are minimal owing to the thick steel yoke construction which provides an effective return path: field intensities are 1000 Gauss and 250 Gauss at 2 cm and 4 cm from the magnet edge respectively.

Figure 5.2: The 7.6-kG magnet constructed of Nd-Fe-B and steel. The longest dimension of the magnet is 28 cm, and the gap width is 2 cm. This magnet weighs 160 pounds, and the force between the poles is 6400 pounds. *(Photo courtesy of Frederick Séguin.)*

Particle trajectories through the magnet are shown in Figure 5.3. The trajectories are calculated using a Runge-Kutta fourth order integration of the Lorentz equation[^6]. The predicted fields, as provided by Dexter Corporation, were found to be very close to those measured (also by Dexter).

Apart from being a compact and field efficient design, the pentagonal shape has the convenient property that the particles appear to emanate, to good approximation, from two virtual source points (see Figure 5.4) – one for particles above 0.6 MeV, the other for particles below 0.6 MeV. This is convenient for designing detector positions since it is important that detectors are placed normal to the incident particles. The tracks formed in materials such as CR-39 can vary in shape and size depending upon the incident angle and, in cases of extreme angles, can actually disappear altogether.

[^6]: Initial trajectory calculations were performed by Dexter Corporation.
Figure 5.3: Predicted particle trajectories through the magnet. The front of the magnet is positioned at 100 cm from the target.

It should be noted that often magnetic spectrometers are designed such that detectors lie along the focus envelope of the particle trajectories. This enables good resolution to be obtained even when the collimator slit is opened wide. The reason that this design was not incorporated here is that such an approach makes for a much less compact instrument: Getting particles at all energies to focus, particularly at the higher energies probed by this experiment, requires a considerably larger, stronger, and more complex-shaped magnet, pushing the limits of available permanent magnet technology. In addition, placing detectors along the focal plane means that the particles are no longer normally incident on all the detectors since the envelope of focal points is not generally perpendicular to the particle directions [81]. The price paid by this design compromise is that the resolution of the instrument is limited by the collimator width, typically between 0.1 – 1.0 cm, which must be adjusted depending on whether particle flux or energy resolution is more important.
Figure 5.4: Particle trajectories appear to emanate from two virtual source points. This property facilitates positioning of the detectors so that all detectors are as normal as possible to particle incidence.

5.1.3 Integrated detector system

Implementation of the magnet and CR-39 spectrometer concept involves four essential tasks: Positioning the track detectors correctly with respect to the magnet, aligning the magnet, collimating the incident particles, and placing the whole system in vacuum. Two spectrometers have been built, the first is situated outside the target chamber (CPS-1), the second one is inside (CPS-2), where CPS stands for "charged-particle spectrometer". The port positions of these instruments allow orthogonal views of the target. Although placing CPS-2 inside the target chamber required construction of a sophisticated re-entrant structure, the design of both spectrometers is otherwise very similar.

Positioning of the CR-39 around the magnet is accomplished by the mounting plate and finger system shown in Figure 5.5. The possible finger attachment positions are illustrated in Figure 5.6. A set of figures showing the energy at each of the finger positions is shown in Appendix C. A cover with 1 – 3 different thickness filters is placed over each piece of mounted CR-39. This is required for three reasons: to shield the CR-39 from thermal plasma blow-off, to range out certain particle species, and, in the case of high energy protons, to range down particles to energies that are detectable after reasonable etch times.
Figure 5.5: Photographs of the mounting plate assembly that is used to accurately position pieces of CR-39 in the dispersion arc of the magnet. Figure (a) shows the mounting plate assembly viewed from the perspective of the magnet. Pieces of CR-39 are positioned in each of the finger structures; these fingers are arranged in a few arcs which cover the dispersion region of the magnet. The finger at the bottom of the photo is positioned to view the target directly. X-ray film is placed at this position in order to ascertain the alignment of the spectrometer. The collimator slit is shown at the top of the photo. Figure (b) shows the loaded mounting plate assembly being lowered onto the magnet (which is obscured) inside the vacuum chamber of CPS-1. After every shot, the mounting plate must be removed, and the CR-39 unloaded. A new, freshly loaded plate must then replace it in preparation for the next shot. *(This design was done by the team. Photo courtesy of Eugene Kowaluk, LLE.*)
Figure 5.6: Plan view of the magnet looking through the mounting plate so that each of the finger positions can be seen. The exit trajectories of various particles and energies are shown. The various finger positions form overlapping arcs to provide flexibility and redundancy in selecting the detector positions. The width of a finger actually spans the distance between two positions so that, should an important particle energy fall in a gap between two pieces of CR-39, the corresponding fingers can be shifted up by one step so that these particles now intersect the center of a piece of CR-39. (Design by M. Romanofsky & R. Petrasso, diagram courtesy of M. Romanofsky.)

Alignment of the magnet and detector system for CPS-1 was achieved by pointing a laser beam down co-aligned pinholes on the magnet axis, and reflecting it off a 4 mm reflective ball placed at the target position. Special mounting screws allowed multiple axis movement of the magnet. CPS-2 was not built with such flexibility. A final check of the system alignment was accomplished by imaging the beam collimator using x-rays from the target. Perfect alignment is not crucial (though convenient) since small deviations can be detected by the film and input into the magnet trajectory codes to predict the resulting particle paths.
The beam collimator is a slit placed in front of the magnet and is typically 1.5 cm high and 0.1 cm wide. Only a single slit is required for collimation because the target is essentially a point source. On low yield shots, a 1.0 cm wide slit can be used to increase the solid angle at the expense of energy resolution. The slit size and distance from the target sets the solid angle subtended by each spectrometer. At 100 cm, the solid angle acceptance of CPS-2 is five times that of CPS-1 which lies at 235 cm.

The vacuum system for CPS-1 consists of a cylindrical chamber with a gate valve leading to a port on the OMEGA target chamber. The gate valve is closed between shots to vent the magnet chamber so that the CR-39 can be replaced. A photograph of the chamber is shown in Figure 5.7. The CPS-2 vacuum system involves a substantially more complicated re-entrant structure illustrated in Figure 5.8 (shown during installation), designed and constructed by the Laboratory for Laser Energetics. The magnet and detectors are housed in a retractable housing that can be pulled back when the detectors need to be accessed. The front of the re-entrant module is filled with the shielding materials consisting of polyethylene (closest to the target) and lead (closest to the detectors).

When electronic detectors are finally used, they will be installed in CPS-2 and the detector housing for this spectrometer has been designed with enough room to accommodate up to six 512×512 CCD’s [92] with their accompanying electronics and cooling systems. Although CPS-2 will eventually run with CCD’s and/or PIN diodes, there will likely be an extensive period where only CR-39 is used. The instrument performance characteristics described below incorporate mostly the issues surrounding the track detector version of this diagnostic. The positions of CPS-1 and CPS-2 on an OMEGA port diagram are shown in Figure 5.9.

Figure 5.7: The CPS-1 vacuum chamber mounted on the outside of the OMEGA target chamber. The magnet entry slit is 235 cm from the target. (Photo courtesy of Eugene Kowaluk, LLE.)
Figure 5.8: Installation of CPS-2 on OMEGA. The re-entrant structure contains 4000 lbs of lead and polyethylene shielding. The magnet slit is positioned at 100 cm from the target. (Photo courtesy of Eugene Kowaluk, LLE.)
5.2 Instrument performance characteristics

5.2.1 Dynamic range

Accessible yields are dictated, to a large extent, by restrictions on countable track densities in the CR-39. For this discussion, a minimum countable density of $\sim 500$ tracks/cm$^2$ (determined by the density of noise tracks in CR-39) and a maximum density of $\sim 1 \times 10^7$ tracks/cm$^2$ (determined by the density above which unacceptable track overlap occurs) will be assumed. The precise value of these numbers is dependent upon the size of tracks being measured so the following description should be considered only as rough guides. To quantify the dynamic range of particle yields, a few specific fusion products will be considered.

The relationship between the track density, $\kappa$, found in a particular detector region (such as a spectral line), and the total yield of particles, $Y$, is given by:

$$Y = \frac{4\pi D^2}{f} \frac{\Delta}{W} \left(1 + \frac{d}{D}\right) \kappa$$

(5.4)

where $D$ is the distance from the collimator to the target (235 cm for CPS-1 or 100 cm for CPS-2), $d$ is the beam path distance from the collimator to the detector (dependent upon energy but typically close to 30 cm), $\Delta$ is the spatial width at the detectors of the feature being measured, $W$ is the width of the

Figure 5.9: OMEGA port diagram, showing the locations of CPS-1 and CPS-2. The angle between the two spectrometers is 101°. Also shown are the positions of Range Filter Packs (RFP) which are simply pieces of CR-39 exposed directly to the target. These allow further symmetry studies by investigating the isotropy of particle yields. (Diagram provided by LLE.)
collimator slit (typically between 0.1 mm and 1.0 mm), and $f$ is the fraction of the total particle yield that contributes to the particles contained within $\Delta$. For all calculations below, $D = 100$ cm, $d = 30$ cm and $W = 0.1$ cm. Although $\Delta$ is actually a convolution of the width of the spectral feature and the width of the collimator slit imaged onto the detector plane, using $W = 0.1$ cm the contributions of "slit broadening" to the spectral features discussed here are small (< 20%); for simplicity they will be neglected.

For $D^3$He protons from a 10 keV plasma, the standard deviation of the spectral line is 242 keV, which corresponds to a spatial width of 0.44 mm for a dispersion of 1/550 mm/keV at the 14.7 MeV detector position. This means that the width, $\Delta$, that contains particles within $2\sigma$ of the mean, is 0.9 mm. The fraction of particles, $f$, within $2\sigma$ is, by definition, 0.67, and thus, using Eq. (5.4) with the minimum track density of 500 tracks/cm$^2$, the minimum measurable yield is $1.1 \times 10^8$.

For fusion products, maximum measurable yields are most likely to be encountered with D-T alpha measurements. For these alphas, born in a 10 keV plasma, the standard deviation of the spectral line is 237 keV, which corresponds to a spatial width of 2.79 mm for a dispersion of 1/85 mm/keV at the 3.5 MeV detector position. Particles within $2\sigma$ are thus spread over 5.6 mm. Using $f = 0.67$, and a maximum track density of $10^7$ tracks/cm$^2$, the maximum measurable yield is $1.4 \times 10^9$.

For D-T knock-on tritons, the high-energy ramp feature (see Figure 3.4) from 7.3 – 10.6 MeV extends over a detector distance, $\Delta$, of 7.5 mm. Using Eq. (5.4) with $f = 1$, and the minimum track density of 500 tracks/cm$^2$, gives a yield of knock-ons within this feature of $6.1 \times 10^8$. The total D-T neutron yield required to produce this number of knock-ons is found using Table 3.2 by assuming a value of the fuel $\rho R$. Taking $\rho R = 5$ mg/cm$^2$ this gives a minimum required D-T yield of $7.2 \times 10^{12}$.

For measurements using CPS-1 instead of CPS-2, where $D = 235$ cm instead of 100 cm, all these yield values are increased by a factor of 4.79. Thus, for example, the maximum measurable D-T alpha yield on CPS-1 is $6.5 \times 10^{13}$.

These results are summarized in Table 5.1. For reference, an estimate of the maximum ablator ion yield above 0.2 MeV is given, although it is difficult to treat the ablator ion spectra in the same qualitative manner. The measurable yield of ablator ions can exceed that of the D-T alphas owing to the considerably larger magnet dispersion at lower energies.
Yield limit

<table>
<thead>
<tr>
<th>Minimum D-3He proton yield</th>
<th>$1 \times 10^8$ ($5 \times 10^6$)</th>
</tr>
</thead>
<tbody>
<tr>
<td>Minimum D-T neutron yield to produce sufficient triton knock-ons</td>
<td>$7 \times 10^{12}$ ($3 \times 10^{13}$)</td>
</tr>
<tr>
<td>Maximum D-T alpha yield</td>
<td>$1.4 \times 10^{13}$ ($7 \times 10^{13}$)</td>
</tr>
<tr>
<td>Maximum ablator ion yield above $\sim 0.2$ MeV</td>
<td>$\sim 10^{15}$</td>
</tr>
</tbody>
</table>

Table 5.1: Summary of various yield limits for CPS-2 using a collimator slit width of 0.1 cm. The minimum yield is fixed by a minimum track count of 500 cm$^2$. Yield values in parenthesis are those calculated for CPS-1.

### 5.2.2 Energy resolution

The detector positions shown in Figure 5.6 cover proton energies from $0.1 - 60$ MeV, though particles of interest are restricted to the region below $\sim 32$ MeV. The close spacing of the fingers allows almost 85% coverage of the spectrum in this range. Two factors affect the energy resolution: the dispersion and the collimator width. A finite slit size ensures that any given point in the detector plane sees a bandwidth of energies (except at perfect focus). This bandwidth, which can be taken as the energy uncertainty (or, strictly speaking, twice that), is larger at higher energies, where the dispersion is less. For instance, using a 1 mm wide slit, the beam width at the 3 MeV proton position is 0.74 mm. At this position, the dispersion is about $1/65$ mm/keV so that the range of energies landing at the nominally 3 MeV position is 48 keV. This gives a resolution of 1.6%. On the other hand, the width of the beam up at 15 MeV is 1 mm, and, with a dispersion of $\sim 1/550$ mm/keV, the energy uncertainty is 575 keV. This gives a resolution of 3.8%. Down at 400 keV, where the beam focus happens to occur at the detector plane, the energy resolution is essentially “perfect”. These cases and some others are summarized in Table 5.2.

<table>
<thead>
<tr>
<th>Particle</th>
<th>Energy (MeV)</th>
<th>Beam Width (mm)</th>
<th>Dispersion (mm/keV)</th>
<th>$\Delta E$ (keV)</th>
<th>Resolution ($\Delta E/E$)</th>
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</thead>
<tbody>
<tr>
<td>Proton</td>
<td>3.0</td>
<td>0.74</td>
<td>0.015</td>
<td>48</td>
<td>1.6%</td>
</tr>
<tr>
<td>Triton</td>
<td>1.0</td>
<td>0.74</td>
<td>0.046</td>
<td>16</td>
<td>1.6%</td>
</tr>
<tr>
<td>Alpha</td>
<td>3.5</td>
<td>0.86</td>
<td>0.012</td>
<td>72</td>
<td>2.0%</td>
</tr>
<tr>
<td>Proton</td>
<td>15</td>
<td>1.04</td>
<td>0.0018</td>
<td>575</td>
<td>3.8%</td>
</tr>
<tr>
<td>Proton</td>
<td>0.4</td>
<td>0.00</td>
<td>0.125</td>
<td>0</td>
<td>0.0%</td>
</tr>
</tbody>
</table>

Table 5.2: Energy uncertainties for measurement of a single particle over some salient energies, using a slit width of 0.1 cm. The values have been calculated for the position of CPS-1 but will not be substantially different for CPS-2. Note how the resolution improves at lower energies owing to the focusing of the beam and greater dispersion.
It must be emphasized though that these numbers represent the energy uncertainty of a *single particle measurement*. All measurements of any significance utilize many particles in a spectral line. The uncertainty in the mean of a spectral line is far smaller than the numbers given above but will of course be dependent upon the number of particles counted. A good first approximation is that the error in measurement of the mean energy of particles in a spectral line is \( \Delta E / \sqrt{N} \), where \( N \) is the number of particles counted. Typically, \( N \) is large enough that systematic errors in the energy calibration will dominate.

### 5.2.3 Signal-to-noise

Neutron-induced noise is dependent upon the sensitivity of CR-39 to neutron interactions and, for applications such as knock-on detection, will usually dominate over inherent noise (which is due to defects in the CR-39). Neutrons can interact either directly with the CR-39, or indirectly by producing charged particles through interactions with surrounding material. To minimize neutron noise in CPS-2, polyethylene and lead shielding material has been used to block the direct line-of-sight between the target and the detectors. It appears that this shielding reduces the neutron noise by a factor of ~5 over the un shielded case.

The relative contributions of signal and noise are best observed by studying a spectrum of signal + noise, along with a spectrum of pure noise (a background spectrum). The particular case of a triton knock-on spectral measurement, as detected by CPS-2, will be examined here. Raw data of knock-on spectra from shot 13817 are shown in Figure 5.10, alongside background spectra. The background spectrum is generated by scanning a region of CR-39 devoid of signal events, such as that above the triton endpoint energy. Analysis of this region gives the density of noise events at various diameters which then allows the background spectrum to be estimated. Based on the assumption that the noise spectrum in a detector region *without* signal is the same as the noise spectrum in a detector region *with* signal (a reasonable assumption if the two detector regions are nearby), this background can be subtracted from the initial signal + noise spectrum.

From these plots, the triton + noise yield in the ramp feature (defined as 6.5 – 11.5 MeV for this spectrum) is \( 5.66 \times 10^9 \), while the yield of pure noise events in this energy range is \( 1.98 \times 10^9 \). This gives a signal-to-noise ratio of 1.9. The fuel \( pR \) determined for this shot (see Section 6.3), using both triton and deuteron knock-ons, is 7.0 mg/cm\(^2\). The signal-to-noise ratio can be increased by widening the collimator slit, which was at 1 mm for this shot. Note that, with successful background subtraction, signal-to-noise levels below unity are still useful, depending of course on the level of statistical uncertainty.
Figure 5.10: Comparison of signal and background for the triton knock-on spectrum of Shot 13817 measured by CPS-2. The background spectrum has been generated using noise events found in the region of CR-39 above the triton endpoint. For this reason, above ~ 11.5 MeV, both spectra are equivalent. In the range 6.5 – 11.5 MeV, the signal-to-noise ratio is 1.9. The fuel $\rho R$ inferred for this shot is 7.0 mg/cm$^2$. 

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6 Results & discussion

Three different classes of particles have so far been detected and quantified: a variety of primary fusion products from D-D, D-\(^3\)He and D-T reactions, fuel knock-ons from D-T reactions, as well as fast protons accelerated from the ablator surface. Table 6.1 summarizes the various particles that have been quantified. Unless stated otherwise, results have been taken on or before September 1998.

<table>
<thead>
<tr>
<th>Class of particle</th>
<th>Reaction</th>
<th>Quantified spectra</th>
</tr>
</thead>
<tbody>
<tr>
<td>Primary fusion product</td>
<td>(D + D \rightarrow T + p)</td>
<td>D-D tritons and protons</td>
</tr>
<tr>
<td></td>
<td>(D + ^3)He (\rightarrow \alpha + p)</td>
<td>D-(^3)He alphas and protons</td>
</tr>
<tr>
<td></td>
<td>(D + T \rightarrow \alpha + n)</td>
<td>D-T alphas</td>
</tr>
<tr>
<td>Knock-on fuel ion</td>
<td>(n (14.1 \text{ MeV}) + D \rightarrow n’ + D)</td>
<td>Knock-on D</td>
</tr>
<tr>
<td></td>
<td>(n (14.1 \text{ MeV}) + T \rightarrow n’ + T)</td>
<td>Knock-on T</td>
</tr>
<tr>
<td>Ablator ion</td>
<td>Protons</td>
<td></td>
</tr>
</tbody>
</table>

Table 6.1: List of particles that have so far been detected and their spectra quantified.

The primary fusion products that have been quantified are D-D protons and tritons, D-\(^3\)He protons and alphas, as well as D-T alphas. Results from a single D-\(^3\)He shot are used to illustrate D-D and D-\(^3\)He reaction product spectra and are shown in Figure 6.1. In addition, hints of the \(^3\)He product of D-D fusion have been seen on some shots, though its low energy means it is usually overwhelmed by ablator ions.

Elastically scattered deuterons and tritons have been detected on a few shots and a spectrum is shown in Figure 6.2. The shape and energies of these spectra are very similar to those predicted from differential cross-sections. On some shots filled with D, T and \(^3\)He fuel, some scattered \(^3\)He ions have been detected, though their numbers have so far been too small to allow spectral measurements.
Finally, fast ablator protons have been seen on virtually all shots. Their spectrum is immensely varied, though there is always a sharp cutoff at the maximum energy. Sometimes the distribution is a reasonably smooth continuum, whereas at other times there are sharply defined and equally-spaced lines. Usually there is a combination of the two. A couple of different spectra are shown in Figure 6.3. In addition, copious numbers of heavier ions, most probably including deuterons, $^3$He and carbon have been detected, though they have yet to be quantified.

Figure 6.1: Spectra from D-$^3$He shot # 13804 of four simultaneously measured spectral lines: D-D protons and tritons with D-$^3$He protons and alphas [93]. Upshifted mean energies are due to accelerating fields in the target corona.
Figure 6.2: D-T alpha spectrum from shot 11739. Deuteron and triton knock-on spectra from shot 13822 [94]. The events out beyond the endpoints of the knock-ons are neutron background. Gaps in the spectra filled in by dotted lines represent the space between contiguous pieces of CR-39 behind the magnet.

Ablator Protons

Figure 6.3: Spectrum of accelerated protons from shot 11739 (on the left) and shot 11771 (on the right) [93]. It is still unclear what precipitates the change from reasonably smooth continua to sharp lines.
A variety of target parameters have been measured so far using the spectrometer. The description below is divided up into sections covering each physical parameter that was determined. Each parameter relies on measurement of usually any of three spectral characteristics: the particle yield, the line width, or the absolute energy. Integrated particle yields determine fusion yield, core $\rho R$ (from knock-ons) and fuel ion temperature (from the ratio of fusion yields), while Doppler line widths of primary fusion products give a second measurement of fuel ion temperature. The absolute energy of fusion products, or, more specifically, a shift from their birth energy, are a probe of ablator $\rho R$’s and temperatures, as well as corona conditions (in particular the electrostatic potential of the target) at burn time. Some of these physical parameters are inter-related by virtue of the spectral characteristic on which they rely. In particular, determining the ablator $\rho R$ from a downshift requires knowing how much the electrostatic potential has upshifted the particle energies. For this reason, corona physics is described before ablator $\rho R$ measurements. Also, measuring the core $\rho R$ requires knowing how accurately the particle yields can be measured. Consequently, the fusion yield studies are described first.

6.1 Fusion yields

There are a number of important issues that need to be addressed using the yield results from various fusion products. These issues can be divided up into three parts: (1) *Comparisons with neutron measurements*. This provides a valuable external check of the CPS results. (2) *Comparisons of associated charged fusion products*, such as D-D protons and tritons, whose yields should be identical. This gives an internal consistency check of the spectrometer. (3) *Establishment of the degree of symmetry/asymmetry of particle fluxes*. This allows the true uncertainty in any single yield measurement to be established. All these results are summarized in Table 6.2.

The ratio of D-D proton to D-D neutron yields for multiple shots is shown in Figure 6.4 and the mean value, using both spectrometers, is $1.035 \pm 0.064$. From Figure 3.1, this ratio should vary between 0.91 and 0.98 over the temperature range $5 - 15$ keV, though it is unclear exactly how well those factors have been established. It is also possible that the neutron results, measured using indium activation, have errors due to unaccounted for scattering effects. The most striking aspect of these data is the considerable, non-statistical scatter that exists. Error bars for each point are approximately $\pm 0.05$, which are determined by the errors due to finite counting statistics for both the protons and neutrons. The standard deviation of this scatter is $0.277 \pm 0.032$ and is substantially larger than can be explained by errors in individual data points. A likely explanation is that the proton fluxes are anisotropic and this issue will be addressed later in the discussion about symmetry.
Figure 6.4: Ratio of D-D proton to D-D neutron yields for both spectrometers over various yields. CPS-1 and CPS-2 data are represented by diamond and triangle symbols respectively. The mean ratio of these data is $1.035 \pm 0.064$. Error bars shown are $\sim 5\%$ and are determined by counting statistics in the proton and neutron measurements. The scatter in the data has a standard deviation of 27.7\% and is considerably greater than can be accounted for by statistical errors.

The ratio of D-T alpha to neutron yields is found to be $0.925 \pm 0.087$, though this has only been measured for 5 shots, all on CPS-1. In principle, the comparison of D-T alpha and neutron yields contains less inherent uncertainty than the comparison of D-D proton and neutron yields. This is because the ratio of alpha to neutron numbers should be exactly one, and the measurement of 14.1 MeV neutrons, using copper activation, is less susceptible to scattering errors than the indium measurements. Once again, the standard deviation of these yield measurements, at $0.194 \pm 0.028$, is much greater than can be explained by statistical uncertainties.

The next step is to compare the yields of associated charged fusion products. The yield ratio of D-D tritons to protons is found to be $1.056 \pm 0.064$ while that of D-$^3$He alphas to protons is $0.972 \pm 0.052$, both of which are within experimental uncertainties of 1.0. This confirms that no systematic errors are detectable (at current uncertainty levels) from measurements of different particles at different energies. These comparisons are taken between measurements from the same spectrometer to avoid issues associated with asymmetric fluxes; however, there is still considerable scatter amongst the data: the standard deviation for the D-D measurements is $0.221 \pm 0.045$ while for the D-$^3$He measurements it is

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7 Although this may seem trivial, there are legitimate concerns regarding the counting of tracks in certain energy ranges. For instance, the D-$^3$He protons are very penetrating particles that must be ranged down before they can be detected by the CR-39. This ranging process causes a spread in the diameter distribution and creates the concern that some particles are being completely ranged out and lost.
0.146 ± 0.028. With more data, it might be interesting to investigate whether this scatter is greater when comparing associated particles from different spectrometers.

<table>
<thead>
<tr>
<th>Comparison with neutron measurements</th>
<th>Ratio</th>
<th>Error (Ratio)</th>
<th>σ</th>
<th>Error (σ)</th>
<th># yields measured</th>
</tr>
</thead>
<tbody>
<tr>
<td>D-D proton / D-D neutron (both CPS)*</td>
<td>1.035</td>
<td>0.064</td>
<td>0.277</td>
<td>0.032</td>
<td>19</td>
</tr>
<tr>
<td>D-T alpha / D-T neutron†</td>
<td>0.925</td>
<td>0.087</td>
<td>0.194</td>
<td>0.028</td>
<td>5</td>
</tr>
</tbody>
</table>

<table>
<thead>
<tr>
<th>Comparison of charged fusion products</th>
<th></th>
<th></th>
<th></th>
<th></th>
<th></th>
</tr>
</thead>
<tbody>
<tr>
<td>D-D triton / D-D proton (same CPS)</td>
<td>1.056</td>
<td>0.064</td>
<td>0.221</td>
<td>0.045</td>
<td>12</td>
</tr>
<tr>
<td>D-³He alpha / D-³He proton (same CPS)</td>
<td>0.972</td>
<td>0.052</td>
<td>0.146</td>
<td>0.028</td>
<td>8</td>
</tr>
</tbody>
</table>

<table>
<thead>
<tr>
<th>Comparison of CPS-1 and CPS-2</th>
<th></th>
<th></th>
<th></th>
<th></th>
<th></th>
</tr>
</thead>
<tbody>
<tr>
<td>D-D proton (CPS-1) / D-D proton (CPS-2)</td>
<td>1.015</td>
<td>0.103</td>
<td>0.273</td>
<td>0.035</td>
<td>7</td>
</tr>
<tr>
<td>D-³He proton (CPS-1) / D-³He proton (CPS-2)†</td>
<td>1.083</td>
<td>0.057</td>
<td>0.181</td>
<td>0.044</td>
<td>10</td>
</tr>
</tbody>
</table>

Table 6.2 A comparison of the various primary yields that have been measured. These are divided into three categories: comparisons with neutron measurements (which are an external consistency check), comparisons of fusion products from the same reactions (which are an internal consistency check), and comparisons between CPS-1 and CPS-2 (which are also an internal consistency check but, in addition, directly measure the yield symmetry). Errors are determined from the finite number of yields that have been measured with equal weighting for each data point. The standard deviation, σ, is much larger than the typical error on a given data point, indicating non-statistical variations in particle fluxes. Unless stated below, data has been compiled from all shots analyzed shots since Jan 1998. (*) Only shots after February 1998 were used here. During January and February, the Indium yield calibration procedure was being modified at LLE and the new procedure was not fully functional. (†) As of this writing, there have been no D-T alpha yield measurements using CPS-2. (‡) Shot # 13803 was removed from this calculation as it showed a particularly low value (0.3) that was clearly an outlier from the other data points.

If the non-statistical fluctuations from shot to shot are purely due to flux anisotropies, variations should be seen between yields from the two spectrometers on the same shot. This is exactly what is seen. Before looking at fluctuations though, the first question to answer is whether there is any systematic asymmetry that appears consistently over many different shots. As shown in Table 6.2, the mean ratio of D-D proton yields of CPS-1 compared to CPS-2 is 1.015 ± 0.103 while the similar ratio of D-³He protons is 1.083 ± 0.057. Thus, any asymmetries that occur on a given shot are averaged out over many different shots, indicating that the anisotropy is due to random fluctuations.

Without much more analysis and a working model it is difficult to say much more about these fluctuations though a semi-quantitative discussion is perhaps in order. Existence of these non-statistical variations is unexpected. The distribution of these fluctuations, for the D-D proton to neutron ratio, is plotted in Figure 6.5 as a histogram of normalized yield difference and, with more data, it will be
interesting to determine the ultimate statistical distribution of these variations. Until then, the simplest measure of these variations is the standard deviation, $\sigma$, listed in Table 6.2. Comparing $\sigma$ of the yield ratio between different spectrometers for D-D protons ($0.273 \pm 0.035$) versus D-$^3$He protons ($0.181 \pm 0.044$) indicates that the lower energy D-D protons show greater fluctuations, indicating greater anisotropies. A simple conjecture may be that the more penetrating particles are less susceptible to the deflections that would cause asymmetric fluxes. This is illustrated schematically in Figure 6.6 where the distributions are taken to be Gaussian.

It is generally thought (and apparently some tests have shown\textsuperscript{8}) that the neutron flux is isotropic. However, analysis of these preliminary data indicate that this may not be case. Comparing the $\sigma$ for the D-D proton to neutron ratio ($0.277 \pm 0.032$) and for the D-D proton ratio between different spectrometers ($0.273 \pm 0.035$) shows remarkable agreement. This is inconsistent with neutron isotropy: If the neutron yields are completely isotropic (and the fluctuations are thus due purely to charged-particle anisotropies), the $\sigma$ for proton to neutron comparisons should be less (by a factor of $\sqrt{2}$) than that for proton to proton comparisons from different spectrometers. In fact, this result is consistent with the notion that neutrons also exhibit the same degree of anisotropy, though it is hard to imagine how this may occur. Perhaps the measurement technique itself is susceptible to fluctuations. More data will be required to study this issue.

In summary, no major surprises have been found when examining the mean yields for various fusion products, though the D-D proton to neutron yield ratio is slightly higher than would be expected. Where they exist, mean yields are symmetric between CPS-1 and CPS-2. In addition, the yields for associated fusion products are identical, to within experimental uncertainties. Unexpected, though, is the discovery of seeming anisotropies in the particle fluxes. These fluctuations show up as non-statistical scatter of $20 - 30\%$ over multiple yield measurements. In other words, any single yield measurement cannot determine the “true” yield to better than $20 - 30\%$, no matter how good the statistics (which in these experiments was typically better than $5\%$). This fact is important when single yield measurements are required, such as in determining core $\rho R$ on a particular shot. The presence of this anisotropy places a limit on the attainable accuracy of any measurements requiring absolute yields.

\textsuperscript{8} Dr. Robert Kremens, private communication; V. Glebov, private communication.
Figure 6.5: A histogram of the normalized yield difference between D-D protons and neutrons for several shots. Results from each spectrometer are treated as separate measurements. The data are centered near zero but have a wide distribution that cannot be explained by statistical uncertainties. A normal distribution with the same standard deviation as the results is over-plotted to indicate how the results might look with improved statistics (if the fluctuations do turn out to be Gaussian).

Figure 6.6: Possible distributions of statistical fluctuations for different particles. Although the Gaussian distribution is assumed, it is plotted with the measured $\sigma$ for each particle. The D-$^3$He proton (at 14.7 MeV) has $\sigma = 0.181 \pm 0.044$, whereas for the D-D proton (3 MeV) $\sigma = 0.273 \pm 0.035$. The supposedly isotropic neutron flux is represented by a delta function. The larger the $\sigma$, the larger is the amplitude of the fluctuations and thus the more anisotropic is the particle flux. Indications are that the more penetrating particles are less susceptible to deflections and thus have a narrower distribution.
6.2 Fuel ion temperatures

Ion temperatures in the fuel have been measured using both the Doppler width of fusion product lines as well as the yield ratio of different fusion reactions. Just as in the case of primary yield measurements, ion temperatures are already determined by well-established neutron techniques, and, using charged-particles to achieve the same goal is a check of charged-particle methods. It should be emphasized that the ion temperature measurements described here are all volume and time-averaged over the fusion burn region, which is localized around the center of the imploded target.

Before discussing the temperature measurements themselves, it is useful to examine the shape of a fusion product line, such as that of D-D protons (Figure 6.7). The fact that this spectrum is well-approximated by a Gaussian indicates that the plasmas are in thermal equilibrium and the temperature measurement techniques described in Section 3.2 are valid. The spectrum of D-\(^3\)He protons is also distinctly Gaussian, as can be seen in Figure 6.8, though the lower dispersion of the magnet at 15 MeV does not provide as high a resolution as that at 3 MeV. It must be pointed out that, although Gaussian line shapes are not always observed, it is likely that the deviations arise from effects other than non-Maxwellian plasmas, which would affect all primary fusion product spectra. Two effects, namely the ranging down of particles, and disturbances due to coronal fields, appear to be responsible for the observed distortions.

The various fusion products that have been detected allow a number of different Doppler widths and yield ratios to be extracted. Of particular interest is the energetic D-\(^3\)He proton which is able to penetrate the thicker shells relevant for high-gain ICF targets. The temperatures determined from Doppler widths of D-\(^3\)He protons is shown in Figure 6.9. For thin shells (which generally means \(\rho R < 0.7 \text{ mg/cm}^2\), or about 2.5 \(\mu\)m of glass), there is rough agreement with the D-D neutron measured Doppler temperature, though the D-\(^3\)He measured temperatures appear to be reading 20 – 30 % higher. It is unclear why this is so. For thicker shells, the width of the protons acquires a significant non-Doppler component, presumably due to ranging through the compressed ablator. Even after only losing 5 % of its energy, this broadening can be detected. Scatter in the data at a given initial shell \(\rho R\) is probably due to a combination of measurement errors and variations in final shell \(\rho R\), which is ultimately what determines the amount of ranging, and thus broadening.
Figure 6.7: Spectrum of D-D protons from Shot 13825. The distribution closely follows a Gaussian profile, providing strong evidence that the reacting ions are in thermal equilibrium. This is further supported by the fact that, when particles undergo negligible ranging, the temperatures inferred from the width of these lines agrees with temperatures derived from the ratio of D-D to D-3He reactions.

Figure 6.8: The spectrum of D-3He protons is also closely approximated by a Gaussian – the lower resolution of these lines compared to that of the D-D proton above is due to the smaller dispersion at higher energies. Before any ranging occurs, the line width is determined by Doppler broadening [95]. With only 5% energy loss (10% in the case shown here), increased line broadening can already be detected.
When ranging occurs, a number of factors could contribute to a line width in excess of that due to pure Doppler broadening. Possibilities include energy straggling, ranging through inhomogeneous matter, or the presence of a distributed burn source. Energy straggling is a consequence of the statistical fluctuations inherent in any ranging process. Inhomogeneous matter, arising from fuel and shell mixing would cause particles to range through differing amounts of material. The presence of a distributed burn source would allow particles to travel different path lengths through the target, producing a broadened energy spectrum, even with homogeneous material.

Calculating the line width requires removing instrument broadening caused by the finite size of the magnet collimator. Now the measured signal is a convolution of the input spectrum and the approximately boxcar response function of the slit. To avoid the difficulties of deconvolution, for now it is assumed that the measured spectrum is a convolution of an initial Gaussian of standard deviation $\Sigma$, with a boxcar function of width $w$, corresponding to the width, in energy units, of the collimated beam when it intersects the detector. The standard deviation, $\xi$, of the measured spectrum is then given by

$$\xi^2 = \Sigma^2 + \frac{w^2}{12}$$  \hspace{1cm} (6.1)$$

which allows $\Sigma$ (and thus the ion temperature) to be extracted from the measured $\xi$. Eq. (6.1) is a consequence of the fact that the probability density function, $h(x)$, of two random, independent probability distributions, $f(x)$ and $g(x')$ is given by the convolution of $f$ and $g$. Thus the variance of $h$ is the sum of the variances of $f$ and $g$. The term $w^2/12$ is the variance of the boxcar function of width $w$.

Using a 1 mm wide collimator, the line width of 15 MeV D$^3$He protons, born in fuel ion temperatures of 15 keV, is broadened by approximately 25%; thus it is crucial to remove the effects of slit broadening. For D-D protons and other lower energy particles on the other hand, the slit contributes less than 3% to the total variance. Clearly problems arise when the initial spectrum is not a Gaussian, as might be the case after significant energy loss. These distributions and others will be the source of future study.

The temperatures inferred from line widths of D-D protons and tritons, as well as D$^3$He alphas compared to the neutron measured Doppler width are shown in Figure 6.10. As with the 15 MeV protons, for shells with an initial $\rho R$ less than 0.7 mg/cm$^2$, the line widths of these particles can be explained by Doppler broadening effects alone. For all particles then, it appears that below 0.7 mg/cm$^2$, ranging processes are unimportant. In other words, shells any thinner than this are completely ablated off by the time the fusion products are generated. A few data points in Figure 6.10 for the D-D protons seem to indicate that increased shell $\rho R$ broadens the line width just as with the D$^3$He protons.
Figure 6.9: The ion temperature of D-3He fuel can be measured using either the yield ratio of D-3He protons to D-D neutrons (or protons), or the Doppler width of the D-3He protons. Results from these two methods are normalized to results from a neutron time-of-flight detector. For thin-shelled targets (initial $\rho R < 0.7 \text{ mg/cm}^2$), temperatures inferred from the Doppler width of D-3He protons are in rough agreement with neutron-determined values; however, for thicker shells, the line width is dominated by other broadening effects and is no longer an accurate temperature measurement. Temperatures inferred from the yield ratio of D-D neutrons to D-3He protons are in rough agreement with neutron results, irrespective of target shell thickness. The temperature range covered by these points (as determined by neutron Doppler and D-3He ratio methods) is 4 – 15 keV, with the highest temperatures achieved using the thinnest-shelled targets [95].

Figure 6.10: Comparison of line width temperature measurements using D-D protons, D-D tritons and D-3He alphas. Once again, the line widths are dominated by Doppler broadening when thin-shelled targets are used. The tritons and alphas are quickly ranged out by thicker shells so it is difficult to how they broaden with ranging. Only limited data for the D-D protons exists for higher $\rho R$ shells but they appear to broaden more rapidly than the higher energy, D-3He protons, as would be expected.
There is considerable scatter in the results for these lower energy particles, and this is probably due to the inherent errors in determining the line width. With greater dispersion at lower energies, it is more likely that a portion of the spectrum will be lost between adjacent pieces of CR-39. The missing spectral portion can cause errors in line width measurements, leading to temperature errors of about 30%.

These uncertainties are much greater than the statistical uncertainties in calculating the line width from a clear profile such as that in Figure 6.7, which are usually less than 5%. At higher energies, where dispersion is weaker, the line width (in physical space) is much narrower, making missed coverage less likely; however, the poorer energy resolution makes uncertainties associated with the finite energy interval of each bin more significant. These subtle issues of error analysis will need to be addressed in the future.

It must also be mentioned that, quite often, the line profiles of tritons and alphas have significant structure present on them. An example of an anomalous triton profile is shown in Figure 6.1. It is no doubt inappropriate to extract a Doppler width from such a distribution and, in the worst cases, such data has been excluded from the data shown in Figure 6.10; however, this structure can appear on many different levels and thus the scatter in the data is likely to reflect this. Perhaps even the D-D protons contain these anomalies as well.

Temperatures inferred from the ratio of D-\(^3\)He protons and D-D neutrons are shown in Figure 6.9 and indicate reasonable agreement with time-of-flight measurements of the neutron Doppler width. This is significant since the ratio method relies on the measured fusion cross sections, whereas the Doppler method does not – the two techniques provide independent verification of ion temperature. The agreement is maintained over a range of different capsule thicknesses, as would be expected. A closer look at these results however, indicates that there is a systematic discrepancy that arises when low D-\(^3\)He yields are measured (see Figure 6.11). For yields above 10\(^{10}\), there is excellent agreement between the two measurement techniques, but deviations appear to steadily increase at lower yields. At present, counting errors of the magnitudes necessary to produce this disagreement have been ruled out. It is also unlikely that neutron measurement errors are responsible since the D-D neutron yields vary by only a factor of ~5, compared to the D-\(^3\)He yields, which vary by a factor of 50 or more.
Figure 6.11: A closer look at the comparison between temperatures determined from the ratio of D-D to D-^3^He yields and those determined from neutron Doppler widths. While there is excellent agreement for D-^3^He yields greater than 10^10, there appears to be a systematic discrepancy at lower yields. The reason for this behavior is uncertain.

Figure 6.12: In D-T fuel, a particular yield ratio of D-T alphas to D-D protons is expected at any given temperature (solid line); however, at these temperatures, the turning point in the predicted ratio makes it impossible to extract an accurate temperature value from measured D-T and D-D yields. By comparing the temperature determined from the D-D proton Doppler width though, the measured ratios are shown to be consistent with the expected values.
In the case of D-T fuel, it is, in principle, possible to extract an average fuel temperature from the ratio of D-T to D-D reactions. Only four shots have been used to investigate this however, and all of them have had temperatures in the 15 keV region, where the yield ratio curve is almost independent of temperature. This makes it very difficult to extract an accurate temperature. In addition, no neutron-measured temperatures were available because of the higher D-T yields which saturated the detectors mounted at the time. What has been plotted in Figure 6.12 then is the D-T to D-D yield ratio versus the temperature inferred from the Doppler width of D-D protons. These results are shown to be consistent with the expected yield ratios determined by the reaction rate coefficients. On the other hand, the line width of D-T alphas was found to be anomalously broad for all four shots and inferred Doppler temperatures disagreed with those from the D-D protons by a factor of almost two. This broadening is reflected in the distorted alpha profiles, an example of which was given in Figure 6.2. It is possible that such distortions are due to time variations of the coronal fields. The fact that some of the D-3He alpha profiles do not show the same broadening, even though the particle energies are similar, may be a consequence of the D-3He burn time being shorter (and thus less sensitive to time variations) than the D-T burn time.

In summary, ion temperature measurements have been made using a variety of fusion products, utilizing both the Doppler width and yield ratio methods. Results from both techniques are within 20 – 30 % of each other in the asymptotic limit of thin-shell targets, where most, if not all, of the shell is ablated by the laser. The width of the spectral lines acquires a significant non-Doppler component when thicker shells are used, and a number of ranging effects may be responsible. Temperatures inferred from the ratio of D-3He to D-D reactions agree with D-D neutron measured Doppler temperatures at high yields but systematically underestimate them at low D-3He yields.

6.3 Core $\rho R$

Core $\rho R$ has been determined by measuring the spectra of deuteron and triton knock-ons elastically scattered by D-T neutrons (see Figure 6.2). This is the first time the characteristic knock-on spectra have been resolved and it is notable how well the distributions, and their endpoint energies, follow those predicted from the neutron differential scattering cross-sections (compare Figure 3.4). As an example of how these spectra are used to extract core $\rho R$, one particular shot will be analyzed in detail.

Shot # 13817 used a 916 $\mu$m diameter target filled with 10 atm of deuterium and 10 atm of tritium, with a shell thickness of 5.0 $\mu$m CH plastic on top of 2.4 $\mu$m glass. The laser energy was 27.1 kJ, delivered in a 1-ns square pulse and the D-T neutron yield was $2.67 \times 10^{13}$. Knock-on spectra were measured both by CPS-1 (using a 10 mm wide collimator slit) and CPS-2 (using a 1 mm wide slit). Results from CPS-2, which had higher counts, better resolution, and lower signal-to-noise, will be emphasized here.
Knock-on spectra from shot \#13817 are shown in Figure 6.13, with the neutron background subtracted. Background levels are determined by scanning the CR-39 out beyond the high energy peak, in a region that contains purely noise events. It is evident that this region contains only noise since the tracks are not restricted to particular diameter bands, as would be the case for charged particles dispersed by the magnet. By assuming that the noise distribution is independent of position, this "pure background" can then be subtracted from the rest of the spectrum, taking care to only select noise events that are in the diameter range occupied by signal events. The number of knock-ons within the high-energy "ramp" feature, as determined from the scattering cross-sections, given in Table 3.2 and repeated here is:

\[
Q_{\text{RAMP}}^{\text{KO}} \text{(Deuterons, } E > 9.6 \text{ MeV)} = 0.0125 \langle \rho R \rangle Y_{DT} \\
Q_{\text{RAMP}}^{\text{KO}} \text{(Tritons, } E > 7.3 \text{ MeV)} = 0.0170 \langle \rho R \rangle Y_{DT}
\]

where \(\langle \rho R \rangle\) is in g/cm\(^2\). The yield of knock-ons within this feature is \(1.93 \times 10^9\) for deuterons (where \(E > 9.0\) MeV) and \(3.68 \times 10^9\) for tritons (where \(E > 6.5\) MeV). Note how the measured energy limits have been downshifted by ranging. Using these values, and the measured neutron yield, the fuel \(\rho R\) calculated using deuterons and tritons is found to be 5.8 mg/cm\(^2\) and 8.1 mg/cm\(^2\) respectively. Taking the average of these two numbers as the estimated \(\rho R\), and knowing that any single yield measurement has a 25% error (there are two measurements in this case), \(\rho R\) is found to be \(7.0 \pm 1.2\) g/cm\(^2\).

A one-dimensional hydrodynamic simulation (courtesy of Radha Bahukutumbi) of the implosion predicts a burn-averaged \(\rho R\) of 11 g/cm\(^2\) and a neutron yield of \(4.13 \times 10^{13}\). The expected knock-on spectra using this simulation are shown in Figure 6.13. Usually the 1-D code simulations predict \(\rho R\)'s several factors higher than actually achieved. This is because 1-D codes, by definition, neglect the Rayleigh-Taylor instability which always reduces compression and yield. A summary of results for shot \#13817 are given in Table 6.3.

<table>
<thead>
<tr>
<th>Shot #13817</th>
<th>Neutron Yield</th>
<th>Knock-On Yield in high energy &quot;ramp&quot;</th>
<th>(\rho R) (mg/cm(^2))</th>
</tr>
</thead>
<tbody>
<tr>
<td></td>
<td>Experiment</td>
<td>Deuterons (E &gt; 9.0 MeV) Tritons</td>
<td>7.0 \pm 1.2</td>
</tr>
<tr>
<td></td>
<td>Simulation</td>
<td>5.21E9 (E &gt; 8.2 MeV) 6.57E9 (E &gt; 5.8 MeV)</td>
<td>11</td>
</tr>
</tbody>
</table>

Table 6.3: Summary of knock-on results for shot \#13817. The knock-on yields have been determined only in the high-energy ramp feature. The fact that the simulated spectrum shows more downshift than the measured means that the energy interval occupied by this ramp is different for the two cases (as is indicated) [94].

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Figure 6.13: The spectrum of deuteron and triton knock-ons from Shot # 13817, a 916 μm diameter target with a 2.4 μm glass + 5 μm CH thick shell, filled to 20 atm with 1:1 D-T gas. The neutron background, determined by the signal level measured in the region above the maximum energy of the knock-ons, has been subtracted. Spectra expected from a 1-D hydrodynamic simulation of the implosion are also shown. Both the core \( \rho R \) (determined by the absolute yield of knock-ons) and the ablator \( \rho R \) (determined by the downshift in the high energy knock-on peaks) are over-estimated by the 1-D simulations, as might be expected [94].

It is instructive to examine the simulated spectra by using Eqs.(6.2) to check how accurately the simplified approach inherent in this equation approximates the knock-on generating process. Note that these equations were based on the “central-burn model”, the assumption that all neutrons are generated at target center and thus get a chance to sample the entire target radius in flight. In the case of the “uniform-burn model” where the neutrons are generated equally throughout the fuel, these equations are modified by a factor of 0.75. For a generalized radial burn profile, a corresponding form factor, \( g \), of order 1 can be determined and the equation may be written:
where $Q(\text{real})$ is the actual number of knock-ons generated and $Q(\text{central})$ is the knock-on yield found from Eqs. (6.2). For a central burn, $g = 1$, while for a uniform burn, $g = 0.75$. The form factor for the simulated spectra is found to be $\sim 0.9$, indicating that the central burn approximation is suitably accurate.

While it is the yield of knock-ons that determines core $\rho R$, the position and shape of their spectra contains information as well. As can be seen, the downshift is predicted to be greater in the simulations than was measured. This indicates that the ablator $\rho R$ (which is where most of the energy is lost) was less than predicted.

A number of other shots have been studied to look for knock-ons. Five of the eleven shots analyzed have produced detectable knock-ons, indicating that the others had $\rho R$'s below $\sim 1 \text{ mg/cm}^2$. Table 6.4 lists all the shots studied so far. As might be expected, the knock-ons have been found mainly on the thicker shell targets, which maintain better integrity during the implosion and thus produce higher $\rho R$'s. There is considerable variability in $\rho R$ among similar targets however and a systematic study would be required to establish these trends more precisely. With the knock-on technique now established, such studies will be undertaken in the near future.
### Table 6.4: D-T shots that have so far been analyzed to look for knock-ons

<table>
<thead>
<tr>
<th>Shot #</th>
<th>Shell Thickness</th>
<th>Pressure (atm)</th>
<th>D-T Neutron Yield</th>
</tr>
</thead>
<tbody>
<tr>
<td>11739</td>
<td>2.3 μm glass</td>
<td>5</td>
<td>1.01E13</td>
</tr>
<tr>
<td>11743</td>
<td>2.3 μm glass</td>
<td>10</td>
<td>3.40E13</td>
</tr>
<tr>
<td>11757</td>
<td>2.0 μm glass</td>
<td>10</td>
<td>1.97E13</td>
</tr>
<tr>
<td>12152*</td>
<td>20 μm CH</td>
<td>17</td>
<td>5.39E12</td>
</tr>
<tr>
<td>12933</td>
<td>2 μm glass + 5 μm CH</td>
<td>10</td>
<td>1.60E13</td>
</tr>
<tr>
<td>12934</td>
<td>2 μm glass + 5 μm CH</td>
<td>10</td>
<td>2.00E13</td>
</tr>
</tbody>
</table>

Knock-on D and T were NOT detected on the following shots

Knock-on D and T WERE detected on the following shots

<table>
<thead>
<tr>
<th>Shot #</th>
<th>Shell Thickness</th>
<th>Pressure (atm)</th>
<th>D-T Neutron Yield</th>
</tr>
</thead>
<tbody>
<tr>
<td>11742</td>
<td>2.6 μm glass</td>
<td>10</td>
<td>4.45E13</td>
</tr>
<tr>
<td>12938</td>
<td>2 μm glass + 10 μm CH</td>
<td>10</td>
<td>3.40E12</td>
</tr>
<tr>
<td>13816†</td>
<td>2 μm glass + 5 μm CH</td>
<td>20</td>
<td>3.78E13</td>
</tr>
<tr>
<td>13817†</td>
<td>2.4 μm glass + 5 μm CH</td>
<td>20</td>
<td>2.67E13</td>
</tr>
<tr>
<td>13822†</td>
<td>2.3 μm glass + 7.5 μm CH</td>
<td>20</td>
<td>2.31E13</td>
</tr>
</tbody>
</table>

(*) The shell on shot # 12152 was probably too thick to allow most knock-ons to emerge. (†) Only on these shots was CPS-2 operational; however, knock-ons were seen on both spectrometers.

### 6.4 Ion accelerations and corona conditions

The experimental conditions for these studies were the same as for others – they were, after all, the same shots. Since the laser conditions are of central importance in the case of coronal physics, the laser and target parameters will be listed again. These experiments utilized 60 beam uniform illumination of spherical targets with 0.35 μm light. Laser energies between 20 and 30 kJ at average on-target intensities of $10^{14} - 10^{15}$ W/cm² were delivered over 1 ns. Mostly square pulse shapes were used, with a few Gaussian pulses. Targets, approximately 1 mm in diameter and filled with D-T, D-D or D-³He fuel, had glass, plastic on glass, or pure plastic ablator materials ranging from 2 up to 20 μm in thickness.

#### 6.4.1 Accelerated ablator protons

Ablator protons, with a continuum of energies up to as high as 1.4 MeV, have been observed. Two examples are shown in Figure 6.14. These spectra exhibit a variety of features, displaying both relatively smooth continua as well as sharp lines. In some cases, particularly intense, discrete lines are found, as shown in Figure 6.15. A feature common amongst all the spectra is a sharply-defined, maximum (endpoint) energy. This endpoint scales with average, on-target laser intensity as shown in Figure 6.16,
with a threshold for proton acceleration at about $10^{14} \text{ W/cm}^2$. The presence of SSD (smoothing by spectral dispersion) does not have any noticeable effect on the maximum proton energy. For the few shots where the spectra have been measured using both spectrometers, the endpoint energies are identical, to within experimental error, though the distributions themselves can be quite different (Figure 6.17). In addition to protons, some other accelerated ablator ions have been identified, in particular deuterons and $^3\text{He}$, as well as what appears to be different charge states of carbon or oxygen; however, it is the protons that are most accessible with this spectrometer.

The intensity threshold for ablative proton acceleration, at $10^{14} \text{ W/cm}^2$ (as shown in Figure 6.16) is consistent with the source of the hot electrons being the two-plasmon instability (see Table 4.1). The onset of this instability occurs at a similar intensity, while stimulated Raman scattering (SRS) begins at higher intensities. Furthermore, the insensitivity of the endpoint energies to beam smoothing by spectral dispersion (SSD) agrees with previous work which showed that two-plasmon decay is largely independent of SSD [65] (unlike SRS).

If the proton endpoint energy is used as a measure of the target potential, it can be concluded that the target is charging up to $\sim 1 \text{ MV}$ for laser intensities of $10^{15} \text{ W/cm}^2$. Since it is the highest energy protons that are produced first, as discussed in Appendix A, this represents the potential at the start of the laser pulse. Symmetry of the endpoint energies is also an indication that the highest energy protons are being accelerated at the beginning of the laser pulse: rapid motion of the hot electrons is likely to produce spherically symmetric fields around the target. Soon after this potential has been created though, it is likely that return current will begin to flow back towards the target from surrounding structures (in particular the stalk, which holds the target), breaking the spherical symmetry of the system. The spectrum of protons at lower energies, generated at later times, are more likely to experience these asymmetries and this could explain the different distributions from CPS-1 and CPS-2 depicted in Figure 6.17.

While the measured accelerations of ions can directly determine the target potential, the more fundamental parameter of interest is the hot electron temperature. Using the static field description of particle acceleration described in Appendix A, a hot electron temperature can be inferred from the endpoint energy of ablator protons. In this model, the potential of the target, and thus what the ion front sees, is determined by the number of hot electrons that escape from the confines of the target. From Figure 4.2, assuming a total hot electron energy of 30 J, measured potentials of 0.5 – 1 MeV indicate hot electron temperatures of 45 – 110 keV. These values are probably overestimates owing to the assumptions, inherent in the model, that the electrons have to completely escape from the target.
Figure 6.14: Two examples of ablator proton spectra [93]. The spectra often show both lines and continua, though the strongest lines tend to be constrained to the lower part of the spectrum. Endpoint energies up to 1.4 MeV have been observed.
Figure 6.15: Occasionally, extremely intense lines appear in the ablator proton spectrum [93]. The lines shown here are particularly intense; usually they show less contrast with the background continua.

Figure 6.16: The proton endpoint scales with laser intensity with an apparent threshold at $10^{14}$ W/cm$^2$ [93]. CPS-1 data are diamonds, CPS-2 are triangles. This threshold and the fact that SSD has no apparent affect on the endpoint energy indicates that the hot electrons are being produced by two-plasmon decay. Some of the scatter in the data may be due to different target conditions: targets with a thin coating of plastic on the outside appear to produce higher energies than just pure glass or thick plastic shells.
Figure 6.17: Ablator proton spectra measured simultaneously on CPS-1 and CPS-2 indicate that the endpoint energies are symmetric; however, the shapes of the spectra can have noticeable differences [93].

Hot electron temperatures can also be determined, somewhat crudely at this stage, using the slope of the ablator proton velocity spectrum. The planar, isothermal, self-similar expansion model described in Appendix A, predicts an ion velocity distribution given by \( \exp(-V/C_s) \), where \( V \) is the particle velocity, \( C_s \) is the ion sound speed given by \((T_e/m)^{1/2}\), and \( T_e \) is the hot electron temperature. Note that on some spectra (such as that in Figure 6.15), it would be inappropriate to use this model since the distributions do not follow this exponential decay. On shot 13402, data from CPS-2 does appear to follow this
distribution, though CPS-1 measurements deviate slightly (Figure 6.18): Comparing data from both spectrometers provides an estimate of the measurement uncertainty. The slopes of the velocity spectra give hot electron temperatures of 40 keV for CPS-1 and 15 keV for CPS-2. In comparison, the endpoint of this spectrum, at 0.85 MeV, leads to an inferred temperature of 85 keV (from Figure 4.2), assuming 30 J of hot electron energy. While there is more than a factor of 2 discrepancy between the estimates of the two methods, they both indicate temperatures in the tens of keV range, rather than the few keV range that has been observed for this $I_\lambda^2 = 10^{14} \text{ W/cm}^2$. These values are more in line with hot electron temperatures measured by Rousseaux et al. [68] for similar laser conditions.

The discrete line structure displayed in Figure 6.15 does not appear to be repeatable: both the line spacings and positions vary from shot to shot. Furthermore, the lines have yet to be repeated from one spectrometer to the next, and, generally speaking, if they appear on CPS-2, they have a much closer spacing than ever seen on CPS-1. The mechanism that produces such features is still a mystery, though R. Short, in his recent presentation at the Conference of the American Physical Society [96], proposed a theory based on a periodically discharging target releasing particles at decreasing energies.

6.4.2 Accelerated fusion products

On thin-shelled targets (2.0 – 2.8 $\mu$m glass), acceleration of fusion products by up to 0.7 MeV per unit charge has been observed. Elevated energies have been seen for all observed fusion products so far: D-D protons and tritons, D- $^3$He protons and alphas, and D-T alphas, and the upshifts are symmetrical to within experimental uncertainties ($\pm$ 50 keV), when viewed from the two spectrometers. These energy upshifts increase with laser energy (or intensity) as shown in Figure 6.19 for the case of D-D protons but are always two or more times less than the simultaneously-measured endpoint energy of the ablator protons. In the case where negligible ranging occurs in the shell (presumably because the shell is completely ablated off), the accelerations of different fusion products are degenerate with the charge $Z$ as shown in Figure 6.20: alphas ($Z = 2$), receive twice as much energy as protons or tritons ($Z = 1$). This is consistent with the picture of electrostatic accelerations. As yet, this has only been observed on a single shot, as all other D- $^3$He targets had greater shell thicknesses, which caused some particle ranging.
Figure 6.18: Velocity distributions of ablator protons shown in Figure 6.17. Taking the slope of these spectra allows the hot electron temperatures to be calculated from the self-similar, one-dimensional expansion model. This leads to temperatures of 40 and 15 keV from CPS-1 and CPS-2 respectively. The discrepancy between the two values is a rough measure of the uncertainty.
Figure 6.19: Fusion products on thin-shelled targets show several hundred keV upshifts in their mean energies [93]. For similar shell thicknesses, these shifts increase with laser energy (intensity).

Figure 6.20: On this target, where the shell was 2.3 μm glass, irradiated with 28.4 kJ, the particle upshifts showed a degeneracy with Z, consistent with the picture of electrostatic accelerations [93].
While the maximum ablator proton energy is a measure of target potential at or near the beginning of the laser pulse, the acceleration of fusion products is a probe of potentials at bang time (the time of fusion burn). The fusion products may be viewed as test particles: their population is several orders of magnitude less than that of the ablator ions and thus their contribution to the fields is negligible. They are in fact precisely gated probes of target potential conditions during the burn period. In thin-shelled, or exploding pusher, targets (2 – 2.5 μm thick shells), irradiated under the conditions described above, bang time occurs at around 700 ps after start of the laser pulse, as predicted by hydrodynamic calculations and measured by the neutron timing diagnostic. This means that the 1 ns pulse is still on when the fusion products emerge. The fact that the upshifts of fusion protons are always less than half the ablator proton maximum energy indicates that the target potential has decayed by a factor of two or more by the time of fusion burn.

For the thicker-shelled targets (~ 20 μm) relevant to future high-density experiments, the bang time occurs after, rather than during, the 1 ns laser pulse. It is reasonable to assume that the fields have decayed more, perhaps even disappeared, by this point. Experiments using shorter laser pulses (~ 400 ps) are currently being undertaken to test this assumption.

It is interesting to investigate the arrival times of the various fusion products at the detector. Shown in Figure 6.21 is a time-of-flight schematic for all of the particles, including an ablator proton front and the hot electrons. As can be seen, the D-3He protons and D-D protons arrive well ahead of the ion front, the D-3He alphas are almost concurrent with it, while the D-D tritons are well behind. Despite these timing differences, all particles appear to experience the same accelerating potential (based on the shot in Figure 6.20). This indicates that, as expected, all of the acceleration occurs very close to the target, when the fusion products are still grouped together and thus experiencing the same potential drop. It should be noted that the D-D tritons are the only fusion products that travel more slowly than the ablator protons of highest energy. This means that these tritons, unlike the other fusion products, have the same velocity as some of the accelerated ions. This may be the cause of the line structures that are sometimes observed on the spectrum of D-D tritons (Figure 6.22), which are never observed on faster particles like the D-D protons. Perhaps the tritons, velocity-matched to some of the ablator protons, are subjected to the same mechanism that produces discrete lines in some of the ablator proton spectra.
Figure 6.21: Time-of-flight to 100 cm for particles from D-3He shot 13786 whose spectra were shown in Figure 6.20. Only the tritons arrive behind the ablator ion front; the alphas arrive just ahead of the front.

Figure 6.22: Often the D-D triton spectra show a regular line structure similar to that observed on the ablator protons. The velocity of these tritons is in the same range as that of the ablator protons.

In summary, ablator proton energies have been measured up to 1.4 MeV and accelerations of the fusion products have been detected up to 0.7 MeV per unit charge. The maximum ablator proton energy indicates that the target is charging up to greater than 1 Megavolt at the beginning of the laser pulse; subsequent decay of this potential by bang time results in the fusion products always receiving less...
energy than the ablator proton endpoint. Hot electron temperatures appear to be in the range of 15 – 40 keV, as determined by the slope of the ion spectrum, or 45 – 110 keV, as deduced from the maximum ion energy. The occasional appearance of discrete, regular lines in the ablator proton spectra is still a mystery. The maximum ablator proton energies and fusion product upshifts appear so far to be symmetric from orthogonal views.

6.5 Ablator $\rho R$

The energy downshift of various particles, in particular D-3He protons and D-D protons, has been used to determine ablator $\rho R$'s from a few mg/cm$^2$ up to ~ 50 mg/cm$^2$. In addition, simultaneous measurements of D-D protons and tritons, along with D-3He protons and alphas, have been used to try to establish a consistent picture of the competing effects of ablator losses (as a function of $\rho R$ and $T$) and coronal accelerations on particle energy shifts.

In the most general case, total energy shift is due to losses in both the core and ablator, along with accelerations in the corona. For many of the conditions encountered, losses in the core can be neglected. This leaves three unknowns: the ablator $\rho R$ and temperature, and the electrostatic potential of the target. Extracting ablator $\rho R$ thus requires elucidating the effects of ablator temperature and coronal accelerations on the measured energy shift. Depending on the particular target conditions and the particle velocity, these parameters can have differing levels of importance.

The highly penetrating D-3He proton is the particle most relevant for diagnosing the thick-shelled targets necessary for high gain. Figure 6.23 shows spectra with mean energy shifts of 0.9 MeV and 1.7 MeV, corresponding to shell thicknesses of 14.6 µm of CH and 18.4 µm CH respectively. Assuming that there is no acceleration and that all energy losses occur in the ablator, the inferred $\rho R$ for these targets based on stopping calculations is 25 mg/cm$^2$ and 55 mg/cm$^2$ respectively, taking ablator temperatures anywhere less than 3 keV. The variation of $\rho R$ with ablator temperature for a given energy loss is illustrated in Figure 6.24. An advantage of using the D-3He proton is that its velocity is high enough above typical electron velocities in the ablator that $\rho R$ is independent of temperature (in the expected regime of 1 – 3 keV) for a given energy loss. This means that a single energy loss measurement can be used to determined $\rho R$.

The ability to extract ablator $\rho R$ from a single energy loss measurement relies on two assumptions: that acceleration effects and core energy losses are negligible. The upshifted spectrum from a target with only a 2.2 µm glass shell is shown for reference in Figure 6.23. Such upshifts can be neglected if the target potential has decayed back to zero by the time the fusion products are generated. As explained in the previous section, for thick-shelled targets where the bang time occurs after the laser pulse has ended, this
is a plausible assumption, though one that needs to be, and is being, tested. With a 1 ns square pulse, bang time for the thin-shelled target is \( \sim 0.7 \) ns after the start of the pulse, whereas for the thicker-shelled target it is at \( \sim 1.5 \) ns.

Typical core \( \rho R \)'s for these targets might be between \( 5 \) – \( 10 \) mg/cm\(^2\) at temperatures of \( 5 \) – \( 10 \) keV. Taking the case that would produce the greatest energy loss (\( 10 \) mg/cm\(^2\) at 5 keV), the predicted downshift for the 14.7 MeV proton is about 0.3 MeV. While this is certainly smaller than the measured shifts of 0.9 and 1.7 MeV discussed above, it is hardly negligible, and this should be kept in mind if measurement accuracies better than \( \sim 30 \% \) are expected.

As with all results from this spectrometer, these measurements are time-integrated and the inferred \( \rho R \)'s are to be understood as yield-weighted averages over the burn period. Target implosion simulations indicate that not only can \( \rho R \) vary during any single burn time, there can also be multiple burn periods corresponding to shock-generated heating (which usually occurs first) and compression-generated heating. Generally, the shock-yield occurs before maximum compression has taken place and so particles generated during this time will probe lower \( \rho R \)'s and exit with higher energies than particles produced near maximum compression. The resulting time-integrated spectrum would then have a broad profile, possibly with multiple peaks corresponding to the different burn times. A simulated spectrum that includes all these effects is shown in Figure 6.25, along with the measured spectrum for that shot. The discrepancy between calculation and experiment is quite apparent, both in mean energy and line width, and future work should focus on adjusting implosion parameters to model the observed spectra. An added complication might occur when 3-D effects play a role. Mixing of the shell and fuel material could produce an inhomogeneous mix which would cause increased broadening of the spectrum.
Figure 6.23: Greater downshifts in the D-³He proton spectrum are observed on targets with thicker shells [95, 97]. Shell thicknesses of 18.4 μm and 14.6 μm of CH result in proton energy losses of 1.7 MeV and 0.9 MeV respectively. These targets were filled with 15 atm of D-³He. Also shown is the upshifted spectrum from a thin-shelled glass target where the protons, undergoing negligible ranging, are accelerated by coronal fields. It is thought that the later bang times of the thick-shelled targets, occurring after cessation of the laser pulse and decay of the accelerating fields, means that fusion products from thicker targets do not experience any upshifts. Laser energies for all these shots were from 26 to 28 kJ.

Figure 6.24: The predicted ablator ρR versus T relationships for D-³He protons ranged down to 13.8 and 13.0 MeV, whose spectra were shown in Figure 6.23. For reference, the curve corresponding to ranging out of the 14.7 MeV protons is also shown. These calculations assume that the CH ablator, with an ion density of 2×10²⁴ cm⁻³, is responsible for all the energy losses, and that acceleration effects are negligible. The independence of these curves for temperatures around 1 – 3 keV (the expected ablator temperature range), means that the ablator ρR can be determined directly from a single energy loss measurement.
Figure 6.25: A comparison between the measured and predicted D-\(^3\)He proton spectra for Shot 13799 [94, 97] (Predicted spectrum courtesy of Radha Bahukutumbi). The higher expected \(\rho R\) of 120 mg/cm\(^2\), compared to a value of 42 mg/cm\(^2\) inferred from the measured spectrum, is reflected in the larger calculated downshift. Structure in the predicted spectrum illustrates the complexity of the implosion process: a shock-induced burn can precede the main burn period near maximum compression, and protons from each of these times will sample different shell \(\rho R\)'s. This is reflected in the broad, multi-peaked nature of the time-integrated spectrum.

The downshift of D-D protons has also been observed on several targets. Spectra with shifts of 0.4 and 0.9 MeV are shown in Figure 6.26 where shell thicknesses were 2.4 \(\mu\)m glass + 5.0 \(\mu\)m CH and 2.4 \(\mu\)m glass + 7.5 \(\mu\)m CH respectively. As a starting point, assume that accelerations in the corona and energy losses in the core can be neglected; this leads to inferred \(\rho R\)'s of 7 and 15 mg/cm\(^2\) respectively, taking an ablator temperature of 1 keV. The variation of \(\rho R\) with temperature for energy shifts of 0.4 and 0.9 MeV are given in Figure 6.27 and it is apparent that, unlike in the case of D-\(^3\)He protons, the inferred \(\rho R\) is sensitive to ablator temperatures in the range of 1 – 3 keV anticipated for these conditions. This is because the velocity of these lower energy protons is comparable to the electron velocities. As a result, for \(\rho R\)'s to be measured accurately, ablator temperatures also need to be determined, necessitating simultaneous measurements of at least two different particle spectra. Furthermore, because the shells in these shots are not as thick as those in the cases studied for the D-\(^3\)He protons, the bang time for these targets is going to be closer to the end of the laser pulse. This means that acceleration effects could play an important role, throwing a third unknown – target potential – into the mix, and requiring that at least three particle spectra be measured.

Since both these shots were filled with D-T fuel, the core \(\rho R\) could be measured from the yield of knock-ons as described in Section 6.3. For shots 13817 and 13822, core \(\rho R\)'s were found to be about 7 mg/cm\(^2\).
which, at a temperature of 5 keV would lead to losses of 0.15 MeV, as compared to the measured total losses of 0.4 and 0.9 MeV. This means that approximately 15 – 30 % of the energy was lost in the core.

The general procedure for extracting ablator $\rho R$, temperature and target potential using measurements of four simultaneous particle shifts has been used on only one shot so far. The most accessible particles for this study are the emissions from a D-3He shot: the D-3He protons and alphas, along with the D-D protons and tritons. In order to allow the tritons and alphas to emerge, a relatively thin shell must be used, which means that, for 1 ns laser pulse lengths, acceleration effects are significant. In addition, the lower velocities of the particles results in the stopping powers being sensitive to temperature in the vicinity of 1 keV. In the limit of a very thin shell that is completely ablated, no ranging occurs and the particles emerge with shifts dependent upon the target potential alone. This was shown in Figure 6.20 for the case of a 2.3 µm glass shell. It is clear that ranging effects are insignificant because the upshifts are degenerate with Z. For the case of a slightly thicker shell (2.7 µm glass), the shifts are no longer degenerate and now carry information about the $\rho R$ and temperature of the ablator (see Figure 6.28).

To compare these different shifts, a potential must first be assumed – a good starting point is the total upshift of the 14.7 MeV proton, or just above. With this potential, $\Delta \varphi$, the energy after downshift of the various particles can be calculated by taking the difference of the measured energy and the acceleration energy, $Z(\Delta \varphi)$. Then, a set of four $\rho R$ versus $T$ curves can be drawn for the calculated energy losses as illustrated in Figure 6.29, where bands, rather than lines, are shown to indicate the uncertainties in each of the energy measurements. In principle, the four bands should overlap at the true $\rho R$ and $T$ when the correct potential is chosen. As can be seen, for neither of the selected potentials, 0.8 MeV or 0.9 MeV, do all four bands intersect at a common point. The potential could not be less than 0.6 MeV (the upshift of the 14.7 MeV proton) and is highly unlikely to be greater than 1.0 MeV (as judged by measured accelerations from thinner shelled targets under the same laser conditions). Absence of a solution is thus a puzzle.

Having four measurements to measure three unknowns is a good way to check the assumptions in the analysis. If only three spectra were measured, a solution could still have been found for certain cases; it is the fourth that indicates there is a problem. The first step however is to examine other shots under similar conditions.
Figure 6.26: Greater initial shell thickness causes larger downshifts in the D-D proton spectrum. Spectral downshifts of 0.9 MeV and 0.4 MeV are seen on targets with shells made of 2.4 μm glass + 7.5 μm CH and 2.4 μm glass + 5 μm CH respectively [94]. These targets were filled with 20 atm of D-T fuel. The spectrum from a 2.0 μm glass shelled target, shifted up by 0.6 MeV is also shown. The acceleration on the thicker shelled targets is probably much less. Shots 13822 and 13817 were filled with D-T fuel, while shot 13778 contained purely D. Laser energies were close to 27 kJ.

Figure 6.27: The predicted ablator ρR versus T relationships for D-D protons ranged down to 2.6 and 2.1 MeV, whose spectra were shown in Figure 6.26. For reference, the curve corresponding to ranging out of the 3.0 MeV protons is also shown. These calculations assume that the glass (SiO₂) plasma (most of the CH having been ablated off), with an ion density of 2×10²⁴ cm⁻³, is responsible for all the energy losses, and that acceleration effects are negligible. The presence of any acceleration would raise the inferred ρR value.
In summary, the energy loss of a variety of particles has been observed, and greater losses have been seen on targets with thicker shells. In the case of the D-^3^He proton with thick-shelled targets, acceleration effects can be neglected and the insensitivity of the stopping power with temperature allows ablator \( \rho R \) to be determined from a single energy loss measurement. On thin-shelled targets, when examining simultaneous shifts of multiple particles whose energy shifts are a function of target potential and ablator temperature as well, it is difficult to reconcile all the measurements. This could indicate a problem with the stopping power model, or perhaps in the simple picture of coronal accelerations through a fixed potential difference. Either way, using measurements of multiple fusion products is an incisive tool for studying the ranging of particles in hot plasmas.

Figure 6.28: A comparison of upshifts for the 4 different charged fusion products generated by a D-^3^He shot when both ranging and acceleration effects are significant [93]. Unlike in Figure 6.20, where the particles suffered negligible energy losses, the addition of ranging effects means that the upshifts are no longer degenerate with \( Z \). The total shift is dependent upon the ablator \( \rho R \) and temperature, as well as the target potential.
Figure 6.29: The downshift of every particle, and its associated error, defines a band in $\rho R$ and $T$ parameter space. In principle, $\rho R$ and $T$ can be determined by finding the overlap region of different bands from multiple particles. The above plots were generated by using the measured energies of $4.3 \pm 0.05$ MeV D-3He alphas, $15.3 \pm 0.1$ MeV D-3He protons, $1.2 \pm 0.1$ MeV D-D tritons, and $3.4 \pm 0.05$ MeV D-D protons as found on shot 13804 (Figure 6.28), assuming accelerating potentials of 0.8 MeV (a) and 0.9 MeV (b). In theory, overlap of all four bands should occur when the correct target potential is chosen. However, never do all four bands overlap, although three bands do occasionally. This illustrates the power of using 4 particles and over-determining the problem, as now the very stopping calculations themselves can be tested. The density of the glass ablator in these calculations was assumed to be $2 \times 10^{24}$ cm$^{-3}$. This analysis was first performed by C. K. Li.
7 Summary

A new charged-particle diagnostic has been designed, constructed and fielded on the OMEGA laser system. The spectrometer, consisting of a 7.6 kilogauss permanent magnet and an array of CR-39 track detectors, explores the energy range between 0.1 – 30 MeV and allows simultaneous measurements of multiple particle species with yields from $10^8$ – $10^{15}$. An automated track scanning system has been developed which is capable of counting up to $10^6$ particles per shot. Scanned data is mapped to the calculated particle trajectories through the magnet, and energy uncertainties are less than 50 keV at the 3 MeV proton position. The symmetry of all measurements is determined using two nearly identical spectrometers mounted with orthogonal views of the target.

Spectra have been obtained for a wide range of primary fusion products, in particular D-D protons and tritons, D-3He protons and alphas, D-T alphas, as well as neutron-scattered fuel ions, or “knock-ons”. Using these particles, a number of important core and ablator parameters have been measured. In particular, D-D proton and D-T alpha measurements are similar, on the average, to corresponding neutron yields, although flux anisotropies of ~ 20 – 30 % are present. Also, ion temperatures from 5 – 15 keV, determined from two independent techniques (fusion yield ratios and Doppler line widths), show reasonable agreement with the neutron measured temperatures. Line broadening of spectra appears to occur for particles losing as little as 5 % of their mean energy. Using well-resolved spectra of knock-on deuterons and tritons up to their maximum energies has allowed direct measurement of core $\rho R$'s of ~ 7 mg/cm$^2$, while ablator $\rho R$’s up to 40 mg/cm$^2$ have been determined from the spectral downshift of 14.7-MeV D-3He protons. The downshift of multiple species from thin-shelled targets has been used to study the regime of charged-particle slowing down where plasma temperatures are important. Puzzling inconsistencies arise when trying to reconcile the shifts of four different particles.

Copious fluxes of ablator protons are observed on all shots with laser intensities greater than $1\times10^{14}$ W/cm$^2$ (the laser wavelength for these experiments was 0.35 μm). These spectra display a wide range of forms, ranging from the exponentially decaying profiles predicted by self-similar plasma expansion, to intense, discrete lines, that are poorly understood. The endpoint energies are usually ~ 1 MeV at intensities of $1\times10^{15}$ W/cm$^2$ and appear unaffected by application of SSD. Simple calculations indicate that hot electron temperatures of tens of keV are required to produce such accelerations. In addition, the charged fusion products have been seen to acquire accelerations of several hundred keV, their energy upshift is usually a factor of two or more less than the endpoint energy of the ablator protons. Both the acceleration of the fusion products and the maximum energy of the ablator protons are similar when measured from both spectrometers.
As diverse as these results may be, it is important to treat them together as a whole, since many of them are interrelated in some way. The most obvious example is the issue of mean energy shifts. If ablator \( \rho R \)'s for ICF-relevant thick-shelled targets are to be measured using the energy downshifts of D-\(^3\)He protons, it is crucial to know how much, if any, upshift was produced by coronal accelerations. Understanding these accelerations requires studying the spectrum of ablator protons and the upshifts of other fusion products, which can only be detected on thin-shelled targets. A consistent picture thus requires measuring spectrum from not only multiple fusion products, but also ablator ions. To this end, the use of track detectors has been crucial for gathering the necessary range of data over a wide variety of target conditions.

With this greater versatility and flexibility comes more data; and with more data comes more questions. There are still a number of unresolved issues. A list might include the inexplicable fluctuations in yield measurements, the difficulties in reconciling energy losses from multiple particles, the discrepancy between temperatures determined by yield ratios at low D-\(^3\)He yields, and the mysterious appearance and disappearance of lines in the ablator proton spectra with their associated effects on fusion product distributions. A key issue to resolve definitively is how rapidly the accelerating electric fields decay during and after the laser pulse. As is common during first operation of a diagnostic, the number of questions quickly becomes greater than the number of answers. On the other hand, these questions yield interesting avenues for further study. Rather than being the culmination of a project then, it is hoped that this is the start of a new field of research, a new window on inertial confinement fusion.
Appendix A: Two simple methods for determining the hot electron temperature from the accelerated ion spectrum

Two simple analytical models will be described that will allow hot electron temperatures to be inferred from measured fast ion spectra. The first model treats the target as an electrostatic sphere that acquires a net charge due to escaping hot electrons; the second method looks more closely at the ion flow process through the one-dimensional fluid equations.

To see the acceleration process in terms of an electrostatic potential acquired by target, it is instructive to consider the ions as infinitely massive (or simply stationary) and the electrons as infinitely mobile so that the system instantaneously achieves electrostatic equilibrium. Electrons with enough energy will escape the restorative pull of the potential they generate, where “escape” is used to indicate reaching the chamber wall or other grounded surface. All others will be trapped within the potential well surrounding the target. The positive potential that is produced means that ions will be accelerated away from the target, acquiring energies equivalent to the target potential multiplied by the ion charge. Clearly, the greater the electron temperature, the larger the number of electrons out on the tail of the distribution that have the requisite escape energies, and thus the larger the potential acquired by the target. The acceleration energy of an ion is thus a measure of the target potential and therefore can be used to infer the electron temperature.

To derive the relationship of target potential to electron temperature, consider a spherical plasma with electrons at a finite temperature and stationary ions. A radius $r_o$ is defined as the boundary of the potential well, dividing electrons that have “escaped” and electrons that have been “trapped”. If $N_{\text{esc}}$ electrons escape, the potential, $\varphi$, measured between $r_o$ and infinity, is simply

$$\varphi = \frac{eN_{\text{esc}}}{4\pi\varepsilon_o r_o} \quad (A.1)$$

Note that by introducing the concept of $r_o$, which is the scale length of the electrostatic potential, the complications involving the precise electron density profiles surrounding the target have been avoided. This crude technique works because in the end, the target potential will not be very sensitive to the precise value of $r_o$. Now the number of electrons that escape is just the number in the electron distribution that have energies greater than $\varphi$. If the electrons can be assumed to have a 3-dimensional Maxwellian distribution, the fraction that escape is given by
\[ \frac{N_{\text{esc}}}{N_{\text{tot}}} = \int_{0}^{\infty} dv \frac{4\pi v^2}{(2\pi v_e^2)^{3/2}} e^{-v^2/2v_e^2} \quad (A.2) \]

where \( N_{\text{tot}} \) is the total number of electrons in the distribution, \( v \) is the speed of the particles, \( v_e^2 = T_e/m_e \), and \( v_0 \) is the escape velocity given by \( v_0 = \sqrt{2e\varphi/m_e} \). As usual, \( T_e \) and \( m_e \) are the electron temperature and mass respectively. Performing this integral gives

\[ \frac{N_{\text{esc}}}{N_{\text{tot}}} = 2 \left( \frac{e\varphi}{T_e} \right) \left( e^{-\varphi/T_e} - \text{erf} \left( \sqrt{\frac{e\varphi}{T_e}} \right) \right) + 1 \quad (A.3) \]

Now the total number of electrons is simply the total energy of the electrons, \( E_{\text{tot}} \), divided by the average energy of each particle, or \( N_{\text{tot}} = E_{\text{tot}} / (3/2T_e) \). Using this along with Eq. (A.1) in Eq. (A.3) thus provides an implicit expression for the target potential in terms of \( T_e, E_{\text{tot}} \), and \( r_0 \):

\[ \varphi = \frac{e}{4\pi e_0 r_0} \frac{2T_e}{3T_e} \left\{ \frac{2}{\sqrt{\pi}} \sqrt{\frac{e\varphi}{T_e}} e^{-\varphi/T_e} - \text{erf} \left( \sqrt{\frac{e\varphi}{T_e}} \right) + 1 \right\} \quad (A.4) \]

From this equation, it turns out that the potential is most sensitive to \( T_e \) and varies only slowly with \( r_0 \) and \( E_{\text{tot}} \). Assuming \( r_0 = 1 \) mm, the target potential is plotted versus electron temperature in Figure 4.2 for \( E_{\text{tot}} \) values of 3, 30, and 300 J.

Although there is a much larger fraction of energy contained in the thermal electrons at \( \sim 0.5 \text{ keV} \), because the target potential is so much more sensitive to temperature than to total energy, it is the suprathermal electrons that determine this potential\(^9\). With a laser output of 30 kJ then, if \( 10^3 - 10^4 \) of the energy is converted to suprathermal electrons with a temperature of 50 keV, the target would charge up to approximately 600 kV. Therefore, protons emerging from the target should be accelerated by 600 keV. It should be noted though that this potential is not maintained indefinitely. Eventually it decays back to zero as currents return from the ionized mounting stalk or from elsewhere in the corona itself. This calculation should thus be viewed as an estimate of the maximum ion energy that can be achieved, early in the laser pulse.

While this simple model predicts the maximum ion energy, it does not give any indication of the spectrum of ions. Calculating such a spectrum requires looking at the time and space dependent details of the plasma expansion and associated ion acceleration, which entails examining the ion-fluid equations. In order to make this problem tractable the following assumptions will be made:

---

\(^9\) Note that since the potential is an order of magnitude greater than the electron temperature, this analysis relies on the high energy tail of the electron distribution being well-represented by a Maxwellian.

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1) The thermal motion of the ions can be neglected. This cold-ion approximation is valid because the kinetic energy acquired by the ions during acceleration is much greater than their thermal energy.

2) Collisionless flow. In the tenuous atmosphere at the plasma-vacuum boundary, the electron-ion collision time is much longer than the time scale for plasma expansion. Note that the assumptions of a collision-free, cold-ion plasma become more accurate as the rarefaction expansion progresses [83].

3) The electron temperature remains constant. Energy input from the laser allows the electrons to maintain a constant temperature, even though they are giving up energy to the ions. The argument usually made [77] is that fresh, hot electrons are produced continuously by the laser in the bulk plasma. Also, the electron thermal conductivity is sufficiently large and the electron-ion energy relaxation time sufficiently small that isothermal conditions prevail.

4) The electrons are in equilibrium with the electrostatic potential generated by the expanding plasma. This is reasonable since the time scale for expansion, governed by the ion motion, is much greater than the electron plasma period. This assumption provides a vast simplification since the electron fluid equations can be replaced by a Boltzmann density distribution.

5) The plasma expansion is planar. This is justified by assuming that most of the acceleration occurs over a scale length less than the target radius (~ 0.5 mm). While it is unclear how true this is, a solution of the equivalent problem in spherical coordinates [75] indicates that the electron temperature inferred from the ion spectrum would be changed by a factor of order 1.

Based on these assumptions, the ion continuity and momentum equations, along with Poisson's equation and the Boltzmann relation, form a closed set. These are shown below:

\[ \frac{\partial n_i}{\partial t} + \frac{\partial}{\partial x} (n_i v) = 0 \]  \hspace{1cm} (A.5)

\[ \frac{\partial}{\partial t} v + v \frac{\partial v}{\partial x} = -\frac{e}{m} \frac{\partial \varphi}{\partial x} \]  \hspace{1cm} (A.6)

\[ \frac{\partial^2 \varphi}{\partial x^2} = \frac{e}{\varepsilon_0} (n_e - n_i) \]  \hspace{1cm} (A.7)

\[ n_e = n_0 e^{\varphi/T} \]  \hspace{1cm} (A.8)

where \( n_i \) is the ion density, \( n_e \) is the electron density, \( n_0 \) is the density of the unperturbed plasma, \( v \) is the ion velocity, \( \varphi \) is the electrostatic potential, and \( T \) is the electron temperature. The initial conditions are that at \( t = 0, n_i = n_e = n_0 \) for \( x \leq 0 \), and \( n_i = n_e = 0 \) for \( x > 0 \). These equations can be solved numerically [73] but in order to obtain an analytical solution, Poisson's equation will be eliminated by assuming quasi-neutrality. Even though it is the presence of space-charge that initiates the whole process of plasma expansion, for a large portion of the flow, quasi-neutrality holds since the characteristic scale length for density variations is large compared to the Debye length. Quasi-neutrality breaks down at the leading
edge of ions, where the ion density drops off over a Debye length, and space-charge effects are significant. One result of assuming quasi-neutrality though is that a maximum ion energy, determined by the energy of the ion front, can no longer be predicted.

Assuming quasi-neutrality and converting to non-dimensional quantities,

\[
\frac{\partial n}{\partial \tilde{x}} + \frac{\partial}{\partial \tilde{x}}(nu) = 0 \tag{A.9}
\]

\[
\frac{\partial u}{\partial \tilde{x}} + u \frac{\partial u}{\partial \tilde{x}} = - \frac{\partial \psi}{\partial \tilde{x}} \tag{A.10}
\]

\[
n = e^\psi \tag{A.11}
\]

where \( \tilde{t} = \omega_p t, \tilde{x} = x / \lambda_D, n = n_i / n_0, u = v / c_s, \psi = e\phi / T, \) and \( \omega_p \) is the ion plasma frequency, \( \lambda_D \) is the Debye length, and \( c_s \) is the ion sound speed \( \sqrt{T / m} \).

Now Eq. (A.11) can be used to eliminate \( \psi \) from (A.10), reducing the system to two equations and two unknowns:

\[
\frac{\partial n}{\partial \tilde{\tau}} + \frac{\partial}{\partial \tilde{\tau}}(nu) = 0 \tag{A.12}
\]

\[
\frac{\partial u}{\partial \tilde{\tau}} + u \frac{\partial u}{\partial \tilde{\tau}} = - \frac{1}{n} \frac{\partial n}{\partial \tilde{\tau}} \tag{A.13}
\]

which are actually identical to the equations for an isothermal, neutral gas. The pressure of a neutral gas and the electric field of a collisionless plasma reduce to the same term on the right-hand side of Eq. (A.13) under isothermal conditions.

These equations yield self-similar solutions [72]. Such solutions are possible because the equations have no characteristic length (once Poisson's equation is removed through quasi-neutrality), and thus quantities can depend on \( x \) and \( t \) through only the ratio \( x/t \) [98]. Some care should be taken though because, although the equations have self-similar solutions, such solutions are only strictly correct if the boundary conditions are self-similar as well. This is not the case for a plasma initially localized, with zero velocity, to the semi-infinite plane \( x < 0 \) – neither the plasma velocity nor density are functions of \( x/t \) alone. However, a numerical integration [73] of the original set of equations (A.5) – (A.9) indicates that, as \( t \to \infty \), the initial conditions are essentially 'forgotten', and the self-similar solution becomes a good approximation.
To find the self-similar solution, assume that \( n \) and \( u \) vary only with \( \xi \), where \( \xi = \frac{x}{t} \). Then (A.12) – (A.13) become:

\[
(u - \xi) \frac{dn}{d\xi} + n \frac{du}{d\xi} = 0 \tag{A.14}
\]

\[
\frac{dn}{d\xi} + n(u - \xi) \frac{du}{d\xi} = 0 \tag{A.15}
\]

Multiplying Eq. (A.15) by \((u - \xi)\) and subtracting (A.14) gives \( u - \xi = \pm 1 \). Taking the positive solution for the case where plasma is expanding in the positive x-direction and plugging this back into (A.14) gives then:

\[
u = 1 + \xi \tag{A.16}
\]

\[
n = e^{-\xi} \tag{A.17}
\]

This solution is valid only for \(-1 \leq \xi \leq \infty\), since the plasma velocity is always in the positive x-direction.

The final step is to obtain the particle velocity distributions (which determine the energy spectrum of ions measured by a detector) from the fluid density and velocity variables that have been determined above. This is possible since the random ion thermal velocities have been assumed to be negligible compared to the velocities acquired by acceleration (a fact confirmed by experiment). When the ions move out of the laser interaction region (or upon termination of the laser pulse), the acceleration process stops, and the particles travel ballistically to the detectors. This means that the ion velocity spectrum, \( dN/dv \), measured by a detector is simply proportional to \( dn/du \). Reverting back to dimensional quantities, and substituting (A.16) into (A.17) gives

\[
\frac{dN}{dv} = K \exp \left( -\frac{v}{\sqrt{T/m}} \right) \tag{A.18}
\]

where \( K \) is a constant. The slope of the \( \ln(dN/dv) \) versus \( v \) curve thus provides a measurement of \( T \), the hot electron temperature. In the case of spherical expansion instead of planar, the spectrum asymptotes to the same functional form, with \( T \) being replaced by \( T/3 \) [75].

Comparing Eqs. (A.17) and (A.18) it can be seen that the distribution of ion velocities (as would be measured by a detector) is the same as the distribution of ions in space; or, in other words, the higher the energy of the ion, the further ahead it is. This means that, for a constant laser input, the highest energy ions are the ones that are accelerated from the plasma at the beginning of the laser pulse.

The self-similar solution forms a starting point for most analytical investigations of plasma expansion. For example, further studies have examined how this spectrum is modified if the electrons have a two-temperature distribution [74, 83], or during the presence of other species [99, 100, 77]. While discussion
of these studies is beyond the scope of this presentation, it should be remarked that presence of a two-
temperature electron distribution causes [74] formation of a collisionless shock wave, provided that the
ratio of hot to cold electron temperatures is greater than $5 + \sqrt{24}$, where this shock wave manifests itself
as a sharp jump in the ion distribution. Presence of a shock wave indicates that the quasi-neutrality
assumption is no longer applicable however, and Poisson’s equation must be solved in the region of the
discontinuity. Studies of collisionless shocks [101] show that strong oscillations can appear in the ion
profile. This is intriguing in light of the experimental results to be described later in this thesis, where the
measured ion spectrum often shows distinct, regularly spaced lines.

From Eq. (A.18) it can be seen that the self-similar solution makes the unphysical prediction that ions
extend to infinite velocities. As mentioned earlier, a full solution of Eqs. (A.5) – (A.8), which include
space-charge effects, predict instead formation of a sharp ion-front, containing the highest energy ions,
that accelerates indefinitely with time. This means that the ions may be accelerated up to, or close to,
electron velocities before the assumption that the electrons are in thermal equilibrium with respect to the
ions no longer holds and these equations break down. Under the physical conditions of laser-plasma
experiments though, this point is probably not reached before the end of the laser pulse terminates the
accelerations, or the ions move out of the laser heating region.

Since the maximum energy, or endpoint, of the ions is a well measured quantity in the experiments to be
described in this thesis, a few remarks should be made about theoretical studies of the endpoint profile.
Eqs. (A.5) – (A.8) predict that the ion spectrum should rise in a ramp just before forming the endpoint
discontinuity [73] [102]. However, References [80] and [76] show that if instead the shape of the
endpoint is determined by a truncation of the high energy electron Maxwellian tail, the ion spectrum will
not show such a rise as the maximum energy is approached. Both these studies conclude that the
endpoint energy of the ions is determined by the maximum energy of the truncated electron distribution.
Appendix B: CR-39 nuclear track detectors

Since CR-39 plays such an important role in this diagnostic, a rather detailed discussion will be given here that provides some background to the track formation process and gives a model that enables track diameters to be predicted using knowledge of a particle’s stopping power.

A solid state nuclear track detector is a material, usually either crystal (such as mica) or polymer (such as Lexan or CR-39), that permanently records the trajectory of heavily ionizing nuclear particles. Chemical etching of the material gouges out conical pits where the particles penetrated the surface, allowing darkened holes to be seen under an optical microscope. Solid state nuclear track detectors are part of a famous tradition of “image” particle detectors – detectors that allow the image of an individual particle trajectory, or at least part of it, to be formed. Cloud chambers, devised in 1911 by C. T. R. Wilson, were the first of the image detectors and were followed by nuclear emulsions, and then bubble chambers (invented in 1952 by D. A. Glaser). All of these played crucial roles in the discovery of elementary particles and the development of twentieth century physics. Solid state track detectors were first discovered in 1958 and, while their impact on physics cannot begin to match those of the earlier devices, they have found wide application in a number of technical fields such as cosmic ray astrophysics, nuclear science and engineering, geology, archaeology, lunar science and the study of meteorites. Their low sensitivity to ionizing radiation, compared to cloud chambers, nuclear emulsions and bubble chambers, prevents their use in many elementary particle physics studies, but enables them to be used in discriminating heavily ionizing particles from an intense background of lightly ionizing radiation. Track detectors respond only to heavily ionizing particles such as fission fragments, alphas, and, in the case of CR-39, protons of a few MeV. They are insensitive to electrons, x-rays and gamma rays. The distinct advantage of track detectors though lies in their low cost, robustness and portability along with their ability to record particle trajectories permanently.

Different track materials have different responses to ionizing particles and CR-39 plastic is one of the few that are sensitive to protons. The nuclear track-recording properties of CR-39 (where the CR stands for Columbia Resin) were first reported in 1978 [103]. Its composition is given by \( \text{C}_{12}\text{H}_{18}\text{O}_7 \). This transparent plastic that is used in many everyday objects, such as eyeglasses, is particularly well suited as a track-recording solid as it is very homogeneous and isotropic, has high radiation sensitivity, is optically transparent, and is etchable using commonly available sodium hydroxide. The uniquely high sensitivity of CR-39 has made it the track detector of choice for light ion studies as the ionization levels of protons

\[ \text{While track detectors do not respond to individual electrons or photons, high fluxes of these particles have been known to alter the etching properties of some track materials, including CR-39 [A. Hussein, KH. Shnishin, A. A. Abou El Kheir, J. of Mat. Sci. 28, 6026 (1993)].} \]
and deuterons are typically below the response thresholds of other detectors. In this experiment CR-39 manufactured by Track Analysis Systems Limited, Bristol, U.K. – referred to as “Tastrak” – has been used throughout. Sample tracks for 3 MeV protons are shown below.

Figure B.1: 3 MeV proton tracks in CR-39 that has been etched for 5.5 hours at 80 °C, in 6.0 molar NaOH. The size of the image is 0.41 x 0.31 mm.

The passage of a heavily ionizing particle through a solid state track detector leaves a narrow (3 – 10 nm) trail of damage which, in polymers, consists of broken molecular chains and free radicals. The amount of damage is an increasing function of the localized energy deposition, or stopping power. There is still debate over the exact mechanism responsible for this damage, whether it is the primary ionization caused by the incident particle itself, or whether it is mainly the ionization caused by the multitude of electrons, generated during passage of the particle, slowing down in the vicinity of particle trajectory; in polymers though, it is generally thought that both mechanisms play some role [104, 105]. Upon chemical etching, the narrow damage trail is more reactive than the surrounding bulk material, resulting in a greater etch rate along the particle trajectory. The shape and surface diameter of the resulting etched cone is a function of the track etch rate to bulk etch rate ratio, \( V_t/V_b \) – the greater this ratio, the greater the size of the track cone. By relating this ratio to properties, in particular the stopping power, of the incident particle, the diameter of etched tracks can be predicted.

\[ V_t/V_b \]

A track diameter is only relevant if the particle is incident normal to the detector surface; only then is a circular track produced – angled incidence creates elliptical tracks. For the purposes of this spectrometer, particles of interest will always be normally incident and, for this reason, a non-circular track can be identified as noise and subsequently rejected.
For the purposes of this spectrometer, it is important to know what track diameters to expect for certain particles at given energies, as this allows each species to be distinguished from others and from noise events. A two-step approach has been taken to achieve this. Firstly, calibrations have been obtained for particles at discrete energies. Secondly, a model for track formation has been constructed whose free parameters were constrained by the available data.

![Diagram of track etching](image)

Figure B.2: A generalized track profile of diameter $D$, after etching for a time $t$ showing surfaces before and after etching. The bulk etch rate is $V_B$. The coordinates of the track surface are given by $(x_T, y_T)$. In the analysis of track growth, etching to the point $(x_T, y_T)$ is assumed to follow the route $(0,0) \rightarrow (0,y') \rightarrow (x_T, y_T)$, where $(0,y')$ is the point where the trajectory leaves the center line.

The geometry of track etching is governed by the simultaneous action of two processes: the chemical dissolution along the particle track at a rate, $V_T$, and a general attack on the etched surface and on the interior surface of the etched track at a lesser rate $V_B$. $V_T$ varies along the length of the particle trajectory since the ionization rate of the particle changes as it slows down; in contrast, $V_B$ is fixed by the etchant concentration and temperature and thus can be assumed to be constant. Specifying $V_T(y)$, where $y$ is the depth into the material, and $V_B$, allows the detailed track profile, its depth and diameter to be uniquely determined. A generalized track profile is shown above, where $x$ is the distance parallel to the surface, $y$ is the depth, and $(x_T, y_T)$ is a point on the track profile. The etching time required to reach the point $(x_T, y_T)$ is given by
\[ t = \int_0^{y'} \frac{dy'}{V_T(y)} + \left[ (y_T - y')^2 + x_T^2 \right]^{1/2} / V_B \]  

(B.1)

The first term gives the time taken to etch to the point \((0, y')\) along the track axis; the second term gives the time to etch from \((0, y')\) to \((x_T, y_T)\). Clearly this time varies depending on where the intermediate point \(y'\) lies. The position of \(y'\) is fixed by the Huygens, or least-time, construction which says that the path of minimum time represents the actual path taken. This is the same method that can be used to determine, for instance, the path of a light ray traveling between a point in air and a point inside a glass block: the path which gives the minimum travel time between these two points fixes the angle of refraction at the glass surface. Taking the derivative of \(dt/dy'\) and setting it to zero gives an implicit expression for the actual position \(y'\):

\[ y' = y_T - x_T / \left[ (V_T(y') / V_B)^2 - 1 \right]^{1/2} \]  

(B.2)

which, together with Eq. (B.1) defines \(x_T\) and \(y_T\) for a given \(t, y', V_T(y)\) and \(V_B\). A more computationally useful form can be found by noting that the angle \(\zeta\) in Figure B.2 is given by \(\tan(\zeta) = x_T / (y_T - y')\). Then,

\[ x_T = [(y_T - y')^2 + x_T^2]^{1/2} \sin(\zeta) \]  

(B.3)

\[ y_T = y' + [(y_T - y')^2 + x_T^2]^{1/2} \cos(\zeta) \]

which can be re-written

\[ x_T = V_B \left[ t - \int_0^{y'} dy / V_T(y) \right] [1 - V_B^2 / V_T^2 (y')]^{1/2} \]  

(B.4)

\[ y_T = y' + V_B^2 / V_T(y') \left[ t - \int_0^{y'} dy / V_T(y) \right] \]  

(B.5)

After a certain etch time, \(t\), the surface of the material has reached a depth \(V_B t\). At the surface though, \(x_T = D/2\), where \(D\) is the track diameter. Thus, to find the track diameter, \(D(t)\), \(y_T\) is set equal to \(V_B t\) and Eq. (B.5) is solved for \(y'\). Usually this will require a numerical procedure since \(V_T(y)\) will not be a simple function. Once \(y'\) is known, it can be substituted into Eq. (B.4) to determine \(D/2\) directly. By calculating \(x_T\) for greater values of \(y'\), the track profile down to the apex can be mapped out. The above approach has been taken from References [104, 105]

Before track profiles can be calculated for a variety of different particles and energies, a prescription must be made that relates \(V_T\) to the energy loss rate of the incident particle. It is difficult to determine the precise form of \(V_T\) but as a simple case, the etch rate ratio (for a particular etchant concentration and temperature) can be assumed to be purely a function of the electronic stopping power, \((dE/dx)_{elec}\):
\[
\frac{V_T}{V_B} = 1 + k \left( \frac{dE}{dx} \right)_{elec}^n
\]  

(B.6)

where \( k \) and \( n \) are free parameters that are to be fixed by comparing the diameter predictions of the track formation model with calibrated data. The stopping power, \((dE/dx)_{elec}\), is determined from the ion stopping code TRIM [106]. The reason why the electronic stopping power is used, rather than the total stopping power, is that it is only the collisions with electrons that produce the ionization responsible for molecular damage. This distinction is really only significant near the end of a particle’s range, where nuclear collisions dominate the slowing down; otherwise, energy losses to electrons are essentially equal to total energy losses. It should be noted here that there is general acceptance [104] that the track etch rate is not purely a function of the electronic stopping power. This is because close collisions of the incident particle with electrons produces high energy electrons which deposit their energy far away from the narrow damage trail around the trajectory. While these energetic electrons are clearly important in the incident particle energy loss calculations, they play a much smaller role in the generation of the damage trail and thus in the determination of the track etch rate. To correct for this, previous workers have replaced the stopping power with either the “restricted energy loss” or the “primary ionization”. The restricted energy loss calculation stipulates that only delta rays with less than \(-350\) eV can cause the necessary damage along the track. Primary ionization assumes that damage is a function not of the energy of electrons (restricted or otherwise) but rather the total number of electrons ionized. Despite these caveats, the behavior of all these quantities is qualitatively similar and for the restricted range of particles and energies required for this spectrometer, Eq. (B.6) is sufficient.

In general, a higher track etch rate ratio results in greater diameters and deeper tracks. Since large, high contrast tracks facilitate track identification and quantification, it is important to understand how best to increase \( V_T/V_B \) by optimizing the etchant temperature and concentration. The etch rate ratio rises monotonically with temperature, making it desirable to run the baths as hot as possible. In contrast, \( V_T/V_B \) peaks at concentrations of \(~ 6.25\) molar NaOH at temperatures of \( > 60 \) °C [107]. Typical operating temperatures for sodium hydroxide (the most common etchant) range from \( 60 – 95 \) °C while concentrations of \( 6.0 – 6.25\) molar NaOH are usually applied [108, 109, 110].

For all experiments described in this thesis, the NaOH was held at \( 80 \) °C and \( 6.0\) molarity. Higher temperatures are more awkward to maintain since the water bath evaporates more rapidly and temperature excursions are more likely. The etch time is adjusted according to the expected track diameters and number densities. Data produced under these conditions for different particles and a variety of energies has been used to constrain this diameter prediction model, which only has two free parameters. A reasonable fit to the available data was found when: \( k = 0.002, n = 1.9\) in Eq. (B.6) where
\( dE/dx \) is in keV/\( \mu \text{m} \). The bulk etch velocity, \( V_B \), was measured to be 1.85 \( \mu \text{m}/\text{hr} \). The predictions of this model and the available data points are shown in an extensive set of curves given below.

In the case where \( V_T \) does not vary substantially over the etch length, \( V_T \) can be assumed to be constant and an explicit expression for the track diameter found from Eqs. (B.4) and (B.5):

\[
D = 2V_B t \sqrt{\frac{V_T / V_B - 1}{V_T / V_B + 1}}
\]  
(B.7)

Substituting Eq. (B.6) into Eq. (B.7), and assuming that \( k(dE/dx)^n \ll 2 \), gives the direct power law relation between track diameter and stopping power that was assumed in Eq. (5.3):

\[
D \approx 2V_B t^{k^{\frac{1}{2}} \left( \frac{dE}{dx} \right)^{n^{\frac{1}{2}}}}
\]  
(B.8)

Using the values \( k = 0.002 \) and \( n = 1.9 \) (where \( dE/dx \) is in keV/\( \mu \text{m} \)) found from the data, and applying the condition that \( k(dE/dx)^n \ll 2 \), Eq. (B.8) only holds when the energy of protons, deuterons, tritons, and alphas are greater than 0.9 MeV, 1.7 MeV, 2.6 MeV, and 23 MeV respectively. Thus this formula is not applicable to fusion-relevant alphas, or to D-D tritons. Note that, implicit in this calculation is the assumption that the stopping power is constant along the etched trajectory. In the limit where \( k(dE/dx)^n >>> 2 \) the track diameter becomes independent of \( dE/dx \) and \( D \approx 2V_B t \), meaning that track diameters for large stopping power particles all saturate at a maximum. This fact becomes important when trying to identify multiple ion species from the accelerated ablator spectra, where heavier ions often have saturated diameters.
Predicted track diameters for protons, deuterons, tritons and alphas in CR-39 etched in 6.0 molarity NaOH at 80 °C for different times. Predictions should not be expected to be more than 30 % accurate. Data points shown are from calibrations using an accelerator, and also from the spectrometer itself. The symbols are: P = proton, D = deuteron, T = triton, A = alpha. Note the scatter in the measured track diameters. The two free parameters in Eq. (B.6) were obtained by obtaining the best fit to these data points. As a result the following were found: \( k = 0.002, n = 1.9 \), where \( dE/dx \) is in keV/\( \mu \)m. The bulk etch velocity was measured to be 1.85 \( \mu \)m/hr. Track diameters below \( \sim 1 \) \( \mu \)m are unrealistic since they are probably not modeled correctly and would be difficult to observe under an optical microscope anyway.
Appendix C: Energy calibration plots for CPS-1 and CPS-2

The energy ranges subtended by each of the CR-39 mount positions for CPS-1 and CPS-2 (compare Figure 5.6). Every point on each detector corresponds to a particular energy value as determined by trajectory calculations of the kind described in Section 5.1.2. The slightly different values for CPS-1 and CPS-2 are due to both alignment and field differences between the two spectrometers. Note the partial redundancies in energy coverage between adjacent mount positions; this is because a single piece of CR-39 spans the distance between two positions.
C9

Distance (cm)

C10

Distance (cm)

C11

Distance (cm)

C12

Distance (cm)

D1

Distance (cm)

D2

Distance (cm)

Proton Energy (MeV)

-1.0 -0.5 0.0 0.5 1.0

-1.0 -0.5 0.0 0.5 1.0

-1.0 -0.5 0.0 0.5 1.0

-1.0 -0.5 0.0 0.5 1.0

Proton Energy (MeV)

-1.0 -0.5 0.0 0.5 1.0

-1.0 -0.5 0.0 0.5 1.0

CPS-1

CPS-2
Appendix D: Experimental procedure and analysis

The process of data gathering and analysis for this spectrometer can be divided into three categories: exposure, etching, and track analysis. The general procedures described here are similar for all experiments performed using this diagnostic. Differences arise mainly in details such as filter selections, etch times and scanning magnifications.

D.1 Track exposure

Before a shot, the specially cut CR-39 pieces must be labeled, cleaned, and loaded into the fingers in the mounting plate. A filter plate is fixed over each piece. Selection of the correct filter is crucial and is highly dependent upon the expected particle types and their energy range at each of the finger positions. A typical set of filters is shown in Table D.1. The mounting plate is then loaded onto the magnet, the diagnostic chamber sealed, pumped down, and finally opened to the target chamber. Following the shot, the gate valve is closed, the diagnostic chamber vented slowly (to prevent damage of the fragile filters) and the mounting plate removed. The CR-39 typically spends less than 1 hour in vacuum. At this stage, a fresh mounting plate with newly loaded track material can be immediately loaded in preparation for the following shot.

D.2 Track etching

In principle there are three controllable variables in the etching procedure: concentration, temperature, and etch time. As explained in Appendix B, track growth is largest at high temperatures and at concentrations close to ~ 6.25 molar NaOH. For practical reasons, etch bath conditions for these experiments are always set at 6.0 molarity and 80 °C, allowing the track sizes to be dictated purely by the etch time.

When deciding upon the correct etch time there are three different issues to consider: the expected particle types, their energies, and their number density. The size of any individual event after a given etch time is fixed by the type of particle and its energy, as explained in Appendix B. Large tracks (and therefore long etch times) are desirable whenever possible so that the CR-39 surface can be rapidly scanned at low magnification. Also, the diameter of large tracks can be determined more accurately than that of smaller tracks (owing to the finite optical resolution of the camera and microscope system) so that multiple particle species, such as 1 MeV tritons and 3 MeV alphas, can be discriminated more precisely.

---

10 Previous studies have indicated that CR-39 exposed to vacuum for longer than 5 hours lose their sensitivity to ionizing radiation [Prohaska et al., "Vacuum desensitization of CR-39 track recording plastic", Rev. Sci. Instrum. v. 65 (9), p. 3020, (1994)]. However, tests with the Tastrak CR-39 used with this spectrometer show that this track material maintains its sensitivity with up to 8 hours of vacuum exposure. On the other hand, Page Moulding CR-39 seemed to begin to lose its sensitivity after ~5 hours.
On the other hand, track sizes need to be restricted when there are high particle densities in order to reduce the probability of track overlap. This typically happens when analyzing D-T alphas and all fast ions. Also, it is generally undesirable to severely "over-etch" tracks, i.e. etch them well beyond the particle range, since their texture can become shallow and difficult to analyze.

With these considerations in mind, the etch times for different particles have been maintained roughly in the ranges shown in Table D.2. In certain conditions (particularly those of high track densities), the etch will need to be performed in several steps in order to find the optimum window when the tracks are big enough to scan and resolve accurately, yet small enough to avoid significant track overlap.

<table>
<thead>
<tr>
<th>Finger position</th>
<th>Proton Energy (MeV)</th>
<th>Upper filter</th>
<th>Expected particles</th>
<th>Lower filter</th>
<th>Expected particles</th>
</tr>
</thead>
<tbody>
<tr>
<td>A7</td>
<td>0.060 – 0.125</td>
<td>0.5 µm Mylar</td>
<td>Protons</td>
<td>0.5 µm Mylar</td>
<td>Protons</td>
</tr>
<tr>
<td>B3</td>
<td>0.105 – 0.170</td>
<td>0.5 µm Mylar</td>
<td>Protons</td>
<td>0.5 µm Mylar</td>
<td>Protons</td>
</tr>
<tr>
<td>B5</td>
<td>0.19 – 0.30</td>
<td>0.5 µm Mylar</td>
<td>Protons</td>
<td>1.5 µm Mylar</td>
<td>Protons</td>
</tr>
<tr>
<td>B7</td>
<td>0.32 – 0.50</td>
<td>1.5 µm Mylar</td>
<td>Protons + some heavier ions</td>
<td>3.0 µm Al</td>
<td>Protons</td>
</tr>
<tr>
<td>C3</td>
<td>0.46 – 0.65</td>
<td>1.5 µm Mylar</td>
<td>Protons + heavier ions</td>
<td>3.0 µm Al</td>
<td>Protons</td>
</tr>
<tr>
<td>C5</td>
<td>0.70 – 0.95</td>
<td>1.5 µm Mylar</td>
<td>Protons + heavier ions</td>
<td>3.0 µm Al</td>
<td>Protons + deuterons</td>
</tr>
<tr>
<td>C7</td>
<td>1.02 – 1.45</td>
<td>3.0 µm Al</td>
<td>Protons + deuterons</td>
<td>3.0 µm Al + 1.5 µm Myl</td>
<td>Protons + deuterons</td>
</tr>
<tr>
<td>C9</td>
<td>1.53 – 2.35</td>
<td>12.5 µm Al</td>
<td>Protons</td>
<td>3.0 µm Al + 1.5 µm Myl</td>
<td>Protons + deuterons + tritons + $^3$He + $^4$He</td>
</tr>
<tr>
<td>C11</td>
<td>2.5 – 4.1</td>
<td>25 µm Al</td>
<td>DD Protons</td>
<td>6 µm Al</td>
<td>DD protons + DD tritons + DT alphas + $^3$He alphas</td>
</tr>
<tr>
<td>C13</td>
<td>4.5 – 8.3</td>
<td>25 µm Al</td>
<td>D$^3$He protons (&lt; 13 MeV)</td>
<td>6 µm Al</td>
<td>DT alphas + D$^3$He alphas</td>
</tr>
<tr>
<td>D8</td>
<td>9.0 – 17.0</td>
<td>900 µm Al</td>
<td>D$^3$He protons (&gt; 14 MeV)</td>
<td>1000 µm Al</td>
<td>D$^3$He protons (&gt; 14 MeV)</td>
</tr>
<tr>
<td>D10</td>
<td>21 – 70</td>
<td>25 µm Al</td>
<td>Triton knock-ons</td>
<td>350 µm Al</td>
<td>Deuteron knock-ons</td>
</tr>
</tbody>
</table>

Table D.1: Typical filter arrangement. Each filter cover usually contains two filter thicknesses – the upper and the lower. This can allow less penetrating particles to be viewed behind one filter and ranged out behind the other. At higher energies, particularly for the 15-MeV protons, it is necessary to have multiple filters to carefully range each part of the spectrum to detectable energies. On D-T shots, D8 will usually have filters to view D and T knock-ons instead of D-$^3$He protons. On D-$^3$He shots, knock-on filters will be removed entirely.
<table>
<thead>
<tr>
<th>Particle</th>
<th>Energy (MeV) after ranging</th>
<th>Finger Positions</th>
<th>Typical number densities (cm$^{-2}$)</th>
<th>Etch Time (hrs)</th>
</tr>
</thead>
<tbody>
<tr>
<td>Protons</td>
<td>0.1 – 1</td>
<td>A7 – C7</td>
<td>$10^6 - 10^7$</td>
<td>0.25 – 0.5</td>
</tr>
<tr>
<td>DT alphas</td>
<td>1.5 – 3.0</td>
<td>C9, C11, C13</td>
<td>$10^6 - 10^7$</td>
<td>0.5</td>
</tr>
<tr>
<td>DD protons</td>
<td>2.0 – 3.0</td>
<td>C11</td>
<td>$10^3 - 10^4$</td>
<td>3 – 8</td>
</tr>
<tr>
<td>DD tritons</td>
<td>0.2 – 0.8</td>
<td>C9, C11</td>
<td>$10^3 - 10^4$</td>
<td>2 – 5</td>
</tr>
<tr>
<td>D$^3$He alphas</td>
<td>1.5 – 3.0</td>
<td>C11, C13</td>
<td>$10^2 - 10^4$</td>
<td>2 – 5</td>
</tr>
<tr>
<td>D$^3$He protons</td>
<td>0.5 – 4.5</td>
<td>D8</td>
<td>$10^2 - 10^4$</td>
<td>6 – 8</td>
</tr>
<tr>
<td>D knock-ons</td>
<td>4.0 – 7.0</td>
<td>D10</td>
<td>$10^2 - 10^3$</td>
<td>6 – 8</td>
</tr>
<tr>
<td>T knock-ons</td>
<td>7.5 – 10.1</td>
<td>D10</td>
<td>$10^2 - 10^3$</td>
<td>6 – 8</td>
</tr>
</tbody>
</table>

Table D.2: Standard etch times used for track pieces from various finger positions. The listed energies are those found after ranging the initial spectrum of energies through the filters shown in Table D.1. This initial spectrum may be broad enough to be spread over a number of filter positions. The range of number densities is representative of the various yields from different shots.

D.3 Track analysis

Once the CR-39 has been suitably etched, the task is to scan the track surface and record the events in digital form. This is done using a computer-controlled, microscope-camera system which captures images, identifies events, rejects noise, and stores data. After the scan, the data file is read by a program which converts the recorded position of events to the appropriate energies using information from trajectory calculations. Statistical corrections and error analysis are performed at this stage.

D.3.1 Microscope scan

The track scanning hardware consists of a microscope, stage and focus controllers, video camera, capture card, and host computer. A Nikon Optiphot-2 brightfield microscope is used with 5 objective lenses: ×4, ×10, ×20, ×40, ×100 to cover the wide dynamic range of track densities. The motorized X-Y stage and focus is controlled by a LUDL controller which responds to commands from either the host computer or a joystick. The stage is driven by stepper motors with a resolution of 0.1 μm per step. Microscope images are recorded using a Hitachi black and white camera and digitized at 8 bits per pixel using an Imascan Precision capture board to form a 640 × 480 pixel image. Using a ×20 objective, the resolution is 0.64 μm per pixel. The host computer is a WIN95, Pentium machine which interfaces with the entire scanning system.

Virtually all the software required to operate this system was written in-house. The scanning program was written using C++ Borland Builder and provides complete control over stage movements, autofocus, video capture and analysis. Stage and focus operation is accomplished via ASCII commands.
between the host computer RS-232 port and the LUDL controller. In order to ensure successful transmission and reception, communications error checking algorithms have been incorporated. Capture of video requires interfacing with the manufacturer supplied Dynamic Link Libraries (DLL's) which allow display of either a captured image or continuous video, at 30 frames per second.

The procedure to scan a piece of CR-39 from the spectrometer starts with aligning it with the x-y directions of the stage and setting the origin of the coordinate system to the center of the piece. The x-axis is taken to be the direction of energy dispersion. Setting up the correct coordinate system is crucial as the position of events relative to this reference frame is ultimately what will be used to determine the particle energies. Scanning is accomplished by stepping through, frame by frame the entire scan area, focusing on each image, analyzing it for events and saving a diameter histogram associated with each image. Once the focal position is found, the most important step – track analysis – can begin. The goal of the analysis is to reduce the information from each image to a diameter histogram of accepted events. This means that following a scan, the reduced data will consist of a 2-dimensional array of diameter histograms, each of which is identified by a particular x-y position on the CR-39. For instance, if a scan covers an area 200×10 images in size, the resulting data array will have 200×10 histograms, each containing, say, 50 diameter bins. Also recorded at every image position is the x-y position (the x position will be mapped to an energy value using the trajectory calculations) and the number of events rejected, for whatever reason.

Event identification and quantification is performed by identifying which tracks are “dark enough” to be real and by defining what constitutes the track edge. This is performed using two parameters: blackness and threshold. The function of both these parameters is illustrated in Figure D.1, which shows a schematic brightness profile of a typical track as seen under a microscope. Note how the region around the center of the track is usually the darkest while the edges blend smoothly back into the background level. In order to measure the track diameter, the edge has to be defined in some way – by a threshold value. All pixels below threshold belong to the event, and all those above do not. An event is thus a contiguous group of pixels which are all below threshold level (a “blob” analysis algorithm developed by Fredrick Séguin identifies all pixels belonging to a contiguous group). If this were the only discriminating parameter, however, many shallow events (due to track defects and other noise) that just dip below threshold, would be spuriously included. Since real tracks are deep pits, their image always contains a group of very dark pixels, whether they be at the center or around the central region. These dark pixels are well below the threshold used to define the edge of the track. By registering only events which contain at least one pixel below a certain “blackness” level (which is always less than the threshold), a large number of noise events are eliminated. The blackness and threshold values are given as percentages of the difference between the darkest possible pixel and the median background level.
Thus 0 % corresponds to completely black and 100 % corresponds to background brightness. Typical values of blackness and threshold are 50 % and 80 % respectively.

**Figure D.1**: Schematic of the brightness profile through a real track and a noise track. A track is first identified by whether or not it possesses a pixel below the requisite blackness level. Blackness is usually set at around 50 %. This step eliminates many faint noise tracks such as the one on the right. Once a track has been identified, all contiguous pixels below the threshold are counted. The periphery of this clump of pixels defines the track diameter.

Before a particular clump of pixels can be accepted as a real event it must pass the “eccentricity” test. Since all signal tracks for this spectrometer are expected to be circular, or very nearly so, (due to normal incidence) all those events with substantial deviations from circularity are taken to be noise. The level of non-circularity is determined by calculating the Fourier amplitudes of the distribution of pixel positions and taking the ratio of the dipole mode amplitude \((m = 2)\) to the circular mode amplitude \((m = 0)\). While 0 % eccentricity corresponds to a perfectly round event \((m = 2\) mode being zero), it is necessary to set the maximum acceptable eccentricity to about 15 % in order to allow for asymmetries caused by finite pixel effects. It is also possible to examine \(m = 3\) and higher modes. This portion of the code was developed by Frederick Séguin.

If an event has the requisite level of blackness and eccentricity, it will be accepted provided it is not touching the boundary. Tracks touching the boundary (including apparently circular ones) are rejected based on the concern that their full size, and thus correct diameter, is not known. In addition, circular tracks that are overlapping will appear as an enlarged, non-circular track and thereby rejected (in most
cases) based on eccentricity considerations. These losses of “true” signal will be statistically compensated later based on the number of tracks that actually were counted.

One of the drawbacks of this track recognition algorithm is that it relies on operator defined parameters (blackness, threshold and eccentricity) that must be determined at the time of the scan. This can be rather subjective as sometimes tracks which can be signal in one context may be seen as noise in another. Thus it is important for the operator to have experience viewing different track types and to recognize the characteristics of a variety of true signals. Without going to a vastly more sophisticated algorithm, it is difficult to imagine how to eliminate the necessity for subjective operator input (see reference [110] for an example of a particularly sophisticated algorithm).

The final step is to measure the diameter of each of the accepted events and generate a histogram of all these events. The diameter is calculated from the total area enclosed within the outer boundary of the pixel cluster. Assuming that this area is the area of a circle, $\pi D^2/4$, the diameter is readily apparent.

**D.3.2 Converting scanned data to energy spectra**

The resulting two-dimensional array of diameter histograms must now be reduced to a particle energy spectrum. This requires mapping of every x-position to an appropriate energy value, summing over the y-direction and selecting a particular diameter range from which to draw the events. Statistical compensation for lost events, outlier elimination, and statistical error calculations are also performed.

Energy conversion utilizes the particle trajectory code described in Section 5.1.2 to calculate the paths of particles over a wide range of energies, determining which of them intersect each piece of CR-39. For an arbitrary position on any CR-39 piece, the energy is determined by interpolating between available particle trajectories. Although a narrow range of energies can actually land at any given point (due to the finite width of the slit), only the energy of the central beam is used. The width of the beam is recorded for use with future deconvolution algorithms.

Clearly, correct determination of the particle energies depends on accurate knowledge of the alignment of the spectrometer as well as confidence that the field structure calculated and measured by the magnet manufacturer is unperturbed by the experimental environment. Use of x-ray film behind the magnet allows any misalignments to be detected and then input into the trajectory calculations. The ultimate test of the energy scale is an absolute, *in situ* calibration using differential ranging filters. This was done for both CPS-1 and CPS-2.
Statistical compensation is performed on each histogram individually. This process takes a histogram and calculates, for each diameter bin, the most likely number of events that were thrown out due to overlapping and boundary clipping. These events are added back to create the "true" histogram. As expected, the correction is largest when there are high track densities and is greater for large events than for small. This algorithm was developed by Jeremy Schnittman.

Any column of histograms can be considered to be a statistical ensemble of measurement samples (energy dispersion only occurring in the x-direction – along a row) and thus it is desirable to scan as many rows as is practical in order to build up counts and reduce statistical errors. Also, a large number of samples allows outliers to be eliminated. Outliers arise when a scratch or dirt patch obscures a particular image and therefore distorts its diameter histogram. Assuming that this obstruction is only present on a minority of images in the scan column, the erroneous histogram can be eliminated by recognizing its deviation from others in the same column.

The next step is to select a particular diameter window and to sum together events within this range from all the histograms (minus the outliers) in a column. Selecting the diameter window is a crucial step in the analysis and represents the crux of this spectrometer. In one step, a large fraction of the noise events (which usually have a continuum of diameters) and, in principle, all other particle types, are removed from the counting process. The presence of the magnet allows this diameter window to be made as narrow as possible. An example of how this is done is illustrated in Figure D.2. Finally, the sum of all these valid events in a given energy bin is converted to units of yield per MeV, by appropriate scaling with the solid angle subtended by the collimator. The error bar at each energy bin is calculated from the square root of the total number of events counted before compensation, and then scaled appropriately to account for compensation. The statistical errors introduced by the compensation process are added in quadrature.
Figure D.2: A contour plot of how diameter histograms of signal tracks can vary over a piece of CR-39 positioned behind the magnet. This sample was obtained from the D8 position observing knock-on tritons in the range 2.3 – 5.0 MeV and deuterons between 3.5 – 7.5 MeV, where the energies are those before ranging through the 25 μm aluminum filter. At each energy position (corresponding to an x-position) the tritons and deuterons are confined to their own, distinct diameter ranges. These diameter ranges vary with energy and trace out the envelopes shown above. It is crucial to select the correct diameter range in order to select the desired particle and eliminate as much of the background noise as possible. The scattered events that do not appear to be confined to particular diameter bands are characteristic of background noise. The band appearing at 3 – 4 μm at x ~ -1.0 cm are small diameter noise events which appear to have similar diameters because of the finite pixel resolution of the camera.

D.4 Laser and target conditions

During initial development and shake-down of this diagnostic, a variety of different target and laser conditions have been utilized, summarized in Table D.3. Data has been acquired from a combination of both dedicated and “piggy-back” shots. The wide variety of shot conditions has been valuable in defining the range of utility of the instruments, as well as in exploring various physics issues.
<table>
<thead>
<tr>
<th>Laser conditions</th>
</tr>
</thead>
<tbody>
<tr>
<td><strong>Wavelength</strong></td>
</tr>
<tr>
<td><strong>Energy</strong></td>
</tr>
<tr>
<td><strong>Number of beams</strong></td>
</tr>
<tr>
<td><strong>Pulse length</strong></td>
</tr>
<tr>
<td><strong>Pulse shape</strong></td>
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<tr>
<td><strong>Laser intensity</strong></td>
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<tr>
<td><strong>Beam smoothing</strong></td>
</tr>
</tbody>
</table>

<table>
<thead>
<tr>
<th>Target conditions</th>
</tr>
</thead>
<tbody>
<tr>
<td><strong>Shape</strong></td>
</tr>
<tr>
<td><strong>Diameter</strong></td>
</tr>
<tr>
<td><strong>Fuel</strong></td>
</tr>
<tr>
<td><strong>Pressure</strong></td>
</tr>
<tr>
<td><strong>Shell Material</strong></td>
</tr>
<tr>
<td><strong>Shell thickness</strong></td>
</tr>
<tr>
<td><strong>Fusion yields</strong></td>
</tr>
<tr>
<td></td>
</tr>
<tr>
<td></td>
</tr>
</tbody>
</table>

Table D.3: Summary of various laser and target conditions for the experiments performed so far on OMEGA.
Appendix E: Compilation of spectra

An assortment of spectra from various different types of shots has been compiled here for reference purposes. The focus is on shots not shown in the main body of this thesis, though some repetition exists. Listed uncertainties include only statistical errors, not systematic errors.

<table>
<thead>
<tr>
<th>Shot #</th>
<th>Target O.D (µm)</th>
<th>Shell Thickness (µm)</th>
<th>Pressure (atm)/Fill</th>
<th>Energy (kJ)</th>
<th>Pulse Width (ns)</th>
<th>SSD</th>
<th>Intensity (W/cm²)</th>
</tr>
</thead>
<tbody>
<tr>
<td>11552</td>
<td>931</td>
<td>2.3G</td>
<td>5D</td>
<td>27.05</td>
<td>1.1</td>
<td>On</td>
<td>9.03E+14</td>
</tr>
</tbody>
</table>

CPS-1; DD protons
YIELD (Gaussian) = 1.87E+011 ± 2.47E+009
\(<E>\) (Gaussian) = 3.623 ± 0.002 MeV
\(\sigma\) (Gaussian) = 0.132 ± 0.002 MeV

<table>
<thead>
<tr>
<th>Shot #</th>
<th>Target O.D (µm)</th>
<th>Shell Thickness (µm)</th>
<th>Pressure (atm)/Fill</th>
<th>Energy (kJ)</th>
<th>Pulse Width (ns)</th>
<th>SSD</th>
<th>Intensity (W/cm²)</th>
</tr>
</thead>
<tbody>
<tr>
<td>11555</td>
<td>931</td>
<td>2.2G</td>
<td>20D</td>
<td>26.64</td>
<td>1.1</td>
<td>On</td>
<td>8.89E+14</td>
</tr>
</tbody>
</table>

CPS-1; DD protons
YIELD (Gaussian) = 3.27E+011 ± 7.87E+009
\(<E>\) (Gaussian) = 3.430 ± 0.005 MeV
\(\sigma\) (Gaussian) = 0.125 ± 0.005 MeV
### Table 1: Shot Parameters

<table>
<thead>
<tr>
<th>Shot #</th>
<th>Target O.D (µm)</th>
<th>Shell Thickness (µm)</th>
<th>Pressure (atm)/Fill</th>
<th>Energy (kJ)</th>
<th>Pulse Width (ns)</th>
<th>SSD</th>
<th>Intensity (W/cm²)</th>
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</thead>
<tbody>
<tr>
<td>11585</td>
<td>882</td>
<td>2.5G</td>
<td>5D/5He3</td>
<td>20.91</td>
<td>1.1</td>
<td>On</td>
<td>7.78E+14</td>
</tr>
</tbody>
</table>

CPS-1; DD protons

YIELD (Gaussian) = 4.11E+010 ± 2.48E+009

\[<E>\) (Gaussian) = 3.245 ± 0.004 MeV

\[\sigma\) (Gaussian) = 0.138 ± 0.005 MeV

CPS-1; D³He protons

YIELD (Gaussian) = 4.67E+010 ± 1.93E+009

\[<E>\) (Gaussian) = 14.903 ± 0.012 MeV

\[\sigma\) (Gaussian) = 0.325 ± 0.011 MeV

CPS-1; DT alphas

YIELD = 1.01E+013 ± 9.42E+010

\[<E>\) = 3.667 ± 0.005 MeV

\[\sigma\) = 0.393 ± 0.002 MeV
CPS-1; DD protons
YIELD (Gaussian) = 2.47E+010 ± 8.90E+008
<E> (Gaussian) = 3.280 ± 0.005 MeV
σ (Gaussian) = 0.151 ± 0.005 MeV

CPS-1; Ablator Protons

<table>
<thead>
<tr>
<th>Shot #</th>
<th>Target O.D (μm)</th>
<th>Shell Thickness (μm)</th>
<th>Pressure (atm)/Fill</th>
<th>Energy (kJ)</th>
<th>Pulse Width (ns)</th>
<th>SSD</th>
<th>Intensity (W/cm²)</th>
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</thead>
<tbody>
<tr>
<td>11742</td>
<td>918</td>
<td>2.6G</td>
<td>5D/5T</td>
<td>21.99</td>
<td>1.11</td>
<td>On</td>
<td>7.48E+14</td>
</tr>
</tbody>
</table>

CPS-1; DT alphas
YIELD (Gaussian) = 4.30E+013 ± 1.82E+011
<E> (Gaussian) = 3.049 ± 0.001 MeV
σ (Gaussian) = 0.367 ± 0.001 MeV
CPS-1; DD protons
YIELD (Gaussian) = 1.06E+011 ± 3.56E+009
$\langle E \rangle$ (Gaussian) = 3.003 ± 0.004 MeV
$\sigma$ (Gaussian) = 0.136 ± 0.004 MeV

CPS-1;
Knock-on deuterons

CPS-1;
Knock-on tritons
<table>
<thead>
<tr>
<th>Shot #</th>
<th>Target O.D (μm)</th>
<th>Shell Thickness (μm)</th>
<th>Pressure (atm)/Fill</th>
<th>Energy (kJ)</th>
<th>Pulse Width (ns)</th>
<th>SSD</th>
<th>Intensity (W/cm²)</th>
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</thead>
<tbody>
<tr>
<td>11743</td>
<td>910</td>
<td>2.3G</td>
<td>5D/5T</td>
<td>23.03</td>
<td>1.11</td>
<td>On</td>
<td>7.98E+14</td>
</tr>
</tbody>
</table>

**CPS-1; DT alphas**

YIELD = 3.11E+013 ± 1.91E+011

\(<E> = 3.601 ± 0.004 \text{ MeV} \\
\sigma = 0.410 ± 0.002 \text{ MeV}

**CPS-1; DD protons**

YIELD (Gaussian) = 9.47E+010 ± 2.24E+009

\(<E> (\text{Gaussian}) = 3.246 ± 0.003 \text{ MeV} \\
\sigma (\text{Gaussian}) = 0.148 ± 0.003 \text{ MeV}

**CPS-1: Ablator protons**
<table>
<thead>
<tr>
<th>Shot #</th>
<th>Target O.D (µm)</th>
<th>Shell Thickness (µm)</th>
<th>Pressure (atm)/Fill</th>
<th>Energy (kJ)</th>
<th>Pulse Width (ns)</th>
<th>SSD</th>
<th>Intensity (W/cm²)</th>
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</thead>
<tbody>
<tr>
<td>11744</td>
<td>944</td>
<td>2G/4CH</td>
<td>1.7D/1.7T</td>
<td>9.45</td>
<td>1.14</td>
<td>On</td>
<td>2.96E+14</td>
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</tbody>
</table>

CPS-1; DT alphas
YIELD (Gaussian) = 1.89E+012 ± 9.33E+009

<> (Gaussian) = 3.675 ± 0.001 MeV

σ (Gaussian) = 0.298 ± 0.001 MeV

1.0 cm collimator slit

CPS-1; DD protons
YIELD (Gaussian) = 5.02E+009 ± 1.18E+008

<> (Gaussian) = 3.214 ± 0.004 MeV

σ (Gaussian) = 0.202 ± 0.004 MeV

1.0 cm collimator slit

CPS-1: Ablator Protons
1.0 cm collimator slit
<table>
<thead>
<tr>
<th>Shot #</th>
<th>Target O.D (μm)</th>
<th>Shell Thickness (μm)</th>
<th>Pressure (atm)/Fill</th>
<th>Energy (kJ)</th>
<th>Pulse Width (ns)</th>
<th>SSD</th>
<th>Intensity (W/cm²)</th>
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</thead>
<tbody>
<tr>
<td>11757</td>
<td>938</td>
<td>2G</td>
<td>5D/5T</td>
<td>20.62</td>
<td>1.1</td>
<td>On</td>
<td>6.78E+14</td>
</tr>
</tbody>
</table>

CPS-1; DT alphas
YIELD (Gaussian) = 1.65E+013 ± 7.59E+010
⟨E⟩ (Gaussian) = 3.093 ± 0.002 MeV
σ (Gaussian) = 0.426 ± 0.001 MeV

CPS-1; DD protons
YIELD (Gaussian) = 4.52E+010 ± 1.00E+009
⟨E⟩ (Gaussian) = 3.028 ± 0.003 MeV
σ (Gaussian) = 0.159 ± 0.003 MeV

CPS-1: Ablator protons
Shot # | Target O.D (μm) | Shell Thickness (μm) | Pressure (atm)/Fill | Energy (kJ) | Pulse Width (ns) | SSD | Intensity (W/cm²)
--- | --- | --- | --- | --- | --- | --- | ---
11763 | 942 | 2.4G | 5D/0.5He3 | 20.95 | 1.09 | On | 6.89E+14

CPS-1; DD protons
- Gaussian fit from 2.82 to 3.59 MeV
- \(YIELD\) (Gaussian) = 2.92E+11 ± 5.62E+09
- \(<E>\) (Gaussian) = 3.180 ± 0.001 MeV
- \(\sigma\) (Gaussian) = 0.144 ± 0.001 MeV

CPS-1; DD tritons
- \(YIELD\) (Gaussian) = 2.27E+11 ± 3.62E+09
- \(<E>\) (Gaussian) = 1.074 ± 0.002 MeV
- \(\sigma\) (Gaussian) = 0.147 ± 0.002 MeV

CPS-1; \(^3\)He protons
- \(YIELD\) (Gaussian) = 6.61E+09 ± 2.27E+08
- \(<E>\) (Gaussian) =14.844 ± 0.010 MeV
- \(\sigma\) (Gaussian) = 0.336 ± 0.010 MeV
<table>
<thead>
<tr>
<th>Shot #</th>
<th>Target O.D (µm)</th>
<th>Shell Thickness (µm)</th>
<th>Pressure (atm)/Fill</th>
<th>Energy (kJ)</th>
<th>Pulse Width (ns)</th>
<th>SSD</th>
<th>Intensity (W/cm²)</th>
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<tbody>
<tr>
<td>12141</td>
<td>900</td>
<td>20CH</td>
<td>10D</td>
<td>21.98</td>
<td>1</td>
<td>Off</td>
<td>8.64E+14</td>
</tr>
</tbody>
</table>

**CPS-1: Ablator Protons**

![Shot 12141: CPS-1](image)

**Shot 12149**

<table>
<thead>
<tr>
<th>Shot #</th>
<th>Target O.D (µm)</th>
<th>Shell Thickness (µm)</th>
<th>Pressure (atm)/Fill</th>
<th>Energy (kJ)</th>
<th>Pulse Width (ns)</th>
<th>SSD</th>
<th>Intensity (W/cm²)</th>
</tr>
</thead>
<tbody>
<tr>
<td>12149</td>
<td>900</td>
<td>2.5G</td>
<td>10D</td>
<td>21.51</td>
<td>1</td>
<td>Off</td>
<td>8.45E+14</td>
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</tbody>
</table>

**CPS-1; DD protons**

YIELD (Gaussian) = 1.90E+01 ± 3.06E+009  
<Φ> (Gaussian) = 3.283 ± 0.002 MeV  
σ (Gaussian) = 0.133 ± 0.002 MeV

**CPS-1; DD tritons**

YIELD (Gaussian) = 2.00E+01 ± 3.46E+009  
<Φ> (Gaussian) = 1.131 ± 0.002 MeV  
σ (Gaussian) = 0.138 ± 0.002 MeV

![Shot 12149](image)
### Shot 12149: CPS-1

CPS-1: Ablator Protons

![Graph](image1)

<table>
<thead>
<tr>
<th>Shot #</th>
<th>Target</th>
<th>Shell Thickness</th>
<th>Pressure (atm)/Fill</th>
<th>Energy (kJ)</th>
<th>Pulse Width (ns)</th>
<th>SSD</th>
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</thead>
<tbody>
<tr>
<td>12152</td>
<td>900</td>
<td>20CH</td>
<td>8.5D/8.5T</td>
<td>21.13</td>
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<td>8.3E+14</td>
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</tbody>
</table>

### Shot 12152: CPS-1

CPS-1: Ablator Protons

![Graph](image2)

<table>
<thead>
<tr>
<th>Shot #</th>
<th>Target</th>
<th>Shell Thickness</th>
<th>Pressure (atm)/Fill</th>
<th>Energy (kJ)</th>
<th>Pulse Width (ns)</th>
<th>SSD</th>
<th>Intensity (W/cm²)</th>
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</thead>
<tbody>
<tr>
<td>12154</td>
<td>900</td>
<td>3.0G</td>
<td>20D</td>
<td>19.35</td>
<td>1</td>
<td>Off</td>
<td>7.6E+14</td>
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</tbody>
</table>

### Shot 12154

CPS-1: Ablator Protons

![Graph](image3)
<table>
<thead>
<tr>
<th>Shot #</th>
<th>Target O.D (μm)</th>
<th>Shell Thickness (μm)</th>
<th>Pressure (atm)</th>
<th>Fill</th>
<th>Energy (kJ)</th>
<th>Pulse Width (ns)</th>
<th>SSD</th>
<th>Intensity (W/cm²)</th>
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</thead>
<tbody>
<tr>
<td>12941</td>
<td>918</td>
<td>2.6G</td>
<td>5D/5He3</td>
<td></td>
<td>29.6</td>
<td>1</td>
<td>Off</td>
<td>1.12E+15</td>
</tr>
</tbody>
</table>

CPS-1: Ablator Protons

Shot 12941

![Graph](image1.png)

Shot 13402

<table>
<thead>
<tr>
<th>Shot #</th>
<th>Target O.D (μm)</th>
<th>Shell Thickness (μm)</th>
<th>Pressure (atm)</th>
<th>Fill</th>
<th>Energy (kJ)</th>
<th>Pulse Width (ns)</th>
<th>SSD</th>
<th>Intensity (W/cm²)</th>
</tr>
</thead>
<tbody>
<tr>
<td>13402</td>
<td>903</td>
<td>3.3G</td>
<td>3D/6He3</td>
<td></td>
<td>26.9</td>
<td>1</td>
<td>On</td>
<td>1.05E+15</td>
</tr>
</tbody>
</table>

DD triton (CPS-2)

Yield (Gaussian) = 6.58E+009 ± 3.82E+009
<E> (Gaussian) = 1.067 ± 0.034 MeV
Sigma (Gaussian) = 0.050 ± 0.035 MeV

D³He alpha (CPS-2)

Yield (Gaussian) = 9.31E+009 ± 4.31E+009
<E> (Gaussian) = 3.903 ± 0.343 MeV
Sigma (Gaussian) = 0.288 ± 0.322 MeV
<table>
<thead>
<tr>
<th>Shot #</th>
<th>Target O.D (µm)</th>
<th>Shell Thickness (µm)</th>
<th>Pressure (atm)/Fill</th>
<th>Energy (kJ)</th>
<th>Pulse Width (ns)</th>
<th>SSD</th>
<th>Intensity (W/cm²)</th>
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</thead>
<tbody>
<tr>
<td>13407</td>
<td>938</td>
<td>2.3G</td>
<td>10D</td>
<td>25.3</td>
<td>1</td>
<td>On</td>
<td>9.15E+14</td>
</tr>
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<tr>
<td>Shot #</td>
<td>Target O.D (µm)</td>
<td>Shell Thickness (µm)</td>
<td>Pressure (atm)/Fill</td>
<td>Energy (kJ)</td>
<td>Pulse Width (ns)</td>
<td>SSD</td>
<td>Intensity (W/cm²)</td>
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<tr>
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<tr>
<td>13410</td>
<td>915</td>
<td>3.2CH/2.3CHCl/9.3CH</td>
<td>3D/6He3</td>
<td>26.6</td>
<td>1</td>
<td>On</td>
<td>1.01E+15</td>
</tr>
<tr>
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</tr>
</tbody>
</table>

CPS-1; DD proton
YIELD (Gaussian) = 8.73E+009 ± 6.35E+008
<\text{E}> (Gaussian) = 3.524 ± 0.010 MeV
\sigma (Gaussian) = 0.166 ± 0.010 MeV

CPS-2; DD proton
YIELD (Gaussian) = 3.81E+009 ± 2.95E+008
<\text{E}> (Gaussian) = 3.555 ± 0.010 MeV
\sigma (Gaussian) = 0.213 ± 0.009 MeV

CPS-2; D³He proton
YIELD (Gaussian) = 9.12E+008 ± 8.83E+007
<\text{E}> (Gaussian) = 14.119 ± 0.035 MeV
\sigma (Gaussian) = 0.407 ± 0.037 MeV
<table>
<thead>
<tr>
<th>Shot #</th>
<th>Target O.D (μm)</th>
<th>Shell Thickness (μm)</th>
<th>Pressure (atm)/Fill</th>
<th>Energy (kJ)</th>
<th>Pulse Width (ns)</th>
<th>SSD</th>
<th>Intensity (W/cm²)</th>
</tr>
</thead>
<tbody>
<tr>
<td>13411</td>
<td>919</td>
<td>3.2CH/2.3CHCl/9.4CH</td>
<td>3D/6He3</td>
<td>25.5</td>
<td>1</td>
<td>On</td>
<td>9.61E+14</td>
</tr>
<tr>
<td></td>
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</tr>
<tr>
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<td>958</td>
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<td>20D</td>
<td>26.8</td>
<td>1</td>
<td>On</td>
<td>9.3E+14</td>
</tr>
</tbody>
</table>

CPS-2; D³He proton
YIELD (Gaussian) = 8.08E+008 ± 8.50E+007

\[ \langle E \rangle \text{ (Gaussian)} = 14.08 \pm 0.039 \text{ MeV} \]
\[ \sigma \text{ (Gaussian)} = 0.420 \pm 0.037 \text{ MeV} \]

CPS-1; DD Triton
YIELD (Gaussian) = 1.43E+011 ± 3.76E+009

\[ \langle E \rangle \text{ (Gaussian)} = 1.545 \pm 0.003 \text{ MeV} \]
\[ \sigma \text{ (Gaussian)} = 0.119 \pm 0.002 \text{ MeV} \]

CPS-1; DD Proton
YIELD (Gaussian) = 1.48E+011 ± 3.72E+009

\[ \langle E \rangle \text{ (Gaussian)} = 3.588 \pm 0.003 \text{ MeV} \]
\[ \sigma \text{ (Gaussian)} = 0.124 \pm 0.003 \text{ MeV} \]
Shot 13778

CPS-2; DD Triton
Gaussian fit from 1.22 to 1.58 MeV
YIELD (Gaussian) = 1.82E+011 ± 1.93E+009
<E> (Gaussian) = 1.460 ± 0.002 MeV
σ (Gaussian) = 0.123 ± 0.002 MeV

CPS-2; DD Proton
YIELD (Gaussian) = 1.80E+011 ± 1.93E+009
<E> (Gaussian) = 3.648 ± 0.001 MeV
σ (Gaussian) = 0.112 ± 0.001 MeV

Shot # | Target O.D (µm) | Shell Thickness (µm) | Pressure (atm)/Fill | Energy (kJ) | Pulse Width (ns) | SSD | Intensity (W/cm²) |
-- | -- | -- | -- | -- | -- | -- | -- |
13795 | 908.2 | 14.6CH | 5D/10He3 | 28.4 | 1 | On | 1.1E+15 |

CPS-2; D³He Proton
YIELD (Gaussian) = 3.00E+009 ± 2.40E+008
<E> (Gaussian) = 13.806 ± 0.026 MeV
σ (Gaussian) = 0.371 ± 0.025 MeV
<table>
<thead>
<tr>
<th>Shot #</th>
<th>Target O.D (µm)</th>
<th>Shell Thickness (µm)</th>
<th>Pressure (atm)/Fill</th>
<th>Energy (kJ)</th>
<th>Pulse Width (ns)</th>
<th>SSD</th>
<th>Intensity (W/cm²)</th>
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</thead>
<tbody>
<tr>
<td>13796</td>
<td>907</td>
<td>14.5CH</td>
<td>5D/10He³</td>
<td>28.4</td>
<td>1</td>
<td>On</td>
<td>1.1E+15</td>
</tr>
</tbody>
</table>

CPS-2; D³He proton
YIELD (Gaussian) = 3.69E+009 ± 2.40E+008
<\text{E}> (Gaussian) = 14.092 ± 0.022 MeV
\sigma (Gaussian) = 0.363 ± 0.017 MeV

<table>
<thead>
<tr>
<th>Shot #</th>
<th>Target O.D (µm)</th>
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<th>Pressure (atm)/Fill</th>
<th>Energy (kJ)</th>
<th>Pulse Width (ns)</th>
<th>SSD</th>
<th>Intensity (W/cm²)</th>
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<tbody>
<tr>
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<td>18.4CH</td>
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<td>1.02E+15</td>
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</table>

CPS-2; D³He Proton
YIELD (Gaussian) = 1.23E+009 ± 1.06E+008
<\text{E}> (Gaussian) = 13.052 ± 0.042 MeV
\sigma (Gaussian) = 0.538 ± 0.037 MeV
<table>
<thead>
<tr>
<th>Shot #</th>
<th>Target</th>
<th>Shell Thickness</th>
<th>Pressure</th>
<th>Energy</th>
<th>Pulse Width</th>
<th>SSD</th>
<th>Intensity</th>
</tr>
</thead>
<tbody>
<tr>
<td>13823</td>
<td>898</td>
<td>2.2G</td>
<td>5D/10He3</td>
<td>26.6</td>
<td>1</td>
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</tbody>
</table>

**CPS-2; DD triton**

YIELD (Gaussian) = 3.19E+010 ± 8.27E+008  
\(<E>\) (Gaussian) = 1.491 ± 0.007 MeV  
\(\sigma\) (Gaussian) = 0.135 ± 0.006 MeV

**CPS-1; DD Proton**

YIELD (Gaussian) = 2.21E+010 ± 1.46E+009  
\(<E>\) (Gaussian) = 3.486 ± 0.007 MeV  
\(\sigma\) (Gaussian) = 0.123 ± 0.007 MeV

**CPS-2; DD Proton**

YIELD (Gaussian) = 2.45E+010 ± 6.51E+008  
\(<E>\) (Gaussian) = 3.537 ± 0.003 MeV  
\(\sigma\) (Gaussian) = 0.131 ± 0.003 MeV
CPS-1; D³He Proton
YIELD (Gaussian) = 4.16E+010 ± 2.22E+009
<E> (Gaussian) = 15.127 ± 0.015 MeV
σ (Gaussian) = 0.330 ± 0.014 MeV

CPS-2; D³He Proton
YIELD (Gaussian) = 4.31E+010 ± 1.12E+009
<E> (Gaussian) = 15.435 ± 0.008 MeV
σ (Gaussian) = 0.378 ± 0.008 MeV

<table>
<thead>
<tr>
<th>Shot #</th>
<th>Target O.D (µm)</th>
<th>Shell Thickness (µm)</th>
<th>Pressure (atm)/Fill</th>
<th>Energy (kJ)</th>
<th>Pulse Width (ns)</th>
<th>SSD</th>
<th>Intensity (W/cm²)</th>
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<td>13823</td>
<td>921</td>
<td>3G</td>
<td>5D/10He3</td>
<td>26.8</td>
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</table>

CPS-1; D³He proton
YIELD (Gaussian) = 4.96E+010 ± 2.19E+009
<E> (Gaussian) = 15.045 ± 0.013 MeV
σ (Gaussian) = 0.332 ± 0.013 MeV
CPS-2; D³He Proton
YIELD (Gaussian) = 4.92E+10 ± 1.03E+009
<E> (Gaussian) = 15.179 ± 0.006 MeV
σ (Gaussian) = 0.347 ± 0.006 MeV

CPS-1; DD Proton
YIELD (Gaussian) = 5.20E+10 ± 2.11E+009
<E> (Gaussian) = 3.197 ± 0.005 MeV
σ (Gaussian) = 0.138 ± 0.005 MeV

CPS-2; DD Proton
YIELD (Gaussian) = 3.50E+10 ± 7.76E+008
<E> (Gaussian) = 3.291 ± 0.003 MeV
σ (Gaussian) = 0.136 ± 0.003 MeV
CPS-2; D^3He Alpha
YIELD (Gaussian) = 3.93E+010 ± 8.32E+008
<E> (Gaussian) = 3.880 ± 0.006 MeV
σ (Gaussian) = 0.330 ± 0.006 MeV

CPS-2; DD triton
YIELD (Gaussian) = 5.52E+010 ± 1.14E+009
<E> (Gaussian) = 1.040 ± 0.006 MeV
σ (Gaussian) = 0.129 ± 0.005 MeV
Bibliography

[53] R. D. Petrasso, private communication
[61] LLRE Review, 73 Dec 1997