Studies of High-Energy Photon Sources from a Laser Wakefield Accelerator

by

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<tr>
<td>CCD</td>
<td>Charge-Coupled Device</td>
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<td>Central Laser Facility</td>
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<td>CPA</td>
<td>Chirped Pulse Amplification</td>
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<td>CUOS</td>
<td>Center for Ultrafast Optical Science</td>
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<td>LPA</td>
<td>Laser-Plasma Accelerator</td>
<td></td>
</tr>
<tr>
<td>LWFA</td>
<td>Laser WakeField Acceleration</td>
<td></td>
</tr>
<tr>
<td>MPI</td>
<td>Multi-Photon Ionization</td>
<td></td>
</tr>
</tbody>
</table>
**OPL**  Optical Path Length

**PET**  PEnTaerythritol

**PWFA**  Plasma WakeField Acceleration

**QED**  Quantum ElectroDynamics

**RAL**  Rutherford Appleton Laboratory

**SHG**  Second Harmonic Generation

**SLAC**  Stanford Linear Accelerator

**TCC**  Target Chamber Center
ABSTRACT

Since the demonstration of a chirped-pulse amplification laser was first made in 1985 [1], the amount of research in laser driven particle accelerators (LPAs) has grown tremendously. Particularly in the past 15 years, there has been significant research into laser wakefield acceleration (LWFA) [2–4] and its potential for high energy radiation production [5–8]. While the tunability, stability, and consistency of these LPAs still lag behind traditional accelerators, recent work has demonstrated successful electron acceleration up to a few GeV [9] and creative ways to generate MeV-level photons [10, 11]. With electron energies reaching the multi-GeV level and photon production ranging from a few keV to 10’s of MeV, laser-plasma accelerators are capable of performing important research on a scale that is thousands of times smaller than traditional accelerator technology.

This thesis describes several experiments carried out on both the HERCULES laser at the Center for Ultrafast Optical Science at the University of Michigan and on the Astra-Gemini laser system at Rutherford Appleton Labs in the United Kingdom. The goal of this thesis is to present empirical evidence of laser wakefield accelerators as an all-in-one tool capable of rivaling the research carried out at linear accelerators, synchrotrons, and Compton gamma ray facilities.

Experiments presented here were carried out to both understand the laser wakefield accelerator as a high energy radiation source and to demonstrate its ability to perform cutting edge research. In the first experiment, measurements of trapped and un-trapped electrons are used to seek a better understanding of electron motion and trapping in the wakefield structure. Additionally, x-rays produced via betatron oscil-
lations in the plasma bubble were studied under a variety of experimental conditions. The working gas was switched between helium and N₂-doped helium showing that ionization injection (a trapping mechanism) on average, results in the production of lower energy photons. This experiment was performed to characterize the radiation produced through laser wakefield acceleration on HERCULES while developing x-ray detection diagnostics to be used on future work.

The next experiment utilized the x-ray detector developed earlier to measure a blue shift in the K-edge of heated aluminum using absorption spectroscopy. The experiment was carried out in a pump-probe setup with a split laser beam being used to heat an aluminum target with one arm and produce betatron x-rays through LWFA with the other arm. By varying the delay between the two beam lines, changes in the aluminum absorption K-edge were measured on the femtosecond timescale. This work demonstrates the benefits to performing pump-probe experiments using an all-optical source.

The final piece of work presented here is an inverse Compton scattering experiment that was performed on the Astra-Gemini laser. A LWFA electron beam was scattered from a counter-propagating laser pulse which resulted in a measured radiation reaction of the electrons along with the production of gamma rays in excess of 100 MeV. A gamma ray spectrometer consisting of an array of cesium-iodide scintillation crystals was developed to calculate the gamma ray spectrum and measure the effects of radiation reaction for the first time.
CHAPTER I

Introduction

1.1 Motivation

High-energy photons have been a vital research tool in the physics community since the discovery of x-rays in 1895 by Wilhem Röntgen [12]. This discovery has led to significant research dedicated to creating tunable light sources generated by oscillating charged particles. Over the past several decades, linear accelerators (LINAC) and synchrotrons have been used for much of the research in charged particle acceleration and high-energy photon production. Synchrotrons and LINACs are conventional particle accelerators that use periodic electric fields along a vacuum tube to add energy to particles. Synchrotrons send the particles into a large circumference storage ring where they cycle continuously and LINACs consist only of a straight tube. One of the primary applications of these accelerated particles has been the generation of high-energy photons. These light sources are capable of producing extremely bright, x-rays as short as 10 femtoseconds in duration. The ultrashort nature and high brightness of the pulses allows these facilities to perform ground-breaking research such as measuring atomic vibrations [13] or determining the atomic structure of a virus [14, 15].

In the past few decades, new methods of particle acceleration have emerged revolving around high intensity lasers. In the past 10 years in particular, accelerating
electrons with high intensity lasers, called laser wakefield acceleration (LWFA), has led to new methods of high-energy photon generation. This is a process by which electrons are accelerated to relativistic energies in the strong electric fields produced by laser-driven plasma waves. A plasma can support electric fields that are several orders of magnitude stronger than linear accelerators and therefore only requires a few millimeters to a few centimeters of acceleration distance to achieve the same energy as the few kilometers required for LINACs. As shown in Figure 1.1, the immense difference in scale and therefore land, material, personnel and monetary resources between these two technologies provides a strong motivation for the research into laser-plasma accelerators.

While synchrotrons and linear accelerators are quite mature technologies with their physics, control of parameters, and limitations having been documented for decades, LWFA is a relatively new technology that will benefit greatly from further research dedicated to control and applications. Linear accelerator facilities, such as

**Figure 1.1:** (a) Three kilometer long SLAC facility [16] and (b) 2.5 centimeter gas cell used to contain the plasma channel in LWFA.
the Stanford Linear Accelerator (SLAC), can create high energy photons by sending the accelerated electron beam through a series of magnets called an undulator that serves to wiggle the beam at a frequency that produces stimulated radiation emission in the x-ray regime. These facilities are known as free electron lasers (FEL). While this gives scientists a significant amount of control over the photon production, it requires significantly more land and resources to achieve this. For example, the Linear Coherent Light Source (LCLS) is an additional 1 kilometer facility built on to the already 2 kilometer long SLAC accelerator.

Some of the other facilities with which laser-plasma accelerators (LPAs) provide potential advantages are synchrotrons and Compton sources. Synchrotrons function by using strong electric fields to accelerate and powerful magnets to contain electrons in a circular orbit rather than in a straight line. These facilities consist of an electron gun, LINAC, and booster region where the electrons gain energy up to relativistic energy before they are moved to a large circumference outer ring where they cycle continuously, emitting radiation as they pass by bending magnets which keep them on track [17]. Although not as large as linear accelerators, synchrotron facilities such as the Diamond Light Source (DLS) at the Rutherford Appleton Laboratory (RAL) in Harwell, England, still have a very large footprint with the DLS having a circumference greater than 500 meters. Unlike the coherent radiation produced by undulators in a free electron laser (FEL), the radiation coming from a synchrotron is broadband and characterized by a critical energy (8.4 keV at the DLS [18]). Due to the broadband nature of the source, double crystal monochromators are used at the DLS to produce a monochromatic x-ray beam between 2 and 25 keV for a greater degree of experimental flexibility.

For research requiring higher energy photons (>1 MeV), Compton backscatter sources such as the High Intensity Gamma-Ray Source (HIGS), operated by the Triangle Universities Nuclear Laboratory, are capable of generating gamma ray beams
between 2 and 100 MeV. The HIGS facility consists of an FEL connected to a storage ring greater than 100 meters in circumference. An electron beam is accelerated and sent through an undulator magnet capable of generating light between 193 and 1064 nm depending on the desired gamma ray energy. The light from the FEL is reflected back with the proper timing to oscillate with the electrons on subsequent passes through the undulator. A secondary beam of electrons is injected and collides head-on with the FEL light; this inverse Compton scattering interaction with the relativistic electron beam generates the desired gamma ray beam [19, 20]. The unique aspect of the HIGS facility, is its ability to create monochromatic gamma ray beams with an energy spread of less than 1% [21]. The high degree of control in the energy and spread of these photon beams makes HIGS an ideal facility for experiments investigating nuclear resonant fluorescence [22, 23]. Further detail about the physics of inverse Compton scattering is discussed Section 2.3.2 and an experiment is presented in Chapter VII.

The drawback to these facilities is their size and cost. The physical space and capital required to build, maintain, and operate these facilities is substantial, which is why they only exist in a few locations worldwide. Some of the early work in LWFA has shown that laser-plasma accelerators are also practical photon sources capable of creating broadband spectra ranging from a few keV to 100 keV [5, 24–27]. Because LWFA is produced with an all-optical source, more creative methods of photon production have been explored in recent years to extend the range of photon energies that can be achieved from a laser-plasma source [8, 10, 11, 28]. One of the ultimate goals of laser-plasma accelerator research is to create a tunable high-energy photon source that is compact, economical and portable. While LWFA technology and optical-based photon sources still lag behind LINACs, synchrotrons and Compton sources for many parameters, Table 1.1 shows that the values for laser-plasma accelerators are comparable in some categories allowing for competitive research to currently be carried out.
### Table 1.1: Comparison of various electron and x-ray parameters achievable from SLAC/LCLS and LWFA sources [7, 9, 26, 29–31].

<table>
<thead>
<tr>
<th>Parameter</th>
<th>LCLS</th>
<th>LWFA</th>
</tr>
</thead>
<tbody>
<tr>
<td>Peak Energy</td>
<td>13.6 GeV</td>
<td>4 - 5 GeV [9]</td>
</tr>
<tr>
<td>Charge per Bunch</td>
<td>0.25 nC</td>
<td>0.5 nC [26]</td>
</tr>
<tr>
<td>Repetition Rate</td>
<td>30 Hz</td>
<td>~0.1 Hz [26]</td>
</tr>
<tr>
<td>Photon Energy</td>
<td>8.3 keV</td>
<td>~10 keV (peak), ~100 keV (max) [7]</td>
</tr>
<tr>
<td>Pulse Duration</td>
<td>70 - 100 fs</td>
<td>30 fs [29]</td>
</tr>
<tr>
<td>Photons per Pulse</td>
<td>1.0 - 2.3 \times 10^{12}</td>
<td>\sim 10^9 [26]</td>
</tr>
<tr>
<td>Peak Brightness</td>
<td>20 \times 10^{-32}*</td>
<td>1 \times 10^{-22}* [26]</td>
</tr>
<tr>
<td>Acceleration Gradient</td>
<td>10 - 50 MeV/m</td>
<td>10 - 100 GeV/m [30]</td>
</tr>
</tbody>
</table>

*Photons per phase space volume, or photons s\(^{-1}\) mm\(^{-2}\) mrad\(^{-2}\) per 0.1% spectral bandwidth.

There are numerous practical applications of high energy x-rays and gamma rays including medical and biological imaging, microscopy, radiography, activation, and homeland security [28, 32, 33]. The development of laser-plasma photon sources allows for more widespread research and applications of this technology at hospitals and universities rather than being confined to large and expensive accelerator facilities.

#### 1.2 High-Intensity Lasers

The scientific breakthrough that allows this unique field of research to exist was the development of high-intensity lasers. Laser is an acronym for Light Amplification by Stimulated Emission of Radiation and the technology has been around since the invention of the ruby laser in 1960 [34]. For years after the ruby laser was invented, researchers continued experimenting with various laser gain (or amplification) media such as CO\(_2\) gas, liquid dye lasers, solid state lasers, and fiber lasers. Today, solid state crystal lasers dominate ultrashort, high-intensity laser research. The highest intensity lasers in the world today are typically Ti:Sapphire lasers pumped by various neodymium:glass lasers and are capable of reaching intensities as high as \(2 \times 10^{22} \text{ W/cm}^2\) [35]. For about 20 years after the invention of the ruby laser, the
Figure 1.2: Diagram of maximum laser intensity plotted against the year. The plateau in the middle is the result of the laser intensities reaching the damage threshold of most lab materials ($\sim 10^{15} \text{ W/cm}^2$). Figure adapted from [36].

The maximum power of lasers was steadily increasing to a plateau until 1985 when Donna Strickland and Gerard Mourou developed a technique called chirped-pulse amplification (CPA) [1]. The plateau of laser intensities between 1970 and 1980 is due to the fact that pushing beyond $\sim 10^{15} \text{ W/cm}^2$ would cause damage to the laser gain media. Chirped pulse amplification is able to circumvent this problem by stretching the pulse in time before amplifying, thus increasing the time frame over which energy is transferred between the gain medium and laser pulse and therefore reducing the intensity in the amplification medium. Figure 1.2 shows the impact that CPA had on achievable laser intensities. This has allowed high energy laser pulses (several joules) to be delivered in an ultrashort pulse duration (10’s of femtoseconds).

Figure 1.3 is a diagram depicting the CPA process. The technique starts out by sending a short pulse produced in the cavity of a laser oscillator into a pulse stretcher, typically a pair of gratings. The gratings disperse the incident light into its frequency components so that the optical path length (OPL) for the red frequencies is shorter.
than that of the blue. This results in a stretched pulse with low power exiting the
first grating pair that can be amplified because its intensity will be well below any
damage thresholds. After amplification, the high-energy, long pulse is sent into a set
of the initial gratings oriented in reverse, in order to undo the stretching and thus
compress the pulse to a very short pulse duration [37]. The result of this is a very
high power laser pulse, delivering several joules in a few 10’s of femtoseconds (fs).

The duration of the final amplified pulse is dependent on several factors throughout
the CPA process, but the primary variable contributing to the output pulse duration
is the bandwidth of the initial pulse coming from the oscillator. The duration of the
pulse can be pushed to shorter timescales by increasing the mode locked bandwidth,
$N\Delta\nu$, where $N$ is the number of locked modes oscillating in phase and $\Delta\nu$ is the
frequency separation of each of those modes. This mode-locked bandwidth or gain
bandwidth is the frequency range over which substantial amplification can occur in
the laser gain medium. The re-compression of the amplified laser is slightly hindered
by gain narrowing in the amplifiers so the final pulse duration will not quite match
the initial bandwidth. For a Gaussian temporal shape, the transform-limited pulse
duration is: $\Delta t = 0.441/N\Delta\nu$. For a Ti:Sapphire laser, this gain bandwidth can be
over 300 nm producing a pulse duration of less than 8 fs from the oscillator and $\sim 20$
fds after amplification.

While high-powered, CPA laser systems are an incredible scientific breakthrough
and have opened the door to studying new and exciting physics, there are several diffi-
cult aspects of these systems that make them a challenging laboratory tool. Chirped-
pulse lasers are highly complex systems that typically require multiple stages of am-
plification, each of which usually has several passes (this is described in Section 3.1).
As the beam is amplified, it becomes necessary to increase the diameter in order to re-
main below the damage threshold of the amplifier crystals and optics. The numerous
amplifier passes and beam expansion stages means that the entire laser system will be
Figure 1.3: A simplified diagram depicting the chirped pulse amplification process. A short pulse is stretched using a pair of gratings to disperse its frequency components. The long pulse is then amplified through some gain medium before compression takes place using backwards propagation through another grating setup. The result is an ultrashort duration, high power laser pulse.

composed of hundreds of individual optics creating a very complex system that will require a significant amount of alignment and attention. The beam size also creates a cost limiting factor as the amplification crystals must be as large or larger than the beam size. The amplification crystals, such as Ti:Sapphire, are typically one of the most expensive components of a high powered laser and their cost of production significantly increases as the beam size is increased to accommodate higher energy pulses.

1.3 Dissertation Outline

- Chapter I is an introduction to the thesis and describes some of the background motivation for the research presented here. There is also a description of high powered lasers because they are the primary tool that makes all this research possible.
• Chapter II describes much of the fundamental physics of laser plasma interactions and provides a theoretical basis for laser wakefield acceleration.

• Chapter III provides information about the laser systems and experimental diagnostics used in the research described in this thesis.

• Chapter IV presents on-axis measurements of the electron beam and discusses un-trapped electrons in a ring distribution.

• Chapter V describes experimental work centered around characterizing the LWFA electrons and betatron x-rays.

• Chapter VI describes a pump-probe experiment on HERCULES in which measurements of the aluminum K-edge transmission were made.

• Chapter VII is about an inverse Compton scattering experiment on the Astra-Gemini Laser.

• Chapter VIII presents conclusions and directions for future work.

The work presented in this thesis is part of a large collaboration of students, research scientists, and professors between the University of Michigan, the Imperial College of London, Lancaster University, and the Astra-Gemini Laser at Rutherford Appleton Laboratory. The author contributed to performing all of the research presented in this thesis as both a team member and project leader depending on the experiment. With the exception of the electron spectrometer analysis in Chapter VII, all of the analysis of experimental data in Chapters IV – VII was performed by the author. The simulation work in Chapters IV and VI were performed by colleagues at Lancaster University and Imperial College of London respectively.
CHAPTER II

Fundamental Physics of Laser-Plasma Accelerators

2.1 Ionization by Intense Lasers

As discussed in Section 1.1, one of the unique and most appealing aspects of laser-plasma accelerators is their ability to support acceleration gradients many orders of magnitude higher than traditional accelerators. This is possible because the electrostatic fields exist in a plasma – a medium consisting of ionized atoms with the electrons and nuclei unbound. Whereas in a traditional accelerator, the electric field is limited to the ionization potential of the surrounding materials. This section covers the fundamental mechanisms by which a high intensity laser pulse can ionize a gas into a plasma.

The simplest method of ionization is the photoelectric effect, which was first described by Albert Einstein in 1905 [38]. This process occurs when a single photon transfers its energy to an electron through an inelastic collision and the electron gains enough energy to free itself from the Coulomb bond holding it to the nucleus as shown in Figure 2.1(a). In order for this to occur, the energy of the photon must exceed the ionization energy of the electron with which it interacts. Since most high intensity lasers operate in the near-infrared to infrared regimes, the photon energies (around 1 - 2 eV max) will not be sufficient for a single photon to raise the electron out of its potential well. However, there exists a nonlinear version of this process called multi-
 photon ionization (MPI) in which several photons can be absorbed by an electron at the same time thus providing enough energy for the electron to escape from the nucleus [39], as illustrated in Figure 2.1(b).

Although MPI exists in the experiments discussed in this thesis, the intensity of the lasers is strong enough for tunneling or barrier suppression ionization to be the dominant mechanism. When the electric field of the focused laser is sufficiently strong, it enters a new regime where the field of the laser distorts or suppresses the Coulomb field that exists between the electron and the nucleus. If the Coulomb field is partially suppressed as shown in Figure 2.1(c), the electron has a probability of escaping or “tunneling” out of the potential well. This change in regime was first described by Keldysh in 1965 and is determined by Equation (2.1) [40]. Where $\omega_L$ is the frequency of the laser, $E_L$ is the electric field of the laser, $E_{ion}$ is the ionization potential, $e$ is the electron charge, and $m_e$ is the electron mass.

$$
\gamma \sim \frac{\omega_L}{eE_L}\sqrt{4m_eE_{ion}}
$$

(2.1)

When the parameter $\gamma$ is less than 1, the tunneling time is less than the laser period and the dominant mechanism is tunneling and when $\gamma$ is greater than 1, the dominant mechanism is MPI. Barrier suppression ionization occurs when the laser field is even
stronger, resulting in full suppression of the Coulomb field. In this situation, the
electron is free to escape as if the Coulomb field was not present.

2.2 Laser Wakefield Acceleration

Electron acceleration on a “table-top” scale with high intensity lasers was first pro-
posed by Tajima and Dawson in 1979 [41], but monoenergetic electron beams were
not produced by this mechanism until 2004 [2–4]. The process of laser wakefield accel-
eration occurs when a sufficiently intense laser pulse enters a plasma and accelerates
electrons by trapping them in the plasma wave created by the laser. There are a few
key physical phenomena that occur when an intense laser pulse enters into a plasma
that allow laser wakefield acceleration to occur. The ponderomotive force and laser
self-focusing are two important aspects of laser-plasma interactions that generate the
plasma waves necessary for LWFA.

2.2.1 Single Particle Dynamics

To better understand how a plasma responds to an incident laser, it is beneficial
to start by considering a single electron. In the case of a fully ionized plasma, the
electrons will move according to internal or externally applied electromagnetic fields.
The Lorentz equation describes electron motion in a plasma, where $u$ is the electron
velocity, $p$ is the electron momentum, and $E$ & $B$ are electric and magnetic fields.

$$\frac{\partial p}{\partial t} + u \cdot \nabla p = -e (E + u \times B)$$  \hspace{1cm} (2.2)

If an external, time-varying electromagnetic field is applied to the plasma, such
as a laser, then we can rewrite Equation 2.2 as Equation 2.3, where $p = m_e u$. For
the sake of illustrating how the laser electric field affects electron motion, a low-field
limit is assumed so the magnetic field will be ignored.

\[
\frac{\partial \mathbf{u}}{\partial t} + \mathbf{u} \cdot \nabla \mathbf{u} = -\frac{e}{m_e} \mathbf{E}(x) \sin(\omega t) \tag{2.3}
\]

Linearizing the above equation to the lowest order (denoted by a 1) and integrating with respect to time yields the electron quiver velocity in the electric field:

\[
\mathbf{u}_1 = \mathbf{v}_{osc} = \frac{eE}{m_e \omega} \cos(\omega t) \tag{2.4}
\]

This oscillation velocity describes fast electron motion and because this scenario involves a spatially-dependent and time-varying electric field, the electron velocity can be divided into a fast \((f)\) and slow \((s)\) component such that: \(\mathbf{u} = \mathbf{u}_f + \mathbf{u}_s\), and \(\mathbf{u}_f \gg \mathbf{u}_s\). When these equations are substituted into Equation 2.3 and time averaged over a laser cycle, the \(\mathbf{u}_f\) term disappears because the average of cosine over one cycle is equal to 0. The resulting equation describes a force felt by the electrons due to the spatial variance of the electric field over the quiver amplitude known as the ponderomotive force:

\[
m_e \frac{\partial \mathbf{u}_s}{\partial t} = F_p = -\frac{e^2}{4m_e \omega^2} \nabla (E^2) \tag{2.5}
\]

When the pulse intensity is sufficiently strong that the electron quiver motion becomes relativistic and \(\gamma\) becomes relevant, the equation for the ponderomotive force is recast for the relativistic correction. Equation 2.6 shows the relativistic form of the ponderomotive force [42].

\[
F_{p,rel} = -m_e c^2 \nabla (\gamma - 1) \tag{2.6}
\]

Because this force scales with intensity, it is helpful to relate \(\gamma\) to the intensity of the
laser through $a_0$, the normalized laser potential:

$$
\gamma = \sqrt{1 + a_0^2}
$$

(2.7)

for circularly polarized light. The normalized laser potential is a dimensionless quantity used to described the intensity of the laser pulse given by: $a_0 = eA_{laser}/(mc^2)$. This parameter can be related to the laser intensity by the following equation:

$$
a_0 = \sqrt{\frac{I[W/cm^2]\lambda^2[\mu m]}{1.37 \times 10^{18}}}
$$

(2.8)

This force acts to push electrons out of its path while leaving the ions relatively unperturbed due to their higher mass. The existence of this force is due to the non-uniform spatial profile of the laser pulse electric field. An electron interacting with the laser field is driven to regions of lower intensity where the electric field is not strong enough to push the electron back to its original position, thus resulting in net electron motion away from areas of highest intensity [43, 44]. The perturbation to the plasma caused by this force creates the plasma waves that are used to trap and accelerate electrons. The perturbation continues for several plasma wavelengths behind the laser pulse, resulting in a traveling plasma wave of electron density peaks and valleys (or “buckets”).

2.2.2 Self-Focusing

When a laser pulse focuses in vacuum, the phase velocity at the edges of the pulse travel faster than the center of the pulse creating a curved phase front. The beam waist of the pulse decreases according to Equation 2.9 and remains roughly focused over the Rayleigh length; $w_0$ is the beam waist before focusing, $z$ is the distance from the focusing optic, and $\lambda_L$ is the laser wavelength. For a Gaussian pulse, the Rayleigh
length is defined as the distance from the focal spot to where the area of the pulse is
twice as large \cite{45} and is given by: $z_R = \pi w_0^2/\lambda L$.

\[
w(z) = w_0 \sqrt{1 + \frac{\lambda_L z^2}{\pi^2 w_0^4}} \tag{2.9}
\]

It is evident from these equations that in order to keep the laser pulse focused over
a longer distance (which is desired in LWFA), it is necessary to increase the size
of the focal spot. This however, will result in a decrease in intensity, so the laser
power must also be increased accordingly. A laser propagating through a plasma is
capable of remaining focused over several Rayleigh lengths due to the self-focusing
phenomenon. Although self-focusing is enhanced by the plasma’s response to the
laser, the phenomenon itself can arise from spatial non-uniformity of the laser pulse
envelope or relativistic effects of plasma electrons. The high intensity of the center of
the envelope results in a nonlinear change of the index of refraction that is different
from the wings of the pulse where the intensity is lower. The modulated refractive
index across the beam waist causes the wings to curve inwards resulting in further
self-focusing of the laser wavefront. In a plasma however, the density is variable and
the depression created by the evacuation of electrons by the ponderomotive force
further enhances self-focusing by creating a local increase in the refractive index.
With a sufficiently intense laser pulse, the relativistic mass increase of the electrons
oscillating in the laser field near to the speed of light also contributes to a change in
the index of refraction of the plasma. This self-focusing phenomenon acts to further
focus the laser pulse beyond the diffraction limit of the focusing optic if the laser
power is above the critical power shown in Equation (2.10) \cite{46–49}:

\[
P_{cr} = \frac{m_e c^5 \omega}{e^2 \omega_{pe}^2} \simeq 17 \left( \frac{\omega}{\omega_{pe}} \right)^2 \text{ [GW]} \tag{2.10}
\]
Figure 2.2: Electron density from a simulation of a ultrashort-pulse laser entering a plasma. The red circle shows the ion cavity created by the ponderomotive force expelling electrons. The red arrows show the static electric field set up by the charge separation between the ion cavity and the surrounding electrons.

where $\omega_{pe}^2 = e^2 n_e / \epsilon_0 m_e$ is the plasma frequency and $c$ is the speed of light. The increased intensity of the pulse strengthens the ponderomotive force creating a positive feedback loop for relativistic self-focusing. As the laser spot gets smaller, diffraction increases and takes over causing the beam to expand to the point that self-focusing becomes dominant again leading to a decrease in the laser spot. This oscillation amplitude is small compared to the overall beam waist of the radiation envelope and keeps the laser guided over the length of the plasma rather than immediately expanding due to diffraction [46]. It is possible, however, to operate at a point where self-focusing and diffraction are in equilibrium, thus eliminating the oscillation of the pulse envelope. This matched condition occurs when the laser focal spot is:

$$w_{matched} \simeq 2 \sqrt{a_0 c / \omega_{pe}} [50].$$

These two features of laser-plasma interactions create an environment for the
generation of relativistic plasma waves. As the pulse is traveling through a plasma, it pushes electrons out of its path both radially and axial, creating plasma waves and leaving a cavity of positive charge trailing the pulse. This charge separation results in a radially symmetric electromagnetic field between the expelled electrons and trailing ions. As the electrons oscillate back together forming the “high electron density” region shown in Figure 2.2 they can become trapped in the potential well created by the charge separation at the back of the cavity. Electrons caught at the back of the “bubble” or cavity are subject to the pull of the longitudinal component of the electric field created by the mass of positive charge between them and the laser pulse.

2.2.3 Laser Wakefield Regimes

One of the main parameters the determines the physics of the LWFA process is the laser pulse duration. There are three main regimes where LWFA can take place: the self-modulated, forced, and bubble (or blowout) regimes [51–53]. The three different regimes are defined by comparing the laser pulse duration to the plasma wavelength. The most effective method for accelerating electrons in the plasma wake occurs in the bubble or blowout regime when the laser pulse length \( c\tau_0 = L_{\text{laser}} \) is less than or equal to half the plasma wavelength \( L_{\text{laser}} \leq \lambda_p/2 \) [41]. Operating in the bubble regime has only been possible in recent years with the implementation of ultrashort (sub-100 fs), high-intensity lasers and is the regime in which all the experiments in this dissertation were conducted. In the highly non-linear regime of short-pulse LWFA, the plasma wave oscillations are so strong that the wave breaks after the first oscillation leaving only a single ion-cavity trailing the laser pulse. The phenomenon of “wave-breaking” in a plasma wave oscillation occurs when the wave form (denoted by the dashed line in Figure 2.2) steepens to the point that the peak of the wave appears to fall into the trough resulting in a collapse of the plasma wave structure and injection of electrons into the ion cavity [53, 54]. For reference, the self-modulated regime occurs
when the laser pulse length is greater than the plasma wavelength ($L_{\text{laser}} > \lambda_p$). This regime is dominated by an instability in which the laser pulse is modulated in time resulting in a series of smaller pulses separated by the plasma frequency. Resonant excitation of the train of laser pulses occurs and there is an enhancement in the acceleration fields over those generated by lasers with $L_{\text{laser}} \approx \lambda_p/2$ assuming fixed laser parameters [51, 55]. The forced regime occurs when the laser pulse length is approximately equal in length to the wavelength of the plasma wave ($L_{\text{laser}} \sim \lambda_p$); this was demonstrated to be an improvement over self-modulated experiments in the early years of LWFA. In the forced LWFA regime, the laser pulse drives a “bubble shaped” plasma cavity longer than $\lambda_p$. However, because the back of the laser pulse is able to travel faster than the front due to the density depression created by the leading edge of the pulse, the compression is an optical shock that can drive the plasma wave beyond its wave-breaking limit [52].

### 2.2.4 Trapping Mechanisms & Acceleration

In order for acceleration to occur, the electrons oscillating in the plasma wave must become trapped by the traveling wave and pulled into the ion cavity. Although there are many different methods to induce electron injection into the bubble, such as density down-ramp injection [56, 57], shock injection [58, 59], and colliding pulse injection [60, 61], self-injection is the most fundamental that can occur without external manipulation of the laser wakefield environment. In a uniform plasma created from a gas with a low ionization threshold such as helium, self-trapping occurs when the electrons enter at the back of the bubble or ion cavity and get caught if they have sufficient forward momentum. In order to properly create a wake with a strong enough potential well for the electrons to remain trapped, the intensity of the laser must be high enough for the ponderomotive force to amply expel electrons and create a large enough bubble. Simulations and analytical calculations have shown that
**Figure 2.3:** Particle orbits in phase space of an electron traveling in a plasma wave with a normalized potential: \( \phi = \phi_0 \cos \omega t \). The solid line is the separatrix and the dashed lines are orbits of trapped and un-trapped electrons [63].

A normalized laser potential of at least 3 \((a_0 > 3)\) is necessary for self-trapping to occur [50, 62].

Since the plasma wave is traveling forward at a speed, \( v_\phi \), the electrons must have the correct orbit and also be traveling at \( v_\phi \) when they pass by the back of the bucket to become trapped in the wake. Figure 2.3 shows a visual representation of electron orbits in a linear plasma wave plotted with the separatrix. In this context, the separatrix is the dividing line between trapped and un-trapped electrons [63, 64]. From this plot it is easier to see two things: the trapping of electrons depends on the electric field of the plasma wave, because this affects the separatrix, and trapping of electrons occurs when their orbit overlaps with the separatrix. In one dimension, this overlap of the separatrix and the electron plasma oscillations occurs at:

\[
E_{WB} = \sqrt{2} (\gamma_p - 1)^{1/2} E_0 \gg E_0
\]  
(2.11)
where $E_{WB}$ is the cold relativistic wave-breaking field, $E_0 = cm_e \omega_p/e$ is the cold nonrelativistic wave-breaking field and $\gamma_p = \left(1 - v_p^2/c^2\right)^{-1/2}$ is the Lorentz factor associated with the phase velocity of the plasma wave [63]. Because the plasma wave is traveling very near the speed of light, it is difficult for the background plasma electrons to reach a speed where their orbit will overlap with the separatrix. For this reason, other methods of injection are being explored among the LWFA community as mentioned previously.

Another trapping method related to this thesis involves using a high-Z gas mixture to lower the trapping threshold via ionization induced trapping [6]. With this method, a high-Z gas or mixture of gases, such as nitrogen, is used as the working medium. Due to the stronger Coulomb forces between the nucleus and electrons in nitrogen, the inner-most electrons are not ionized by the pre-pulse or leading edge of the laser pulse as is the case with helium. Ionization of the inner most electrons occurs with the peak of the pulse and therefore results in electrons being “born” on to a trapped orbit (dashed line in Figure 2.3) where they can receive a longitudinal push from the back portion of cavity’s electric field, thus lowering the threshold necessary for electron trapping [6].

Though there are numerous schemes that have been explored to control electron trapping through the use of multiple lasers, multiple pulses or injector beams [60, 61, 65–67], the two methods described above are the most fundamental and simplest forms of trapping that don’t require any external manipulation of the LWFA system.

### 2.2.5 Laser Wakefield Scalings

Once the electrons become trapped at the back of the cavity, they will experience a very strong longitudinal electric field resulting from the bubble of positive charge between the laser and the newly trapped electrons. The structure of the field in the longitudinal direction along several wake buckets can be seen in Figure 2.4(a). As
depicted in this figure, dephasing occurs when the electrons reach the center of the bubble and experience an equal and opposite electric field pushing them backwards. This means that the energy gain of the electrons will only occur over the first half of the bubble. It is possible to determine the energy gain by first deriving the dephasing length because $\Delta E \propto eE_z L_{deph}$, where $E_z$ is the longitudinal electric field.

Starting with the phase velocity of the wake traveling through the plasma [68]:

$$v_\phi = c \left( 1 - \frac{3\omega_{pe}^2}{2\omega_0^2} \right)$$  \hspace{1cm} (2.12)

To determine the dephasing length, it is necessary to find the difference in speed between the electrons and the wake. Since the laser, and therefore the wake, is traveling slower than the speed of light and the electrons can outrun the laser pulse, they can be assumed to roughly be traveling at the speed of light: $v_e \approx c$. Therefore, the difference in speed between the electrons and plasma wake is going to be:

$$\Delta v = v_e - v_\phi = c \left( 1 - 1 + \frac{3\omega_{pe}^2}{2\omega_0^2} \right) = c \frac{3\omega_{pe}^2}{2\omega_0^2}$$  \hspace{1cm} (2.13)

Working with the matched focal spot condition, the shape of the bubble can be assumed spherically symmetric and the radius of the bubble can be assumed equal to
the beam waist. Knowing that the electrons dephase once they reach the middle of the bubble, the time to dephasing can be calculated using the radius of the bubble:

$$t_{\text{deph}} = \frac{r_b}{\Delta v} = \frac{2\sqrt{a_0} \frac{c}{\omega_{pe}}}{3\omega_{pe}^2 \omega_{pe} \sqrt{a_0}} = \frac{2\lambda_p \omega_0^2}{3\pi c\omega_{pe}^2 \sqrt{a_0}}$$  \hspace{1cm} (2.14)$$

where $r_b$ is the radius of the bubble and $\lambda_p = 2\pi c/\omega_{pe}$ is the relativistic plasma wavelength. Simply multiplying this by the speed of the electrons, $c$, yields the dephasing length ($L_{\text{deph}}$) of electrons accelerating in a matched bubble:

$$L_{\text{deph}} = \frac{2\lambda_p \omega_0^2}{3\pi \omega_{pe}^2 \sqrt{a_0}}$$  \hspace{1cm} (2.15)$$

Now that the length over which the electrons travel to reach maximum energy is defined, it is easy to calculate their energy gain by calculating the strength of the electric field that causes the acceleration. As shown in Figure 2.4(a), the longitudinal field in 1D is linear and can therefore be written as: $\partial E_z/\partial z = e n_p/\epsilon_0$, where $n_p$ is the plasma density; integrating this along $z$ yields: $E(z) = e n_p z/\epsilon_0$. At the back of the bubble where $z = r_b$, the field strength is: $E(r_b) = 2\frac{en_p}{\epsilon_0} \sqrt{a_0} \frac{c}{\omega_{pe}} = 2\omega_{pe} \frac{mc}{e} \sqrt{a_0}$. This is the maximum field strength felt by the electrons which decreases linearly as they are move towards the center of the bubble. To find the energy gain of the electrons, it is necessary to use the average electric field, which in the linear case will simply be $\frac{1}{2}E_{\text{max}}$. It is more useful to look at the energy gain of a 3D nonlinear wakefield, in which case the average electric field is much more complicated to derive as the magnetic field contribution is not negligible. From Lu et al. [50], the average accelerating field for the 3D nonlinear case differs by a factor of $\frac{1}{2}$: $E = \frac{1}{2}\omega_{pe} \frac{mc}{e} \sqrt{a_0}$. Using all of this together, the energy gain of the electrons is given in Equation 2.16 for the 3D nonlinear case [50]. The energy gain for different regimes can be found in the last column of Table 2.1.
\[
\Delta E \simeq \frac{2}{3} m_e c^2 \left( \frac{\omega_0}{\omega_{pe}} \right)^2 a_0 \simeq m_e c^2 \left( \frac{P}{m_e^2 c^5/e^2} \right)^{1/3} \left( \frac{n_c}{n_p} \right)^{2/3} \]

\[
\Delta E \text{ [GeV]} \simeq 1.7 \left( \frac{P[TW]}{100} \right)^{1/3} \left( \frac{10^{18}}{n_p \text{ [cm}^{-3}] \right)^{2/3} \left( \frac{0.8}{\lambda_0 \text{ [\mu m]}} \right)^{4/3} \quad (2.16)
\]

Since dephasing occurs when the electrons remain trapped past the point of maximum energy gain, it is ideal to have the electron beam exit the plasma as soon as maximum energy is achieved. This however, is not so easily controlled so significant work has been performed on understanding the physics of plasmas longer than the dephasing length. When electrons remain trapped past the dephasing length, they undergo longitudinal oscillations as they are slowed down in the front half of the bubble and then re-accelerated in the back half. The electrons also interact with the laser pulse at the front of the bubble if they remain trapped for too long. Although interacting with the laser pulse ruins the quality of the electron beam, the electric field of the laser will force the electrons to undergo additional oscillations thus producing more x-ray photons [69].

Another important aspect to consider in LWFA experiments, is the depletion length. Depletion occurs when the laser pulse has lost enough energy that it can no longer sustain the creation of plasma waves and the ion cavity fades [70]. Simulations performed by W. Lu et al. shown in Figure 2.4(b) – (d) depict the progression of the wakefield bubble as pump depletion occurs over several millimeters of propagation. If the beam remains in the surrounding plasma past the depletion length, the electron beam generated by LWFA can continue to drive a wake of its own. This process, known as plasma wakefield acceleration (PWFA), occurs when the electron beam becomes the plasma wave driver and can lead to additional trapping and acceleration.
Table 2.1: Dimensionless scalings of various laser wakefield parameters. The values in the $a_0$ column describe the range for which linear, 1D nonlinear and 3D nonlinear are valid and the remaining columns describe how each of the parameters scale in that regime. $k_p w_0$ is the scaling for the ion channel force, $\epsilon_{LW}$ is the average accelerating field, $k_p L_d$ and $k_p L_{pd}$ are the dephasing and depletion length scalings, respectively and $\Delta W$ is the normalized energy gain of the electrons [50].

<table>
<thead>
<tr>
<th>Regime</th>
<th>$a_0$</th>
<th>$k_p w_0$</th>
<th>$\epsilon_{LW}$</th>
<th>$k_p L_d$</th>
<th>$k_p L_{pd}$</th>
<th>$\Delta W/ (mc^2)$</th>
</tr>
</thead>
<tbody>
<tr>
<td>Linear:</td>
<td>&lt; 1</td>
<td>$2\pi$</td>
<td>$a_0^2$</td>
<td>$\frac{\omega_0^2}{\omega_p^2}$</td>
<td>$\frac{\omega_0^2}{\omega_p^2} \frac{\omega_p \tau}{a_0^2}$</td>
<td>$a_0^2 \frac{\omega_0^2}{\omega_p^2}$</td>
</tr>
<tr>
<td>1D Nonlinear:</td>
<td>&gt; 1</td>
<td>$2\pi$</td>
<td>$a_0$</td>
<td>$4a_0^2 \frac{\omega_0^2}{\omega_p^2}$</td>
<td>$\frac{1}{3} \frac{\omega_0^2}{\omega_p^2} \omega_p \tau$</td>
<td>$4a_0^2 \frac{\omega_0^2}{\omega_p^2}$</td>
</tr>
<tr>
<td>3D Nonlinear:</td>
<td>&gt; 2</td>
<td>$2\sqrt{a_0}$</td>
<td>$\frac{1}{2} \sqrt{a_0}$</td>
<td>$\frac{4}{3} \frac{\omega_0^2}{\omega_p^2} \sqrt{a_0}$</td>
<td>$\frac{\omega_0^2}{\omega_p^2} \omega_p \tau$</td>
<td>$\frac{2}{3} \frac{\omega_0^2}{\omega_p^2} a_0$</td>
</tr>
</tbody>
</table>

2.3 Radiation Production from Laser Wakefield Accelerators

This section describes some of the methods in which high energy photon radiation can be produced from electron sources based on laser wakefield acceleration. The photon sources described in this section were used as investigative tools during the
2.3.1 Betatron Oscillations

One of the unique and notable aspects of a laser wakefield accelerator is that hard x-rays are produced during the acceleration process. Electrons trapped in the bubble can undergo transverse oscillations known as betatron oscillations due to the transverse focusing fields of the bubble as described in Section 2.2. The x-rays produced through this oscillation have been coined “betatron x-rays” and are often described by the dimensionless parameter, $K$, which is a measure of the ratio between the oscillation amplitude and wavelength as shown in Figure 2.5. This value is used to denote the strength of the electron oscillations and if $K \gg 1$, the electrons are oscillating in the “wiggler” regime and if $K \ll 1$, the electrons are in the “undulator” regime [71], as shown in Figure 2.5. This difference in oscillation strength will subsequently affect the number and spectrum of the x-rays produced. The wavelength of the produced radiation is provided by the following equation [71]:

$$\lambda = \frac{\lambda_u}{2\gamma^2} \left( 1 + \frac{K^2}{2} + \gamma^2 \theta^2 \right)$$  \hspace{1cm} (2.17)

where $\lambda_u$ is the wavelength of the sinusoidal oscillation in the rest frame of the electron.

Following the derivation described by Corde et al. [71], it is possible to see how
Figure 2.6: (a) The synchrotron-like spectrum from a single electron (solid) undergoing betatron oscillations at a constant amplitude and a spectrum integrated over a Gaussian distribution (dashed) of betatron amplitudes. (b) The same spectra as (a) plotted on a logarithmic scale to show the effect on higher energy photons [73].

The strength parameter, $K$, will have an effect on the produced x-ray spectrum. For an electron traveling forward while oscillating perpendicular to the direction of travel, it is necessary to consider the effect of the transverse momentum on the longitudinal momentum. When $K \ll 1$, the transverse momentum is small enough that the effect on the longitudinal momentum is negligible, meaning that $\bar{\beta}_z \simeq \bar{\beta}$ and $\gamma_z = \sqrt{1 - \beta_z^2} \simeq \gamma$, where $\bar{\beta}$ and $\bar{\beta}_z$ are the total and longitudinal velocity normalized to the speed of light. With these assumptions, the system is equivalent to an oscillating dipole in which case the resulting spectrum consists of monochromatic x-rays at the oscillation frequency, $\omega$ that have received a double-Doppler upshift by the relativistic speed of the electron [72]. The resulting frequency of the emitted radiation ($\omega_c$) can be written as,

$$\omega_c = 2\gamma^2\omega_u$$

which comes from simplifying Equation 2.17 for $K \ll 1$ and $\theta = 0$. When the oscillations are strong enough such that $K \gg 1$, it is no longer appropriate to assume that the transverse momentum does not have an effect on the longitudinal velocity. In the ultrarelativisitc limit, the harmonics of the oscillating electron merge to generate the
“synchrotron” spectrum shown in Figure 2.6. The derivation of this spectrum is identical to an electron in a circular orbit producing synchrotron radiation as described by Jackson [74] and the radiated power per unit frequency is shown in Equations (2.19 - 2.22).

\[
\frac{dP}{d\omega} = \frac{P_\gamma}{\omega_c} S\left(\omega/\omega_c\right) \quad (2.19)
\]

\[
S(x) = \frac{9\sqrt{3}}{8\pi} \int_x^\infty K_{5/3}(\xi) d\xi \quad (2.20)
\]

\[
P_\gamma = \frac{e^2 c \gamma^4}{6\pi \epsilon_0 \rho^2} = \frac{2e^2 \omega_c^2}{27\pi \epsilon_0 c \gamma^2} \quad (2.21)
\]

\[
\omega_c = \frac{3}{2} \frac{\gamma^3 c}{\rho} \quad (2.22)
\]

The first equation describes the synchrotron spectrum, given in radiated power per unit frequency, \( K_{5/3} \) is a modified Bessel function of the second kind, \( P_\gamma \) is the radiated power (\( P_\gamma = \int (dP/d\omega) d\omega \)), and \( \omega_c \) is the critical frequency of the synchrotron spectrum. To put Equation 2.22 in a more usable format, the radius of curvature of the electron orbit, \( \rho \), is written as: \( \rho = \gamma \frac{\lambda_u}{2\pi K} \) which results in the following equation describing the critical energy of a synchrotron spectrum [71]:

\[
\omega_c = \frac{3}{2} K \gamma^2 \frac{2\pi c}{\lambda_u} \quad (2.23)
\]
2.3.2 Inverse Compton Scattering

With regular Compton scattering, a photon undergoes an inelastic collision with a charged particle resulting in the scattered photon having a different energy than the incident photon. By contrast, inverse Compton scattering (ICS) occurs when high energy photons are generated by scattering charged particles (electrons in this case) from photons. This can be carried out by colliding a laser wakefield accelerated electron beam with a counter-propagating high intensity laser pulse [75, 76]. This method of gamma ray production was first demonstrated on SLAC where the LINAC accelerated electron beam was collided with a laser pulse [21, 77]. More recently, successful implementation of inverse Compton/Thomson scattering on all-optical sources has also achieved multi-MeV gamma rays from the scattering interaction [8, 10, 11].

As with the radiation produced by betatron oscillations described in the previous section, the inverse Compton scattering photons are produced because a relativistic electron beam interacts with an oscillating electric field. In this case, however, the very intense electric fields are generated by the counter-propagating laser rather than quasi-static charge separation as in a plasma wave. This interaction falls into the regime where the wiggler parameter, $K$ is much greater than 1 and so the electrons lose a significant amount of their energy during the interaction and produce a broadband, synchrotron-like spectrum.

Estimating the photon energy capable of being produced by all-optical systems can be done by using common values from LWFA systems in Equations (2.24 - 2.25) [71]. This calculation shows the potential for very high energy radiation production from an optical source while maintaining a high degree of energy tunability by changing the energy of the counter-propagating laser. A visual representation of this equation can be seen in Figure 2.7(a) for various laser intensities and electron beam energies.
\[ \hbar \omega_c \, [\text{MeV}] = 3.18 \times 10^{-6} \gamma^2 \sqrt{I \text{[10}^{18} \text{ W/cm}^2]} \] (2.24)

\[ \hbar \omega_c \, [\text{MeV}] = 3.18 \times 10^{-6} (1500)^2 \sqrt{1 \times 10^3 \text{[10}^{18} \text{ W/cm}^2]} \approx 225 \text{ MeV} \] (2.25)

With today’s laser technology capable of producing focused intensities with an \( a_0 \) exceeding 20 and relativistic electron beams with gamma factors well over 1000, inverse Compton scattering is highly nonlinear and pushes the envelope of classical theory [78]. Although the laser vector potential is Lorentz invariant, the strength of electric field (or \( K \) parameter) is not invariant, which means the energy of the electron beam can have a significant effect on whether or not the interaction is quantum electrodynamically strong. Hartemann and Kerman [79] solve the Dirac-Lorentz equation [80] for an inverse Compton scattering scenario in the classical regime. They have developed equations for the transverse and longitudinal electron motion as well as the backscattered spectrum that agree with the Lorentz equation for small \( a_0 \). For the parameters relevant to the work in this thesis, the electron-laser interaction is strong enough that it is necessary to consider semi-classical corrections to the equations that describe inverse Compton scattering.

Unlike the oscillations in a LWFA bubble, the inverse Compton scattering interaction can be strong enough for radiation damping to affect the velocity of the electrons. This is an effect known as a radiation reaction in which the oscillating particles feel a recoil force from emitting radiation [81]. The extent to which these effects are relevant are determined by the relativistic and gauge-invariant parameter, \( \chi_0 \), which compares the electromagnetic field that exerts the Lorentz force on the electron to the critical electromagnetic field of quantum electrodynamics (QED). The critical field is the same as the Schwinger limit and is defined as a field capa-
ble of separating a virtual electron-positron pair providing an energy that exceeds
the electron rest mass energy over an acceleration length as small as the Compton
wavelength, \( \lambda_C = \frac{h}{m_e c} \approx 3.9 \times 10^{-11} \) cm \([82]\). The following equations show the
\( \chi_0 \) parameter \([78]\) and the critical field \([83]\), where \( E_0 \) is the amplitude of the laser
electric field.

\[
\chi_0 = \frac{2\hbar}{m_e c^2} \omega_0 \gamma_0 a_0 \sim \frac{\gamma E_0}{E_{cr}} \quad (2.26)
\]

\[
E_{cr} = \frac{m_e^2 c^3}{e \hbar} = 1.3 \times 10^{18} \text{ V/m} \quad (2.27)
\]

One of the corrections that needs to be applied in the QED model relates to the
emitted radiation. The classical description of an oscillating charged particle fails
in this case, as the radiated energy of a single photon can be greater than that of
the particle. However, in the QED approach, a term is added to the equation of
motion to approximately account for the overestimation of the radiated power. This
correction \( g(\chi_0) \) can be added to the expression for the radiation reaction force as
seen below \([78, 79, 84]\).

\[
\frac{d(\gamma \beta)}{dt} = -\frac{e}{m_e c} (E + \beta \times B) - g(\chi_0) \frac{2e^4 \gamma^2}{3m_e^3 c^5} |E + \beta \times B|_{\perp}^2 \quad (2.28)
\]

For this equation, it is simple enough to understand that the right hand side represents
a balance between the force applied by the electromagnetic field (first term) and the
radiation-reaction or damping force (second term). The left hand side is simply
the acceleration of a particle where \( \gamma \) is the Lorentz factor and \( \beta = v_e/c \). The
factor \( g(\chi_0) \) is added to the damping force to vary the effect that the radiation
reaction force has on the motion of the electron \([78]\). This term can be defined by
performing a polynomial fractional fit to the data provided in \([85]\) such that \( g(\chi_0) = \)

30
Figure 2.7: (a) Contour plot of spectral critical energies from scattering interactions of various electron beam energies and laser intensities. (b) Emission spectra for classical theory and various levels of $\chi_0$ displaying the sharp cutoff that appears due to electron energy loss [82].

\[(3.7\chi_0^3 + 31\chi_0^2 + 12\chi_0 + 1)^{-4/9}\]. It is easy to see from this definition that as $\chi_0 \to 0$, $g(\chi_0) \to 1$ and the equation returns to the original form without the correction factor. Figure 2.7(b) shows the physical reality of this factor. As $\chi$ increases, the radiation reaction gets stronger and acts to decelerate the electron, thus lowering the maximum energy photon that can be radiated.

With interaction physics leaving the classical regime, there will be many experiments in the coming years dedicated to measuring phenomena that are described by quantum electrodynamics. In an inverse Compton scattering experimental setup, “straggling” and stochastic spreading of the electron beam are two related QED phenomena that would be of interest [86]. When a relativistic electron beam collides with a counter-propagating laser, the electron will have radiated away much of its energy by the time it reaches the peak field intensity of the laser due to the spatial intensity profile of the laser. When QED becomes relevant, the electron has a finite probability of “straggling” into the peak field of the laser pulse without having radiated much energy [86–88]. This effect manifests itself as an increase in hard gamma ray production at the tail end of the spectrum. As a result of straggling and due to the
probabilistic nature of photon emission from an electron in the QED regime, there should be a significant increase in the energy spread of the electron beam after scattering [86–88]. In the classical limit, a monoenergetic beam of electrons experiencing the same field would radiate equally and thus the beam would also be monoenergetic after the interaction. This transition from the classical to the QED regime may prove rather difficult to measure on a laser wakefield accelerator as it would require a very stable and monoenergetic electron beam.
CHAPTER III

Experimental Methods

The experiments discussed in this thesis were conducted on two similar high intensity laser systems: the Hercules Laser system at the Center for Ultrafast Optical Science (CUOS) at the University of Michigan, Ann Arbor, and the Astra-Gemini Laser System, of the Central Laser Facility (CLF) at the Rutherford Appleton Laboratory in Didcot, UK.

3.1 The Hercules Laser System

The Hercules (High Energy Repetitive CUOS LasEr System) Laser is a solid-state system using sapphire crystals doped with titanium (Ti:Sapphire) as the lasing medium. The laser is capable of delivering 300 TW pulses (10 Joules in 33 fs) at a repetition rate of 0.1 Hz [35]. Hercules is a chirped pulse amplification system that operates at a central wavelength of 800 nm and is composed of 5 different amplification stages as shown in Figure 3.1. The Hercules pulses originate from the oscillator, a Kerr-lens mode-locked cavity producing a 75 MHz train of pulses with nJ-scale energy and a duration of \(~12\) femtoseconds (fs). Exiting the oscillator, the train of pulses are stretched and spectrally shaped by a Dazzler and sampled down to 10 Hz by a Pockels cell and a set of polarizers before receiving an energy boost up to the \(\mu\)J level in the pre-amplification stage. The pre-amp stage along with the regenerative amplification
Figure 3.1: Block diagram of the HERCULES Laser system. The system is divided into three main parts, the front end where the pulse receives 1000× amplification before stretching, the amplification stages where the pulse undergoes several stages of energy gain and the compression stage where the pulse is re-compressed down to the femtosecond duration.

stage and 10 TW amplification stage are all pumped by neodymium-doped yttrium aluminum garnet (Nd:YAG) lasers frequency doubled to 532 nm in order to satisfy the pumping requirements for Ti:Sapphire. Before further amplification can occur, the pulse must be stretched by several orders of magnitude as to not damage optics and crystals during the energy gain. Immediately after pre-amplification, the train of 10 Hz pulses is stretched by a folded Martinez pulse stretcher to 0.5 ns [89]. The train of pulses now have a long enough duration that they can safely travel through the rest of the amplification stages. In the next stage, the regenerative amplifier provides enough energy to bring the pulses up to 15 mJ after 20 round-trip passes in the Ti:Sapphire crystal for that stage. At this point, the 10 Hz laser amplified by the regen can be used for low-power alignment. Since this alignment path still goes through the compressor, the 10 Hz train of pulses are on the femtosecond time-scale allowing for these mJ laser pulses to be used in setting up timing diagnostics.

For full power shots, the regenerative amplified laser is sent into 3 final amplification stages: 10 TW, 30 TW and 100+ TW. The names of the stages roughly indicate
the output power of the laser after each individual amplification stage. During the final amplification stage, the laser is increased to a maximum of 17 J by four flash lamp pumped Nd:Glass lasers before being sent into the compressor chamber. The Hercules compressor consists of two pairs of gold-coated holographic gratings with 1200 lines/mm aligned in a Treacy setup [90]. Exiting the compressor, Hercules delivers a maximum of 9 J in a minimum of 30 femtoseconds to one of two experimental chambers. The two experimental chambers are designed to facilitate either gas or solid target experiments. Much of the solid target work is centered around ion acceleration while the gas target work, which is presented in this thesis, is centered around electron acceleration.

3.2 The Astra-Gemini Laser System

The Astra-Gemini Laser system is part of the Central Laser Facility on the RAL campus in Chilton, England. Similar to Hercules, the Astra-Gemini laser is a Ti:Sapphire CPA system pumped by frequency doubled Nd:YAG and Nd:Glass lasers. The system is split into two laser beamlines with the Astra component capable of supplying up to 0.75 J in 30 fs at 2 Hz on target. The Gemini portion of the laser is a final amplification stage that was later added to the standalone Astra laser. Before the compressor of the Astra laser is a pulse picker that splits the 10 Hz beamline into a 5 Hz beamline for the Astra target area and a 5 Hz beamline for the Gemini amplification system. The Gemini amplification system consists of two separate beamlines each capable of amplifying the incoming 1.5 J up to 25 J using separate 60 J Nd:Glass pump systems. The separate beamlines in the Gemini amplification stage means there are two separate compressors each capable of delivering up to 15 J in 40 fs to the Astra-Gemini target area. The dual 500 TW beamlines on Astra-Gemini along with a very large target chamber provide excellent opportunities to further explore LWFA applications through pump-probe experiments or colliding beams. A
Table 3.1: Comparison of HERCULES and Astra-Gemini by the numbers.

<table>
<thead>
<tr>
<th>Parameter</th>
<th>HERCULES</th>
<th>Astra-Gemini</th>
</tr>
</thead>
<tbody>
<tr>
<td>Central Wavelength</td>
<td>800 nm</td>
<td>800 nm</td>
</tr>
<tr>
<td>Repetition Rate</td>
<td>0.1 Hz</td>
<td>0.05 Hz</td>
</tr>
<tr>
<td>Number of Beamlines</td>
<td>1</td>
<td>2</td>
</tr>
<tr>
<td>Beamline Energy (each)</td>
<td>9 J</td>
<td>15 J</td>
</tr>
<tr>
<td>Minimum Pulse Duration</td>
<td>30 fs</td>
<td>40 fs</td>
</tr>
<tr>
<td>Beamline Max Power (each)</td>
<td>300 TW</td>
<td>500 TW</td>
</tr>
</tbody>
</table>

side-by-side comparison of the two laser systems can be found in Table 3.1.

3.3 Laser Diagnostics and Alignment Methods

In this section the various laser diagnostics and measurement methods that were used on the HERCULES Laser System are described.

3.3.1 Laser Power Monitoring

Monitoring the laser power on each shot is an easy and essential diagnostic for high-powered laser operation as these experiments require sufficient high power to achieve self-focusing and drive plasma waves as described in Section 2.2. In addition to needing high power to properly run experiments, monitoring the level on each shot indicates whether or not the entire system is operating normally. A sudden drop in power typically indicates the failure of one of the many pump lasers that make up the HERCULES Laser System as described in Section 3.1.

The power is monitored on a shot-to-shot basis using a photodiode to measure the leak-through of a dielectric mirror near the compressor chamber. Dielectric mirrors used on HERCULES have >99% reflectivity with the remaining energy refracted through the mirror capable of being detected on the back side. The signal obtained by the photodiode is displayed using an oscilloscope as a voltage. In order to convert this voltage to a given power value, a calibrated power meter is used simultaneously.
to measure the laser power so that the voltage on the oscilloscope can be correlated to laser pulse energy and therefore laser power. This photodiode is therefore a measurement of the laser energy and operates with two main assumptions: the pulse duration does not change significantly on a shot-to-shot basis and that the voltage measured by the photodiode scales linearly with laser energy. Both of these assumptions are verified by performing autocorrelation measurements and calibrating the photodiode to keep the voltage well below saturation.

### 3.3.2 Alignment Methods

Ensuring proper and consistent alignment of the laser into the experimental chamber is necessary for data obtained during different experiments to be compared. Significant variations in the laser beam pointing will change the direction of the electron beam, thus changing the calibration of the electron spectrometer and misaligning many of the diagnostics.

In the gas target experimental chamber on HERCULES, there are two main alignment steps that are required for any experiment. First of which is the pointing alignment – this meant to ensure that the laser is properly entering into the chamber from the compressor the same way each time. This means that the beam is both parallel to the optics table and points directly down the center of the entrance port without skewing to either side. The pointing diagnostic involves measuring the light that leaks through the first dielectric mirror in the experimental chamber as shown in Figure 3.2. The pointing imaging system is fixed and control of the beam entering the chamber comes from a mirror after the compressor. Because the pointing imaging never moves, there is certainty that the laser is entering into the first mirror centered and parallel every time.

The next step in the process is to bring the beam into alignment on target chamber center (TCC). The TCC point in the chamber is the location of the gas target with
Figure 3.2: A schematic of the gas target chamber on HERCULES displaying the primary beam path along with the two main points of alignment: pointing and focal spot cameras.

the beam running parallel down the longitudinal axis of the chamber. The alignment to this point is achieved by inserting small “pick-off” mirrors into the focal path to intercept the beam and point it to the stationary “Focal Spot Camera” as shown in Figure 3.2. In this situation, a stationary camera is used to mark a specific point in space that corresponds to the laser moving down the axis of the chamber parallel to the floor.

3.3.3 Focal Spot Characterization

In addition to proper alignment of the laser pulse in the experiment, the laser pulse duration and focal spot quality must be monitored to have optimum laser performance. The focal spot of the beam at best focus is measured during the alignment to TCC. An 8-bit charge-coupled device (CCD) Watec camera is used to view the
Figure 3.3: Line out of a focal spot on the HERCULES laser with the inset showing an image of the focal spot itself. Due to no electro-optics being used in generating the focal spot, there is a portion of the energy existing in large wings of the pulse.

beam profile through a 10× magnification objective and displayed using Coherent BeamView software. Figure 3.3 shows an example of the laser focus from the f/20 parabola measuring approximately 34 µm. This is slightly larger than the diffraction limited diameter of an f/20 focusing system of 20 µm. Although focal spot size is important in a LWFA experiment to aid in coupling energy into the plasma waves [91], it is possible to successfully perform LWFA experiments without opto-electronic devices such as deformable mirrors (DM) through careful alignment of the optics and use of high quality focusing off-axis parabolic mirrors. It is not only necessary to monitor the quality of the beam profile at best focus before each experiment but it is also beneficial to scan to either side of best-focus and measure any astigmatic qualities of the beam. For this reason, the 10× objective of the focal spot imaging system is mounted on a translation stage capable of scanning through best-focus or shifting to a new point in space to account for any new optics in the beam path.
3.3.4 Pulse Duration Measurements

As was discussed in Section 2.2.4, the regime for LWFA depends largely upon the pulse duration of the laser pulse being used to drive the plasma waves, and for this reason it is essential to monitor the duration of the laser pulse. On the HERCULES laser, the pulse duration is typically measured from the low-power regenerative amplifier stage and a second-order autocorrelator. A second-order autocorrelator is a method of measuring the pulse duration of an ultrashort laser by crossing two halves of the pulse in a crystal to generate second harmonic light with a spatial distribution proportional to the duration of the input pulse [92].

Laser pulses like those produced by HERCULES are the shortest events created by human technology and thus it is impossible to resolve their temporal structure with any form of standard detector or camera. Due to this, it is only possible to measure these sub-picosecond pulses with the pulse itself. In order to achieve this, it is necessary to split the input pulse into two using a beam splitter and sending one of the two pulses into a delay stage as depicted in Figure 3.4(a). The two pulses are then focused into a non-linear crystal that will produce second harmonic light as a result of the input beams interacting in the crystal. It is possible to calibrate this second harmonic light to determine a pulse duration by delaying one beamline with
respect to another. Changing the optical path length affects the overlap of the pulses in the second harmonic generation (SHG) crystal and therefore changes the spatial distribution of the second harmonic light on the detector [92]. By taking note of the spatial shift that accompanies the temporal delay change and assuming a pulse shape, we can calculate the duration of the laser pulse. This is sufficient for the majority of LWFA experiments on the HERCULES laser and an example of an autocorrelation scan can be seen in Figure 3.4(b) showing pulse durations as low as 33 fs.

Although second-order autocorrelators are very useful in determining the duration of a pulse, it is often very important to also know the shape of the pulse or measurements of any features that exist in the pre-pulse or post-pulse. Despite the intensity autocorrelator being a one dimensional problem, several unique solutions exist, meaning multiple different input pulses can generate the same autocorrelation trace. Because second-order autocorrelators are symmetric, they cannot provide any information about the front or back of the pulse and tend to wash out any asymmetric structure that may be present in the pulse envelope. Therefore, it is often necessary to use a different method of pulse duration measurement such as a third-order autocorrelator, which is an asymmetrical autocorrelation measurement capable of providing information about the pre-pulse and post-pulse structure, or a Frequency-Resolved Optical Grating (FROG) which provides both temporal and frequency resolved measurements of the pulse. Rather than measuring signal energy vs. delay as is done with an autocorrelator, the FROG measures the signal spectrum vs. delay [94]. With this method, a FROG is capable of delivering a more well-rounded picture of the input pulse by revealing both the phase and intensity vs. time.

3.4 Gas Targets

In this section the various gas targets used over the course of the research for this thesis are discussed.
In order to create a low density plasma target for LWFA, it is necessary to sustain a plasma channel where the laser can remain guided long enough to trap and accelerate electrons in the plasma waves. The simplest way to do this is to sustain a column of gas as a target for the focusing laser. In the first successful LWFA experiments on HERCULES, supersonic gas jets a few millimeters in length and narrow capillaries provided this steady-state column of gas [96, 97]. Leaving the gas jet open for a few milliseconds allows the gas column to reach a steady-state with the surrounding vacuum environment with respect to the femtosecond laser. Despite the gas jet being very robust, simple, and effective, it does not provide the most stable target for creating a consistent plasma channel free of shocks and turbulence and it is necessary to expel a significant amount of gas into vacuum on each shot to reach the required...
densities. Using a capillary on the other hand allows for very successful and stable guiding of laser pulses over long distances, however it is not without its own challenges. Using a capillary requires very precise alignment, potentially several times over during an experiment, as it is not uncommon for the laser pointing to drift slightly. One must also be very careful not to focus the laser on the opening of the capillary as the thin plastic can be destroyed by a single shot.

On HERCULES, various designs of 3D printed gas cells are used instead of jets or capillaries, and have resulted in significant improvement of the electron beam pointing stability, divergence, energy spread, and maximum energy [95]. Gas cells allow for operation at significantly lower backing pressures than the jets; and due to the relative containment of the gas, a more stable gas environment that is linearly proportional to the backing pressure is created. Figure 3.5 shows an example of the change in electron beam quality over several different gas cell designs. The cells themselves can be seen in Figure 3.6. The first iteration of gas cells involved a simple, single-stage shown in Figures 3.6(c) and (d) and provided a more consistent and stable gas environment. Additional iterations of the gas cell design involved adding a second
stage and variable length control. Figures 3.6(a) and (b) show an example of a two-stage, variable length gas cell. The two-stages were implemented in an attempt to control electron injection and the variable length is made possible by the $45^\circ$ slope at the back of the cell allowing vertical movement to change the plasma length. The extra height of the variable length gas cell was designed for more accurate interferometric measurements. The reference area for the interferometer described in Section 3.5.3 is above the plasma channel, and unlike the single-stage gas cell, the reference now incurs the same background phase shifts, due to glass windows and gas, as the interaction region. With the two-stage gas cell, the only difference between the reference and data regions is the plasma. This means that any phase shift incurred in the data region is due only to the plasma.

### 3.5 Experiment Diagnostics

This section describes the various diagnostics used during experiments to monitor the laser-plasma interaction, the electron beam energy and profile, and the x-ray signal.

#### 3.5.1 Electron Spectrometer

The electron spectrometer (eSpec) is one of the most important tools in laser wakefield acceleration experiments because it provides information about the maximum energy of the electron beam, the energy spread, and the beam divergence. The eSpec set up in the gas target chamber of the HERCULES laser involves a dipole magnet and a LANEX scintillation screen. As the electron beam is ejected from the back of the plasma, it enters into a 0.8 Tesla dipole magnet with a separation of 2 cm between the plates and a length of 15 cm. The magnetic field lines oriented vertically normal to the direction of propagation in order to direct the electrons in the plane parallel to the chamber floor. The magnetic spectrometer was designed so that the
Figure 3.7: Calibration curve for the electron spectrometer used to translate electron energy to pixel value on the imaging CCD.

deflecting radii are considerably larger than the plate separation, thus providing a homogeneous field and minimizing the effect of the fringe fields. When the electrons experience the strong magnetic field, they bend in a circular arc with a radius equal to the relativistic Larmor radius shown in Equation (3.1), where $e$ is the elementary electron charge, and $B$ is the strength of the magnetic field.

$$r_L = \frac{\gamma m_e c}{eB}$$

(3.1)

When the electrons leave the dipole magnet, they travel straight into the LANEX scintillation screen with the high energy electrons having only altered their course slightly due to the large radius of curvature and the low energy electrons having
been turned significantly. The LANEX fluoresces at about 545 nm with an intensity proportional to the amount of charge that struck the screen. An example of various electron spectra obtained from this diagnostic can be seen in Figure 3.5.

To calibrate the electron spectrometer, particle tracking and geometrical measurements of the chamber were made to transfer the trajectory of various electron energies. The calibration of the spectrometer is highly dependent upon the location and angle of the LANEX screen. Figure 3.7 shows the electron energy translated to CCD pixel number so that energy lines could be plotted on the CCD images of the electron beam. It is evident from the figure that the resolution at lower energy is much better than at high energy. This is partially due to the physics of relativistic Larmor gyration and partially due to the geometric constraints of the experimental chamber.

3.5.2 Electron Profile

The electron profile or “eProfile” diagnostic is another LANEX screen set up to image the accelerated electron beam. For the experiments discussed in this thesis, the LANEX was set up on-axis, perpendicular to the electron beam trajectory to measure the profile. The LANEX was shielded from the front using aluminum and imaged by a CCD camera viewing the back. This is an important diagnostic because it provides information about the structure of the electron beam and surrounding area. In conjunction with the eSpec, the eProfile serves to give a more complete picture of the electron beam structure as it exits the gas cell.

It has been used in several experiments on HERCULES including measurements of an accelerated ring structure that appeared under certain conditions [59] and divergence measurements of electrons that aren’t trapped in the wake as described in Chapter IV.
3.5.3 Michelson Interferometer

One of the most important diagnostics used on most experiments is the Michelson interferometer. This diagnostic not only shows the size and shape of the plasma channel but also provides density information about the plasma formed by the laser-gas interaction. A Michelson interferometer works by overlapping or interfering the interaction region of the plasma with a reference region of no plasma. The density of the plasma can be obtained by measuring the relative phase accumulation gained by the light passing through the plasma channel. The total phase accumulated by light ($\Delta \phi$) of wavelength $\lambda_0$, passing through a medium of length $L$, with a change in index of refraction given by $\Delta \eta = \Delta \left(1 - \frac{\omega_{pe}^2}{\omega^2}\right)^{1/2}$, is given by Equation 3.2 [98].

\[
\Delta \phi = \frac{2\pi L}{\lambda_0} \Delta \eta
\]  

(3.2)

On the HERCULES gas target chamber, a 2 $\mu$m thick nitrocellulose pellicle is used to reflect 4% of the laser as it enters the chamber. The interferometer beam path is made of 2-inch optics and has to be the same length as the main beam line in order to measure the laser-plasma interaction at the moment the laser pulse reaches the back of the gas cell or jet. A single, two-inch beam passes transversely through the gas cell immediately after the laser interaction as shown in Figure 3.8. With a single beam passing into the interferometer, the reference region and interaction region come from the same beam path. After passing transversely through the gas cell, the beam is imaged on to the Michelson interferometer where it is split into two parts, one of which travels to a delay arm and the other of which is reflected from a roof prism. The roof prism flips the image of the gas cell shown in Figure 3.8 so that the reference region can be overlapped with the laser axis. This is a unique and helpful variant of a Michelson interferometer because it only requires one beam
Figure 3.8: Interferometer setup for (a) a single stage gas cell with the reference region above the cell and (b) a two stage gas cell where the reference region is within the cell in order to accumulate the same phase as the interaction region. (c) A standard interferometer image of a single stage gas cell showing the fringe shift down the axis and the analyzed density map of the fringe shift region.

path before the interferometer itself, and in the case of the two stage gas cell shown in Figure 3.8(b), the reference region acquires all the same phase as the interaction region except for the plasma. With the single stage gas cell shown in Figure 3.8(a), it is necessary to account for the fact that the interaction portion accumulates phase from the glass gas cell windows and the reference portion does not.

The output of this measurement is shown in Figure 3.8(c) with the fringe shift due to the phase accumulation of the plasma channel. Below the output of the interferometer diagnostic is an example of the image analyzed to show the density map distribution of the plasma channel. To extract the density from the 2D phase image, an Abel inversion algorithm was written in MATLAB that relates the phase map, $\Delta \phi(x, y)$ to the plasma density using the following equation [99, 100]:

$$
\Delta \phi(x, r) = \frac{2\pi}{\lambda} \left( \sqrt{1 - \frac{n_e(x, r)}{n_c}} - 1 \right) dl \approx \frac{\pi}{\lambda} \frac{n_e(x, r)}{n_c} dl \quad (3.3)
$$

The interferometer analysis starts to break down when plasma densities begin to exceed $5 \times 10^{19} \text{ cm}^{-3}$ as discontinuous phase jumps appear in the Fourier-transformed
Figure 3.9: (a) Quantum efficiency curves for a variety of ANDOR cameras; the highlighted green curve shows the camera used in gas target experiments. (b) The conversion between photoelectrons and x-ray energy for the ANDOR camera used in single photon counting.

phase map. Operating with densities exceeding $1 \times 10^{19}$ cm$^{-3}$ is rare and typically only done when trying to maximize x-ray flux without care for the quality of the electron beam.

3.5.4 X-Ray Deep-Depletion CCD

Many of the experiments on the HERCULES gas target chamber include measuring x-rays from betatron oscillations. The primary diagnostic used for x-ray detection is an ANDOR iKon-M SO deep-depletion CCD capable of measuring x-rays between roughly 0.10 keV and 30 keV. The quantum efficiency curve across this energy range can be seen in Figure 3.9(a). This camera model has a CCD chip with $1024 \times 1024$ pixels, each of which are $13 \ \mu m \times 13 \ \mu m$, which means the chip is approximately $13.3 \ mm \times 13.3 \ mm$.

This ANDOR camera uses a silicon-based CCD chip which produces a number of photoelectrons proportional to the energy of the incident photon as shown in Figure 3.9(b). The accumulated charge on each pixel is converted to a “counts” value that can be turned back into photon energy if the correct conversion is known. Although the CCD is designed to detect x-ray energies, the camera is also sensitive to laser light and therefore needs to be protected by filters. Typically a thin aluminum or beryllium filter is used to shield the camera from laser light without absorbing too
<table>
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<th>λ [Å]</th>
<th>θ_i [degrees]</th>
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</tr>
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<td>2.50 (Ti)</td>
<td>21.91</td>
<td>1</td>
</tr>
<tr>
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<td>1</td>
</tr>
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<td>19.84</td>
<td>7.95 (Al)</td>
<td>53.27</td>
<td>2</td>
</tr>
</tbody>
</table>

Table 3.2: Various crystals and the Bragg diffraction variables associated with each that were used during x-ray measurements on **HERCULES**.

3.5.5 Crystal Spectroscopy

In addition to the ANDOR CCD camera, crystal spectroscopy was used on various experiments in an attempt to complement the results of the single photon spectroscopy from the ANDOR camera. Three different types of crystals were used in the experiments: highly-oriented pyrolytic graphite (HOPG), highly-annealed pyrolytic graphite (HAPG) and pentaerythritol (PET). The different crystals were used for their different reflection coefficients and angles of refraction.

The limiting factor of the crystals used in x-ray experiments on **HERCULES** is measurable bandwidth. With crystal spectroscopy, the bandwidth of a spectrum capable of being reflected is depending upon the incident angle of the x-rays. Due to the crystals size (∼ 5 cm across) and the low divergence of the betatron beam, there is little difference in the incident angle of the x-rays across the surface of the crystals. The half angle divergence of a betatron x-ray beam is approximately 2 degrees. With an incident angle of θ_i ± 2°, the resulting bandwidth of the reflected x-rays will only cover ∼ 1 keV if the whole beam is reflected. Although the bandwidth is small, the resolution of this detection method is much higher and makes it possible to measure finer structures in an x-ray spectrum that may be lost in single photon counting.

Ultimately, the single photon counting x-ray camera was used as the primary...
diagnostic for x-ray detection in the work presented in this thesis for a number of reasons. Due to space restriction of the vacuum chambers, it was only possible to use an image plate to detect x-rays diffracted from the crystals when an ideal setup would utilize a spectroscopy camera. The spectroscopy camera would allow measurements to be made without ever breaking vacuum as opposed to the use of image plates that require development after accumulating several shots. To create an entire spectrum from crystal diffraction, it is necessary to rotate the crystal several degrees to change the angle of incidence of the x-ray beam and properly stitch together the signal. Due to the changing flux of the x-ray beam and image plate development, it was not possible to measure anything beyond the 200-400 eV window allowed by the various crystals listed in Table 3.2.

3.6 Single Photon Counting

One of the primary reasons the ANDOR iKon M camera was chosen is because it is possible to recreate the incident x-ray spectrum from the output of the camera under the proper conditions. The method used for generating a spectrum from an x-ray CCD image is known as single photon counting and ideally works in a regime when any one of the 1,048,576 pixels on the chip is struck by one, and only one, photon. As was discussed in the previous section, there is a linear relationship between the incident photon energy and the “counts” output of the chip on the camera. Using data provided by the camera manufacturer, a conversion of roughly 15.25 eV/count (above background) can be used to determine the energy of the photon that deposited its energy in a particular pixel.

3.6.1 Double Hit Probability

Although single photon counting is a relatively easy way to obtain a shot-by-shot x-ray spectrum, the process is not without its challenges. The primary consideration
Figure 3.10: (a) Low occupancy Monte Carlo simulation (8.2% of pixels hit) showing good agreement between the theoretical spectrum and the spectrum calculated by single photon counting; the predicted error agrees with the actual double hits. (b) A high occupancy Monte Carlo simulation (30.4% of pixels hit) showing that the reconstructed spectrum begins to fail by overestimating the high energy end of the spectrum due to double hits counting as a single hit. (c) Using the predicted error, the calculated spectrum is adjusted to match the theoretical spectrum.
in single photon counting is limiting the flux to the x-ray camera. This analysis method only works under the assumption that each pixel is hit by only one photon. In reality, it is very difficult to consistently operate in a regime where there is enough flux to produce a statistically significant single shot spectrum but not too much that double hits exist. Determining the probability of a double hit for a given flux can be done by treating the x-ray hits on the camera as a Poisson distribution. Equation 3.4 shows the probability, \( P(n) \), of \( 'n' \) events occurring per interval with \( \lambda \) as the average number of events per interval.

\[
P(n) = \frac{\lambda^n e^{-\lambda}}{n!}
\]  

(3.4)

For the case of x-rays hitting a \( 1024 \times 1024 \) CCD chip, \( n \) will be 2 or 1 hit per pixel and \( \lambda \) is the number of pixels hit by an x-ray divided by the total number of pixels, or the chip occupancy. By comparing the probability of a double hit to the probability of a single hit, it is possible to relate the double hit probability (\( DHP \)) to the occupancy of the chip as was done in Equation 3.5, showing that the probability of a double hit event occurring is equal to half of the occupancy. This means that if 10% of the pixels were hit by a photon, the probability of one of those pixels having been struck by two photons is only 5%.

\[
DHP = \frac{P(2)}{P(1)} = \frac{\left[ \frac{\lambda^2 e^{-\lambda}}{2!} \right]}{\left[ \frac{\lambda^1 e^{-\lambda}}{1!} \right]} = \frac{\lambda}{2}
\]  

(3.5)

It is also possible for higher order hits (3 or 4 per pixel) to occur as well; however, due to the \( \lambda^n \) term in the numerator and the \( n! \) term in the denominator, these probabilities are very small at any acceptable occupancy and have therefore been
ignored in all analyses. To investigate the effect of double hits on a recreated spectrum a Monte Carlo code was written to simulate an x-ray burst on the camera with various levels of occupancy. The code also attempted to correct for the error accrued by double hits registering as single hit by predicting the double hits spectrum. The results of this analysis are shown in Figure 3.10. Figure 3.10(a) shows the results of a low occupancy simulation; with only 8% of the pixels filled, only 4% of the photon hits are double hits. The single photon counting algorithm is able to accurately reproduce the theoretical spectrum and the estimated error (i.e. double hits spectrum) agrees with the actual double hits across each energy bin. In these calculations, the double hits spectrum (DHS) was calculated by performing a convolution of the calculated spectrum with itself. This convolution:

\[
DHS = \frac{\lambda}{2} \sum_{i} S(i) S(i - j)
\]

serves to sum the signal of each energy bin \((S(i))\) with every other energy bin to calculate the probability of a double hit event occurring across the spectrum. As the number of photons impinging on the CCD chip is increased and the occupancy rises towards 30%, Figure 3.10(b) shows that a 15% probability of double hits is enough to distort the calculated spectrum. Due to several lower or mid-energy photons hitting the same pixel and being recorded as a single photon, the spectrum is over-estimated at high energy end and lacking in hits between 0 and 15 keV. The estimated error however still agrees reasonably well with the actual double hits and can therefore be used to correct the miscalculated spectrum. Figure 3.10(c) shows that the adjusted spectrum, which was calculated using the estimated error, agrees quite well with the theoretical spectrum. This Monte Carlo simulation clearly shows that it is possible to correct spectra up to \(\approx 30\%\) occupancy using the estimated error. This greatly reduces the strict occupancy limit on this detector and allows more data to
Figure 3.11: (a) A portion of the x-ray CCD camera showing various photon hits. (b) One of the hits enlarged to show energy deposited into several pixels.

3.6.2 Charge Spreading

Another challenging aspect of using this x-ray CCD camera is the issue of energy spreading. When an x-ray impinges on the silicon detector of this camera, it rarely deposits all of its energy into a single pixel due to multiple collisions within the silicon depletion layer. Figure 3.11 shows that even when x-ray flux is too low for double or neighboring hits to be statistically possible, the hits on the CCD are spread across several pixels rather than being contained to just one. If the image in Figure 3.11(a) were to be read in without any corrections applied, the resulting spectrum would be distorted with significantly more low energy photons than in reality and the high energy photons would be registered a few keV lower than in reality. The figure also shows the inconsistency of energy deposition across the CCD; some of the hits spread out over 3 pixels and some spread out over 5 or more pixels.

There are a few ways to remedy this issue, one of which is to only accept hits where the charge didn’t spread across multiple pixels. This has the benefit of creating a
Figure 3.12: Single photon calibration of the x-ray CCD camera using an Fe-55 radiation source. The calibration shows the extent of the charge spreading (red line) along with the algorithm’s ability to correct for the charge spreading (black line).

spectrum from true hits only but it requires a very high number of pixels to be discarded and therefore produces noisy spectra. In addition, the energy spreading that occurs is likely more common among high energy photons and so this method would distort the shape of the spectrum. Another simple method is to bin the image into $2 \times 2$ or $3 \times 3$ bins. This has the added benefit of summing together the divided energy back into a single hit for proper recording, however, it will significantly lower the resolution of the camera thus lowering acceptable flux levels on the camera. Binning a shot that is too densely populated can result in forcing false double hits if two shots land within the same binning radius. The best way to account for the energy spreading is to “anti-diffuse” the charge by solving:

$$\frac{\partial Q}{\partial x} = \kappa \nabla^2 Q$$  \hspace{1cm} (3.6)

where $Q$ is the charge and $\kappa$ is an arbitrary constant. Integrating $\nabla \cdot (\nabla Q)$ to get
∇Q provides the gradients across the chip which indicates the sources from where the charge spread. From this point it is possible to take all the charge and deposit it up the gradients in order to sum diffused charge back into the primary hit. The algorithm steps through this charge-summing process at least 3 times to account for all the energy leakage. Figure 3.12 shows the ability of this special binning method to recollect spread charge back into the primary pixel from which it originated.
CHAPTER IV

Annular Distribution of Un-trapped Electrons

4.1 Introduction

Since laser wakefield acceleration was first proposed [41], experiments have been primarily centered around studying the electrons that get trapped and accelerated [2–4]. However, there is much more happening in the wakefield bubble beyond the trapped electrons. As research progresses for LWFA as a competing light source, efficiency becomes more important. With much of the fundamental physics of LWFA having been measured and modeled, it is important to turn towards optimization of the process. One way to gauge the efficiency of LWFA is to determine the fraction of trapped electrons. It is possible to calculate this numerically based on the parameters of the laser and plasma [63] but a diagnostic is required to measure this in real time.

Measurements of un-trapped electrons with high transverse momentum forming an annular distribution around the linear beam are presented in this chapter. These electrons start out following the same path as electrons that get trapped. However, they do not have the proper momentum matching required to become trapped and therefore only receive a small forward momentum kick. Recent work has shown that this annular beam can have an energy of a few MeV and 100’s of picoCoulombs of charge [101]. Since using the main beam to generate high-energy photons is one of the goals of LWFA research, it is important to understand the physics of these annular
electrons as they can carry away a significant amount of energy from the laser that may otherwise be put into acceleration.

While these electrons are likely not going to be used as the driver of a light source, they may have applications in single-shot imaging or dosimetry – due to their ability to induce fluorescence over a large area [101] – or they may be a useful diagnostic for LWFA experiments. Measuring the divergence and charge of the electron beam on each shot will reveal information about the wakefield structure or trapping potential as the orbit of these electrons is very closely related to the self-trapping mechanism. The ratio between the trapped and un-trapped electrons is not currently measured in the gas target chamber of Hercules, but on-axis imaging has revealed that un-trapped electrons form a high-divergence ring around the beam. Detecting both the ring electrons and the beam may help to optimize experimental goals by making adjustments to increase or decrease the amount of charge that becomes trapped. Implementing this on Hercules would be a first step towards improving the efficiency of the system. The work presented in this chapter, however, is only an initial measurement of these rings; further discussion about the potential use of this measurement can be found in the Future Work section of Chapter VIII.

4.2 Methods

This experiment was carried out on the Hercules laser in the gas target chamber. A schematic of the experimental setup can be seen in Figure 4.1. Laser wakefield acceleration experiments were performed with an $f/20$, off-axis parabola and a 5 mm gas jet using helium as the working gas. The electron beam quality was first optimized using a 0.8 T magnetic spectrometer to produce consistent, low-divergence beams. Measurements of the electron beam profile were made on axis by placing a LANEX screen 7.5 cm from the exit of the gas jet. The LANEX was shielded with aluminum foil and imaged from the back with an 8-bit CCD camera. The setup was
Figure 4.1: Experimental setup of the on-axis electron beam profile measurements.

designed such that moving the dipole magnet out of the beam path moved the on-axis LANEX into place. To obtain better imaging of the electron beam, one shot was taken with an image plate placed in the same position as the LANEX. The difference in the electron profiles when using image plates and LANEX can be seen in Figures 4.2(a) and (b)–(d) respectively.

Analysis of the ring size was performed by tracing the ring profile on the LANEX images as shown in Figure 4.2(d)–(e) and obtaining the vertical and horizontal size of the trace. The vertical and horizontal divergences were averaged together to get an average divergence of the whole ring. The vertical error bars in Figure 4.2(a) are the result of this averaging. While this analysis method is a bit subjective, the low contrast of the images and incomplete rings makes it difficult to quantitatively define the location of the ring at all points.

To confirm the source of these electrons, 3-dimensional particle-in-cell (PIC) simulations were carried out in OSIRIS [102]. Two simulations were run at 0.5% critical density ($n_{cr}$) and 0.125% $n_{cr}$ with two particles per cell and a laser $a_0$ of 4. In both cases, the laser spot size was matched to the wakefield; i.e. the lower density simulation had a laser spot size exactly twice as wide as the higher density simulation.
4.3 Results and Discussion

The data taken during this experiment were a series of on-axis LANEX images some of which showed the existence of a large-divergence ring surrounding the central electron beam profile. Figure 4.2 shows on-axis ring profiles of various sizes measured at different plasma densities. The circular nature of this feature is due to the spherical shape of the ion cavity. The ring is formed by un-trapped electrons that receive a forward momentum boost by traversing behind the bubble. This population of electrons can be seen in Figure 2.2 labeled as “high electron density.”

The 3D OSIRIS simulations that were carried out as part of this work confirm that the origin of these rings is from un-trapped electrons traversing the shell of the bubble and receiving a forward momentum kick as they pass behind the first bucket. Figure 4.3(a) shows a slice of the 3D space with several electron trajectories mapped on to the wake structure. The particles that appear to enter straight through the middle of the bubble are actually traveling over or under the bubble as the figure is a projection of 3D space on to 2D. In the frame of reference of the traveling laser
pulse, the electrons originate in front of the pulse at the right side of the window, travel around the ion cavity, and are kicked out the side of the simulation window. The electrons exiting the simulation window nearly perpendicular indicates that they are traveling at close to the phase velocity of the wake; this means that they received a strong forward momentum kick as they passed through the region of high electron density at the back of the bubble. The radial symmetry of this population of particles is also highlighted in Figure 4.3(b), showing electrons with high transverse momentum.

As stated in the Methods section, the simulation was carried out at two different densities to determine if there was a significant effect of density on the ring formation. Figure 4.5 shows the profile of the electron beam exiting the plasma. The ring shape is very clear in both of these simulations with the lower density simulation showing the appearance of two concentric rings. The secondary, inner ring was also seen on some data shots and is due to electrons traveling behind the second bucket as shown in Figure 4.3. Although the simulations showed the appearance of rings under both simulation parameters, there does not appear to be a strong relationship between the ring size or shape and the density.

Analyzing the varying size and shape of the rings from the experimental data reveals a slightly inverse trend between the divergence of the ring and the density of the plasma. Figure 4.4(a) is the result of measuring the average divergence of the on-axis ring at different plasma densities. The LWFA scalings in Section 2.2.5 can help provide an understanding of the relationship between the size of the on-axis ring and LWFA experimental conditions. Equation (4.1) shows the relationship between the bubble size, $r_b$, and the plasma density, $n_e$ [71].

$$r_b = \frac{2c}{\omega_p} \sqrt{a_0} = \frac{2c^{3/2}}{e} \sqrt{\frac{m_e a_0}{n_e}} \quad (4.1)$$
Figure 4.3: (a) Two-dimensional projection of the wake structure with electron trajectories superimposed. The electrons that appear to pass through the middle of the bubble are an artifact of projecting 3D onto 2D. (b) Three-dimensional plot of the wakefield structure with electrons having a high transverse momentum highlighted in blue.
This equation states that as the density of the plasma increases, the ion cavity will decrease for a laser with the same incident \( a_0 \). The result of a smaller bubble is a sharper wakefield and therefore may create a lower divergence ring on-axis as is shown with the experimental data in Figure 4.4(a). Although the equation states that the bubble radius scales as \( n_e^{-1/2} \), this trend does not fit the data. This means that either the ring divergence does not depend on the bubble size or there are other unknown factors that change the scaling of the divergence with density. Additionally, recent work with 2D and 3D PIC simulations outside of those performed for this work showed no strong relationship between ring divergence and plasma density or laser \( a_0 \) [101]. While Kostyukov’s model [103] predicts some variation in the divergence across plasma parameters, the authors attribute the lack of a trend to “the small range of momenta acquired by electrons not injected into the bubble [101].” A portion of Kostyukov’s model from the work done by Yang et al. is plotted with the data in Figure 4.4. The slope does not provide a very good fit to the data from this work so it worth properly analyzing this model for HERCULES parameters in the future. Plots by Yang et al. of the divergence from the PIC simulations can be seen in Figure 4.4(b) & (c). The magnitude of the divergence of the rings from Yang’s simulations agree with the divergence measured in this experiment.

4.4 Conclusions

In summary, on-axis measurements of the electron beam were made, showing the existence of a large divergence ring surrounding the main beam. Three-dimensional simulations were carried out in OSIRIS to confirm that the source of the large-divergence ring is from un-trapped electrons gaining forward momentum by traversing between the wakefield buckets. Although no strong trend appeared between the two simulation densities, the experimental data indicates that the size of the ring is weakly correlated to the density of the plasma. Previous simulation work confirms that the
Figure 4.4: (a) Half angle divergence of the on-axis ring measured against shot density. Each diamond represents a single shot with the vertical error bars representing the difference in size in the horizontal and vertical directions. The error bars for density are due to inaccuracies in the interferometer analysis. The yellow dotted line represents Kostyukov’s model, taken from the yellow dashed line in (c). (b) PIC simulations in 2D and 3D showing the relationship between ring divergence and laser $a_0$. (c) PIC simulations with $a_0 = 3$ showing the relationship between plasma density and ring divergence. The K1 and K2 dashed lines show Kostyukov’s model for 1 and 2 bubbles respectively. Figure (b) and (c) adapted from [101].
Figure 4.5: On-axis ring profiles of the 3D OSIRIS simulations at (a) 0.125% critical density and (b) 0.5% critical density. The hot spots in the ring are also present in the experimental data, but these may also be due to low resolution of the simulation.

trend is not apparent across varying plasma parameters, however, models of the wake structure suggest a relationship does exist [101, 103]. Additional simulation and modeling work needs to be done with HERCULES parameters to properly compare the data set.

It is thought that the trend in the data may come from a decrease in the bubble size with increasing plasma density. Although LWFA scalings indicate that the bubble size is related to the plasma density, it is possible that this relationship does not extend to the divergence of the rings. The divergence of the rings is likely dictated by the ratio of the transverse to longitudinal velocity of these un-trapped electrons which may not depend on the size of the bubble. Experimental and simulation work on this annular electron distribution is relatively new and would benefit greatly from future work. Measuring these rings has the potential to serve as a diagnostic to judge LWFA efficiency and trapping success. It might also provide insight into self-trapping as this phenomenon is already challenging to fully understand.
CHAPTER V

The Effect of Ionization Injection on the Betatron X-ray Spectrum

5.1 Introduction

Advances in laser technology have led to the use of high powered lasers to accelerate electrons to relativistic energies. These tunable “table-top” accelerators not only produce quasi-monoenergetic beams of electrons [2–4] but are also excellent sources of synchrotron radiation [7, 24, 26, 104–109]. In previous studies, it has been shown that these x-rays produced via small betatron oscillations [108] have a very high degree of spatial coherence, which is necessary for phase contrast imaging [110], and can resolve features less than 3 μm [26]. And although the origin of these x-rays is well known, there are a variety of experimental conditions such as laser power, plasma density [111], and plasma species [112] that can affect the production of these x-rays.

When performing experiments with an x-ray source, it is often very useful to have the capability to measure the x-ray spectrum of the source on a shot-to-shot basis because it means certain parameters can also be measured on each shot. One such parameter that is useful in describing radiation from betatron oscillations is the critical energy $\hbar \omega_c$. The critical energy is parameterized by both the oscillation frequency of the electrons in the bubble and the electron energy and describes a point
near the peak of spectral intensity [73],

$$\hbar \omega_c = 3 \hbar \gamma^2 \alpha \beta \omega \beta,$$

(5.1)

where $\gamma$ is the Lorentz factor associated with the emitting electron, $\alpha \beta$ is the wiggler parameter and $\omega \beta$ is the oscillation frequency of the electron in the plasma ‘bubble’.

The x-rays produced will likely vary significantly from facility to facility as the laser conditions and gas delivery dynamics will alter the wakefield environment and therefore affect the x-ray production. On Hercules, there are a variety of one- and two-stage gas cells that are used in conjunction with different gases as described in Section 3.4. The goal of this experiment is to characterize the betatron source from LWFA on Hercules and obtain a better understanding of how the different experimental conditions may affect the resulting spectrum.

In LWFA experiments, helium plasma is typically used as the primary medium to accelerate electrons. Previously, it was found that using helium doped with 2.5% nitrogen leads to a higher amount of charge collected by the wakefield through ionization injection [6, 113]. The ionization injection mechanism that occurs when nitrogen is present allows electrons to be born inside the ion bubble as described in Section 2.2.4. By contrast, with pure helium as the working gas, the injection mechanism is self-trapping, meaning a mechanism relying on evolution of the bubble structure [114, 115]. As the electrons swing back together behind the bubble after the laser pulse has passed, they can become trapped and accelerated into the bubble due to the static electric field caused by the charge separation of the ions and the electrons in the bubble shaped wakefield.

As equation 5.1 shows, the critical frequency (and therefore the energy) is dependent upon the strength parameter, $\alpha \beta$. This parameter describes the strength of the betatron oscillations, and depending on where in the bubble the electrons are born,
they can undergo undulator motion (small amplitude oscillations, $\alpha_\beta \ll 1$) \cite{108} or wiggler motion (large amplitude oscillations, $\alpha_\beta \gg 1$) \cite{24}. For small amplitude oscillations, the electrons will emit mostly monoenergetic photons at the fundamental wavelength given by $\lambda = \lambda_\beta / (2\gamma^2)$ (where $\lambda_\beta$ is the betatron wavelength), but the changing electron energy will result in a broad spectrum of emission. For large amplitude oscillations, the electrons will emit instantaneously synchrotron-like radiation that falls in intensity exponentially after the critical energy \cite{24, 73}. As a result, one may expect that the injection mechanism will influence the initial amplitude of the betatron oscillations and therefore the x-ray spectrum. In this chapter, single-hit x-ray spectroscopy \cite{105} is used to compare experimental measurements of the x-ray spectra from betatron oscillations in a laser wakefield accelerator with and without ionization injection (i.e., by including N$_2$ additive to the helium). Measurements are performed with single-stage gas cells and two-stage gas cells \cite{95}, the latter of which enables decoupling of the ionization injection and acceleration.

5.2 Methods

These experiments were performed at the HERCULES laser facility at the University of Michigan. The experimental configuration is shown in Figure 5.1(a). One to three joules of laser pulse energy was focused by an $f/20$ off-axis parabolic mirror onto the entrance hole of various 3D printed plastic walled gas cells, up to 1 cm in length (Figure 5.1(b) and (c)). Density information from the plasma channel was obtained using a probe beam that passes transversely through the gas cell and into a Michelson interferometer as described in Section 3.5.3. The single-stage gas cell was used in experiments utilizing an average power of $53 \pm 9$ TW and the two-stage gas cell experiment was at an average power of $72 \pm 8$ TW. The full width at half maximum (FWHM) of the focused laser spot is $38 \, \mu m$, which provided an intensity in the range $5.1 \times 10^{18} \, W/cm^2 \leq I \leq 6.4 \times 10^{18} \, W/cm^2$. 

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Figure 5.1: (a) Experimental setup of the interaction region. (b) Single-stage gas cell used in this experiment, 7 mm in length. (c) Two-stage, variable length gas cell used in this experiment. The length is varied by changing the vertical position at which the laser passes through the cell. The skinny blue arrows coming from the bottom of the gas cells show the path of gas flow and the thick red arrows passing horizontally from front to back show the laser path.

Both a single-stage gas cell with one gas input and a two-stage gas cell capable of using two separate gases were used in this experiment. The single-stage gas cell was 7 mm long. The first stage of the two-stage gas cell was 1 mm long and filled with helium or helium doped with 2.5% N\textsubscript{2}. The second stage varies between 5 and 10 mm in length and was filled with pure helium. Accelerated electrons produced during the interaction were sent through a 0.8 Tesla dipole magnet spectrometer and measured on a scintillating LANEX screen while the betatron radiation was allowed to propagate 2.55 meters before being detected by an ANDOR iKon-M camera. Aluminum foil of thickness 15 µm was placed directly in front of the CCD camera to be used as a filter to lower the photon flux to enable single photon counting and to protect the camera.

With the experimental parameters of the HERCULES laser system, the betatron oscillations produce x-rays up to about 20 or 30 keV in energy. Using 3D printed gas cells [95] and helium doped with 2.5% N\textsubscript{2} gas produced monoenergetic, high charge, stable electron beams [6]. The greater electron beam stability led to greater betatron
flux stability and therefore improved single-shot spectrum results.

Single hit spectroscopy with a CCD camera was used to detect the photons and generate a spectrum between 1 and 30 keV. The high detection efficiency of the camera and high brightness of the x-ray beam allowed for the calculation of a spectrum on a shot-to-shot basis as opposed to integrating the signal over several shots. Additional details about the x-ray camera and the single photon counting algorithm are discussed in Sections 3.5.4 and 3.6.

5.3 Results

Results were obtained during two experimental runs. The gas cells that were used were both single-stage, 7 mm gas cell and two-stage 6 – 11 mm variable length gas cells (though all shots were taken at 7 mm). Electron spectra that were produced by the single-and two-stage gas cells were typically broad in energy spread as shown in Figures 5.2(a) and (b). As has been shown before [95, 116], the electron beams produced by the two-stage gas cell had lower divergence and were more monoenergetic compared with the single-stage gas cells. The two-stage gas cell produced more consistent and stable electron beams than the single-stage gas cell but the average energy of the electron beams remained between 80 and 120 MeV for both gas cells with a maximum energy of 500 MeV. Although the single-stage gas cell accelerated more electrons and therefore produced beams with higher charge, this did not result in more x-ray hits on the CCD camera compared with the two-stage gas cell. This is most likely due to the fact that electron beams from the single-stage gas cell had a much larger angular divergence compared to the thin, narrow divergence beams of the two-stage gas cell. The large divergence electron beams tend to produce large divergence x-ray beams and therefore result in fewer hits on the CCD chip.

In both types of gas cells, using pure helium tended to generate more monoenergetic electron spectra. For both types of gas cell, there was a clear linear correlation
Figure 5.2: (a) Typical electron spectra from a 7 mm single-stage gas cell and (b) typical electron spectra from a variable length, two-stage gas cell. Electron spectra are from both pure helium shots and N₂-doped shots because there was no constant discernible difference in the spectra.

Figure 5.3: Correlation of measured photon flux $dN/d\Omega$ with measured electron charge. (a) single-stage gas cell and (b) two-stage gas cell.
between accelerated charge and the total x-ray flux generated, as shown in Figure 5.3.

Having the capability to measure the x-ray critical energy on a shot-to-shot basis can yield information about the laser interaction and electron oscillations. During various experiments with both pure helium and 2.5% N$_2$-doped helium, measurements of the critical energy revealed a difference between the spectra of the two gases. To characterize the spectrum in terms of an effective critical energy\(^1\), an exponential function was fitted to the photon energy distribution \(dN/d(\hbar\omega)\) of the form \(dN/d(\hbar\omega) \propto \exp[-2\hbar\omega/\hbar\omega_c]\) [73] and \(\hbar\omega_c\) was extracted from the fit. Many of the data points were averaged over several shots, so error bars were added by finding the standard error of the mean. The points without error bars were single shots rather than averaged shots.

The raw spectra were calculated using an algorithm that converts CCD charge into individual photon hits while accounting for the possibility of charge spreading and double hit pixels [117]. Raw spectra can also be adjusted for the quantum efficiency of the camera and any filters the x-ray beam passes through, producing a spectrum with a resolution as low as 150 eV (see Section 3.6). Typical raw and corrected spectra for (blue) with and (red) without ionization injection are shown in Figures 5.4(a) and (b). When heavy filtering is used, it can be more beneficial to look at the raw spectra rather than artificially distorting the low energy end of the spectra with the response functions. When obtaining the slope of the photon emission spectrum, however, the corrected spectrum needs to be used.

As shown in Figure 5.5(a), when a single-stage gas cell (of 7 mm in length) was used, the spectra produced by a pure helium plasma resulted in consistently and substantially higher critical energies than when nitrogen was present. Each data point was averaged over 1 – 4 shots depending on how many shots were taken that produced

\(^1\)The actual spectral shape may be more complicated than the standard synchrotron function, because the electrons are simultaneously being accelerated while emitting [111].
Figure 5.4: (a) Example of betatron spectrum calculated from raw data on CCD chip and (b) the same spectrum with quantum efficiency and filter corrections applied. Solid blue lines show an x-ray spectrum generated with ionization injection and dashed red lines without. The straight lines show the linear best fit that was used for effective critical energy calculations.

an acceptable photon flux. Figure 5.5(b) shows results from a similar experiment for a two-stage gas cell where the first stage was 1 mm long and filled with N₂-doped helium and the second stage was filled with pure helium and ranged from 5 – 10 mm in length. The graph shows that the same clear correlation did not exist for a two-stage gas cell and that the difference in effective critical energy was unique to a single-stage cell. This difference can also be seen in the spectra in Figure 5.4(a) with the red curve characteristic of helium in the single-stage gas cell and the blue curve characteristic of the N₂-doped helium in the single-stage gas cell.

Figures 5.5(b) and (d) show the same critical energies as a function of average electron energy simultaneously measured on the electron spectrometer. These figures indicate that the effect of a substantially lower effective critical energy with ionization injection was not due to a difference in the maximum energies achieved by the electrons. Since the variable, $\gamma$ in equation 5.1 is dependent upon the transverse and longitudinal electron momentum, a correlation between the electron energy and critical energy was expected. Expression 5.1 suggests that the difference must therefore be due to a difference in wiggler parameter, through the amplitude of oscillations or the betatron frequency.
Figure 5.5: Average spectra effective critical energy for measurements with pure helium (red squares) and nitrogen-doped helium (blue diamonds) in (a) and (c) single-stage gas cell and (b) and (d) a two-stage gas cell (only first stage contains mixture). (a) and (b) are as a function of density. (c) and (d) are as a function of average electron energy.
5.4 Discussion

In a 7 mm single-stage gas cell, such as the one used in this experiment, there is only one cavity and one gas input within the cell. This means that the laser pulse is propagating through a uniform plasma along the entire 7 mm length of the cell as opposed to the two-stage gas cell where the ion species present in the plasma are different in the first millimeter than in the rest of the cell. This is an important aspect of the plasma accelerator because the ion species of the plasma can define the dominant injection mechanism. The injection mechanism describes how electrons become trapped in the wakefield bubble before acceleration occurs. The two different types of injection mechanisms that are present in these experiments are ionization injection and self injection. The injection mechanism is responsible for the differences in effective critical energy seen between a pure helium plasma and a plasma with nitrogen.

When pure helium is used in the cell, the main mechanism for injection of electrons into the bubble is self injection. As the electrons orbit around the ion bubbles in the plasma wave set up by the ponderomotive force of the laser, there is a finite probability that they will become trapped if they receive sufficient forward momentum to match the phase velocity of the bubble [2–4]. When these electrons are pulled into the bubble, they are typically pulled in from far off-axis and therefore undergo large transverse oscillations in the wiggler regime, which produce the synchrotron-like radiation detected by the CCD camera.

For the case where nitrogen is present in the gas, the dominant injection mechanism is ionization injection. In this phenomenon, the inner shell electrons of nitrogen are not ionized until the peak of the laser pulse and therefore, they are “born” at rest within the bubble itself and experience energy gain due to the potential difference between the edge and interior of the bubble [6]. Since these electrons are “born” inside the bubble, it may be expected that their initial oscillation amplitude is significantly

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less than the self-trapped electrons that travel along the bubble sheath. In the experimental data for single-stage gas cells, the x-ray effective critical energy was greater on average with pure helium compared to a gas mixture that would cause ionization injection, consistent with this interpretation. The data further indicates that this may not be the full explanation. The difference in effective critical energy would be expected to be observed for the two-stage gas cell as well, where ionization injection is limited to only the first 1 mm stage. This lack of correlation however, may be due to potential prolific self-injection occurring in the second stage in the two-stage injector/accelerator configuration or that x-ray emission is dominated by oscillations occurring in the second stage and are less affected by the composition of the first stage.

5.5 Conclusions

In conclusion, the effect of the use of ionization injection compared with self-injection on the synchrotron-like x-ray spectra generated by betatron oscillations of electrons in a laser wakefield accelerator for single-stage and two-stage gas cells was measured. In the case of a single-stage gas cell an average increase in the effective critical energy of the x-ray spectrum was observed. This could be due to the effect of the injection mechanism on the initial transverse momentum. This trend, however, was not evident in a similar experiment performed with two-stage gas cells and may mean that injection mechanism — and therefore the transverse momentum — is not confined to the first stage of the gas cell. These results offer some evidence that there may be a difference in transverse emittance on injection of electrons into the wake bubble with ionization injection compared with self-injection, but further study is needed.
CHAPTER VI

K-edge Absorption Spectroscopy of Laser Heated Aluminum with a Laser Wakefield Generated X-ray Pulse

6.1 Introduction

X-rays are vital to the research of the plasma physics community and are used in a plethora of ways, including spectroscopy to identify ionic states [118], ultrafast diffraction to measure biological processes [119], and absorption measurements to observe chemical reactions [120]. Absorption of x-rays by the electrons in the K-shell is a common and well documented phenomenon [121] that occurs across a broad energy range depending on the element. In order to study a variety of materials, it is beneficial and necessary to have a bright and broadband x-ray source. Conventional x-ray sources include wiggling electron beams from linear accelerators in an undulator, broadband x-ray radiation from synchrotron accelerators, and laser-heated foils. Due to the ultra-short nature of laser pulses used in laser wakefield acceleration, the resulting electron beam and x-ray bunch also have a duration on the femtosecond time-scale, making the x-rays ideal for probing ultra-fast phenomena. One of the most appealing experiments for a LWFA is a pump-probe design. Because laser wakefield acceleration is conducted on an entirely optical source, it is possible to con-
duct experiments with a very high degree of temporal resolution by splitting the laser pulse into a pump beamline and a probe beamline. Previous pump-probe experiments using optical sources have shown how the ultrafast nature of lasers and acute timing accuracy allow for measurements such as the disassociation of a molecule [120] or appearance of ionization states as a metal is heated [122].

In this chapter, transmission measurements around the aluminum K-edge using betatron x-rays from a laser wakefield accelerator to probe laser-heated aluminum are presented [123]. Data are obtained using a high-energy detection, deep-depletion CCD camera at various time delays between the pump and probe beams, ranging from 50 fs to 400 fs.

6.2 Methods

The experiments reported here were carried out on the HERCULES laser system at the University of Michigan. Upon entering the experimental chamber, the single, 73 ± 3 TW beam is split using a 75/25 beam splitter with 75% of the energy reflected from an f/20 parabolic mirror to drive a laser wakefield accelerator and the remaining 25% focused by an f/3 off-axis parabolic mirror to heat the aluminum target. This results in an intensity of about $5 \times 10^{18}$ W/cm$^2$ for the wakefield arm and $1 \times 10^{16}$ W/cm$^2$ for the heater arm. The gas target for this experiment was a 3 mm gas cell which was selected over a gas jet in order to provide additional stability and consistency to the electron beam [95]. The timing between the two beam paths was controlled using a delay stage on the pump beamline allowing timing control between the two beamlines to as short as 10 fs. The solid target chosen for this experiment was a 4 µm aluminum foil. Aluminum was chosen due to the fact that there is significant x-ray flux around its K-edge for our parameters, and a thickness of 4 µm was chosen to ensure that a significant level of absorption can occur without it being mistaken for noise or fluctuations. Further discussion about the target chosen for this experiment
can be found in Section 8.2 of Chapter VIII.

The pump and probe beam lines were set up perpendicular to one another with the target at a 45 degree angle to both the pump and probe beam paths as shown in Figure 6.1(a). Because of fluctuations in x-ray flux from betatron oscillations, it was necessary to obtain data for a reference spectrum and an aluminum K-edge absorption spectrum in a single shot. In order to accomplish this, the aluminum targets were set up to cover half of the x-rays that hit the camera as shown in Figure 6.1(c). Due to the geometry of this experiment, shooting the aluminum target at the focal point of the $f/3$ heater beam would have resulted in the projection of a small heated spot causing the camera to detect x-rays that passed through both heated and unheated aluminum. Figure 6.1(b) shows how the $f/3$ focal spot was defocused to $\sim 500 \mu m$ in order to ensure the projection of the heated target would overfill the CCD chip.

The aluminum target was placed in a circular, rotatable, plastic holder with 90 holes 3.5 mm in diameter to hold the targets. The aluminum was placed so that it covered exactly half of each 3.5 mm target hole as shown in Figure 6.1(c). The spatial alignment of the beams was performed by centering a 50 $\mu m$ cross hair on each of the beam lines using a high magnification system. Doing so ensured that the spatial overlap of the beam centers were set to within 50 $\mu m$ of one another. To achieve temporal overlap, it was necessary to have both beamlines imaged onto the same high magnification system. A piece of mylar 2 $\mu m$ thick was placed into one of the 3.5 mm holes on the target wheel to be used for temporal alignment. The transmission of the $f/3$ beamline was measured through the mylar and the $f/20$ beamline was reflected off the mylar to achieve beam overlap. The delay arm on the $f/3$ beamline was scanned until strong fringes appeared on the monitor. The fringes existed over roughly one pulse duration allowing the timing of the beams to be confidently set within one 30 fs pulse duration. The $t_0$ timing was set in vacuum at the start of the experiment and adjusted before each shot series; it was not rechecked.
Figure 6.1: (a) The experimental setup showing the orientation of the pump and probe beam lines relative to the solid aluminum target wheel. (b) The overlap of heated and un-heated aluminum with the betatron x-rays showing how the experiment was designed to ensure that heated x-rays would always hit the CCD sensor.
before each new induced delay.

The geometry constraints of this work were the most challenging aspect to balance. Placing the camera far from the source helped to stabilize the x-ray flux on each shot and between shots. The distance allowed the betatron x-ray beam to spread out, thus limiting the x-ray flux density on camera and allowing for more accurate spectral reconstruction. The distance the camera could be moved away from the source was limited by the laboratory setup as well as the foil magnification. The magnification of the foil’s edge on the camera is the ratio between the source to the image and the source to the object. Increasing the distance between the camera and the source increased the magnification and made pointing fluctuations in the x-ray beam to move the foil edge off the camera more likely. To mitigate the increase in magnification for future experiments, it would be helpful to move the target further from the source; unfortunately this is not possible with the current gas target chamber on HERCULES. While this would serve to reduce the magnification and maintain a consistent x-ray flux on the camera, it also means that the heated portion of the aluminum target would need to be increased. This is shown in Figure 6.1(b); it is evident that moving the aluminum target closer to the camera would require a larger focal spot for the $f/3$ to ensure proper overfill. If the focal spot of the heater beam grows, the laser power must also increase or the target thickness must decrease to account for the lower intensity on target and to ensure uniform heating.

The CCD images were analyzed to generate spectra using a single photon counting algorithm that attempted to account for charge spreading and double hits as described in Section 3.6. The x-ray flux on the camera was limited by placing it 2.5 meters away and using 20 $\mu$m of beryllium as a light filter. Only shots that had a fill fraction of less than 30 percent were accepted as this was shown to be the maximum occupancy that gave acceptable data after compensation for double hits. The best way to handle the issue of double hit events was to operate in a regime where the x-ray flux was low.
enough that double hits were statistically improbable. This approach however, can be difficult due to unstable x-ray flux and provided less available data for calculating a spectrum. The technique used for double hits correction is described in Section 3.6.1 and was implemented in the analysis of this work. The charge spreading was corrected for by implementing the anti-diffusion method mentioned in Section 3.6.2.

This method was tested by applying it to an iron-55 source spectrum recorded with the same CCD as used in the experiment. Figure 3.12 shows that it recovers Gaussian peaks around the manganese K-alpha and K-beta lines as well as the fluorescence and escape peaks which could not be achieved with $3 \times 3$ cell charge collection methods, and lends confidence to the efficacy of the method described in Section 3.6.

6.3 Results

Data were obtained using an Andor iKon-L deep-depletion x-ray CCD camera. Figure 6.2(a) shows an example of the half-half images that were obtained for various time delays. The dark half of the square in Figure 6.2(a) is from x-ray absorption in the aluminum and the light half is x-rays passing through vacuum. The figure shows how the aluminum edge may not always align perfectly straight or down the center of the CCD. Due to the previously mentioned geometry constraints of the chamber and the need to lower x-ray flux on the camera by moving it far from the source, the magnification of the edge on the camera had to be set between 50 and 60 times, causing the edge to sometimes move off of the camera.

With each experiment, it was possible to acquire up to roughly 70 data images before needing to replace the aluminum on each of the 3.5 mm plastic holes. Due to several shots begin discarded because of high flux, low flux, or no edge present, the experiment required considerable amount of time to obtain a reasonable data set.

In order to measure a change in the K-edge of aluminum on the femtosecond time scale, each data image, like the one in Figure 6.2(a), was analyzed into two
Figure 6.2: (a) Example CCD image with a region of interest selected over each half that would be used to create a separate reference and data spectrum. Half the camera is darkened due to x-ray absorption by 4 μm aluminum. (b) The reference and cold aluminum absorption spectra that were calculated from each half of (a) plotted with 4 μm of aluminum as a reference.

separate spectra: one reference spectrum with no aluminum in between the source and camera, and one data shot with heated or cold aluminum present between the source and camera. Figure 6.2(a) shows how a region of interest was selected from each half of the image along with the resulting spectra in 6.2(b). The region of interest was kept constant across the two halves of the image to ensure the amount of flux for each side was proportional. It is worth noting however, that between different data shots, the ROI size changed depending on how well the aluminum metal edge split the CCD image. Figure 6.2(b) shows one of the issues incurred while performing single photon analysis with the data. As mentioned, there is a balance in the x-ray flux between having sufficient x-ray hits to see an edge feature on camera and not having too many hits such that accurate reconstruction of the spectrum is impossible. Because it is very difficult to obtain the perfect x-ray flux, the vast majority of the data, has an occupancy near 30% which results in distorted amplitudes of the spectra. This can be seen by the fact that the data in Figure 6.2(b) does not match perfectly with the reference spectrum × 4 μm thick aluminum. The higher flux on the reference
side of the images results in more data lost to double hits or charge-spreading and in reality it is likely that the aluminum spectrum has a more accurate magnitude than the reference × tabulated data. Although there are some corrections that can be applied, it is not a perfect solution and additional work on the algorithm may yield better agreement to theory. Despite the disagreement with theory, all the data was analyzed the same way and the flux levels were not seen to change the location of the K-edge. Section 8.2 discusses possible fixes to stabilizing the x-ray flux in future experiments.

To obtain a measurement of the transmission spectra for each of these timing delays between the pump and probe arms, the reference side of the image was divided by the absorption side of the image. The result of this was a transmission curve for each pump-probe time step to compare to one another, to the cold aluminum case, and to tabulated data. Figure 6.3(a) shows a few of the aluminum transmission curves superimposed with the cold transmission spectrum. To properly show the shift of the K-edge, the transmission curves of the different time steps have been shift down from their original positions to be in-line with the reference curve.

6.4 Discussion

It is evident from Figure 6.3 that heating the aluminum caused the K-edge absorption energy to shift by a several eV to higher energies. Since the aluminum target was heated by an intense laser pulse, it was necessary to confirm that the shift in the edge was not due to any self-emission of the heated aluminum plasma. To do this, a measurement was performed with the heater beam arriving at the aluminum target after the x-rays had probed. The most negative point on Figure 6.3(b) shows that the location of the K-edge matches that of the cold K-edge within reasonable error, thus confirming that the x-ray absorption is different for the hot plasma. The point at -200 fs which appears to be red-shift from the cold K-edge reveals another
Figure 6.3: (a) Transmission spectra of aluminum at various delays of the probe beam in reference to the pump (or heater) beam. The transmission spectra were adjusted down to show the shift relative to the cold and CXRO reference. The position of the original cold edge is shown by the dashed line. Shifting the spectra does not change the location of the K-edge. (b) The shift in the K-edge measured from the transmission spectra at different pump-probe timing delays. The furthest point on the left with negative time delays has the pump arriving after the probe to ensure K-alpha emission did not create false signal. Points without error bars consist of only one acceptable CCD image.

challenge of accurate analysis of the data. If the camera is not sufficiently far away from the x-ray source, it is possible for the profile of the photon beam to be very non-uniform across the face of the chip. Since the energy of the betatron beam is not flat across the profile of the beam [7], it is possible for there to be an uneven distribution of photon energies across the the aluminum and reference sides of the CCD image. The beam profile is strongly visible on the shot at -200 fs and so it is believed that the uneven distribution of photon energies across the chip is creating a false red-shift at this point. Additional work on the single photon algorithm may mitigate this effect. The point at -100 fs appears heated in the data even though the pump laser was timed to arrive after the probe. This is likely due to the inability to measure the exact timing between the x-rays and the pump pulse. As mentioned, timing was set with the two laser beamlines in vacuum, however, with a group delay
Figure 6.4: One dimensional simulations of laser heated aluminum performed in EPOCH. (a) The hot electron temperature plotted as a function of time and penetration depth. (b) The thermal electron temperature plotted as a function of time and penetration depth.

of the probe laser through the plasma on the order of 70 fs, feedback error on the delay stage motor of $\sim 30$ fs, and timing certainty set to within one pulse duration ($\sim 30$ fs), it is probable that the x-rays arrived 100 fs later than the original vacuum timing.

One-dimensional simulations were performed in EPOCH to determine the thermal and hot electron temperatures of the laser-heated foil. The results of the simulation (Figure 6.4) show that when laser heating of the aluminum occurs, the thermal and hot electron temperatures increased on the femtosecond timescale. The simulations indicate that the incident laser could have created a thermal population of 100’s of eV electrons over a region up to 250 nm thick, which indicates that the majority of the target would remain cold. The hot electrons, however, reach several keV and penetrate greater than 1 $\mu$m into the target in less than 100 fs. Several experiments in similar regimes confirm that the pump laser will create hot electrons on the order of several keV [124–126] in the aluminum target. Because this is happening on a timescale of about 100 fs, the aluminum is in a state in which most of the metal is cold but there exists a population of hot electrons caused by the incident laser that is propagating through the target creating an ionization front [127]. The hot
electrons will preferentially create ionization events in the L and M shells of the aluminum whereas the x-ray radiation produced by the laser-aluminum interaction can cause K shell ionization. The creation of holes in the K, L, and M shells of aluminum will increase the electric potential between the nucleus and the remaining electrons, resulting in a blue-shift of the K-edge. At longer time scales, the build up of surrounding plasma would result in ionization potential depression (IPD), also known as continuum lowering [128, 129]. Lowering of the continuum would result in a red-shift of the K-edge as less energy is required for K shell electrons to reach the continuum.

In the results presented here, the blue shift in the K-edge seems to agree with previous work of isochorically heated aluminum [130–132]. The shift is between 40 and 75 eV with the magnitude of the shift increasing at longer timescales. The blue shift that would result from the ionization events caused by the laser and the subsequent propagation of hot electrons is likely causing the blue shift seen in this data. To better understand how the hot electrons propagate through and affect the 4 µm aluminum target, simulations that capture the physics of hot electrons must be performed. These simulations would also serve to indicate how fast the hot electrons travel through the material. Without simulations, it is difficult to say whether the rising magnitude of the shift with time is due to a physical effect or merely within the error of the measurements.

One of the potential missing parts of this explanation is the lack of resonant absorption in the spectral data. If the aluminum is partially ionized in the L and M shells, the K-edge would experience a blue shift; however, an absorption feature below the K-edge should also be present due to the K shell electrons absorbing x-rays and becoming excited to the M shells. In order for the aluminum to experience a blue shift in the K-edge without resonant absorption, the aluminum would have to preferentially ionize in the K and L shells while the M shell remained full. This would
create a blue shift at the K-edge without the K→M resonant absorption feature that should show up below the edge.

6.5 Conclusions

A measurable shift in the K-edge transmission of aluminum on the femtosecond time scale is presented here. The reaction is initiated by heating a 4 \( \mu m \) piece of solid aluminum with an ultrafast, high-intensity laser. The results show a blue shift at the K-edge of the heated aluminum when compared to the cold K-edge transmission. These results agree with previously published work indicating that a blue shift should occur if the heating is rapid and relatively isochoric \([130–132]\). Hydrodynamic simulations indicate that the majority of the plasma remains cold which further confirms that IPD could not yet have an effect for this experiment. Additional simulations that capture the physics of hot electron production are needed to obtain a better understanding of how they might propagate through the material and cause enough ionization events to blue-shift the K-edge. The simulation work would also serve to provide a time-resolved picture of the hot electron population.

Though using a CCD camera for this experiment was very convenient because it allows spectra to be obtained on a single-shot basis and does not require the processing of image plates, charge spreading and double hits limit the resolution capabilities of obtaining a spectrum through single photon counting. For future work, it would be beneficial to use single photon counting in conjunction with crystal spectroscopy to obtain higher resolution transmission curves.
CHAPTER VII

Inverse Compton Scattering

7.1 Introduction

Since the first empirical observation of laser wakefield acceleration in the blowout regime in 2004 [2–4], the technology has been used as an alternate photon source to linear accelerators and synchrotrons. Although laser wakefield accelerators produce synchrotron radiation with femtosecond duration through betatron oscillations, the photon source is limited in peak energy and brightness. Typical photon energies from LWFA are in the range of a few 10’s of keV with a brightness of up to $10^{22}$ photons per second per mrad$^2$ per mm$^2$ per 0.1% bandwidth [26]. Though these x-rays have significant applications [33, 133], the lower peak energies limit the usability of betatron x-rays in high density measurements, nuclear activation and radiography. Over the past decade, laser power has been increasing and electron beam stability from wakefield acceleration has become more controllable, leading to new methods of high energy photon production from laser-plasma sources.

One such method is inverse Compton scattering of a relativistic electron beam, produced by LWFA, with a counter-propagating, high-intensity laser. Compton scattering as a laboratory tool has been a useful technique for decades, [134, 135] however it has only been a few years since technology has progressed to the point that lasers can be used to generate MeV-level gamma rays through inverse Compton scatter-
Inverse Compton scattering on a laser wakefield accelerator opens the door for much higher energy photon production from an all-optical source. It is desirable to study the creation of bright, multi-MeV photons from a laser source because they are significantly smaller and cheaper than conventional linear accelerators or synchrotrons. One of the challenges of using these large devices for real-world applications such as cancer radiotherapy [136, 137], radiography of dense objects [135, 138], isotope identification by nuclear resonant fluorescence [139], and active interrogation for homeland security [140, 141] is that every material being investigated must be brought to one of the few facilities worldwide. As discussed in the motivation section of this thesis, development of an all-optical source opens the door for a greater degree of location flexibility while still having a tunable photon source.

While current scattering technology operates in the linear regime and has a broad range of applications, experiments of nonlinear Compton scattering can serve to provide an empirical foundation to the physics that govern strong-field QED phenomenon.
such as electron-positron pair cascades \cite{142, 143}. In this experiment, a relativistic
electron beam produced by laser wakefield acceleration was collided with a counter-
propagating laser with a peak, focused $a_0$ of 24 \cite{144, 145}. The goal was to measure
the radiation reaction of the electron beam due to the intense oscillation it experienced
from the focused laser. The radiation reaction was to be confirmed by measuring the
change in energy and energy spread of the electron beam after the scattering interac-
tion. Laser technology is now at the point that QED phenomena can be measured in
the laboratory and this work may be a first step in obtaining experimental data to
compare to assumptions made in radiation reaction models.

Figure 7.1 shows a contour plot of the $\chi_0$ values that could be reached given
the laser and electron parameters in this work. The maximum value that might be
expected from this experiment is just over 0.1, which is significant enough to see a
change in the gamma ray spectrum from radiation reaction as shown in Figure 2.7(b).
As was described in Section 2.3.2, the maximum energy of scattered gamma photons
may be as high as 300 MeV if perfect overlap between a 1 GeV electron beam and
the peak $a_0$ of the scattering laser can be achieved.

7.2 Methods

This section details the various techniques employed in calculating a gamma ray
spectrum in this inverse Compton scattering experiment. The experimental setup is
covered as well as the analytic methods for converting the data into reliable spectra.

7.2.1 Experimental Setup

This work was carried out on the Astra-Gemini laser system at Rutherford Ap-
pleton Laboratory in the UK; details about this laser are discussed in Section 3.2.
Both beam lines were used in this work; one of the pulses was collided with a rela-
tivistic electron beam produced by the other laser pulse. The “south” beam line was
**Figure 7.2:** Experimental setup of the Compton scattering experiment showing the $f/2$ scattering beam focused to the back of the gas jet where electrons are accelerated by the $f/40$ wakefield driver.

**Figure 7.3:** Results of a raster scan of the scattering beam focal spot. The grid shows the integrated CsI signal for various alignments of the scattering beam focal spot with respect to vacuum spatial alignment. No data exists for the black areas.
Figure 7.4: Energy deposition curves of a monoenergetic photon beam impinging on a CsI detector array from 47 different Monte Carlo simulations performed in Geant4.

focused to the edge of a 15 mm diameter gas jet using an $f/40$ spherical mirror to a peak intensity of $10^{19}$ W/cm$^2$. This beam drove plasma waves through the gas target produced by the 15 mm nozzle, resulting in the acceleration of electrons up to 1 GeV. The “north” beam line was focused to the opposite edge of the 15 mm nozzle by an $f/2$ off-axis parabola reaching a peak intensity of $7 \times 10^{21}$ W/cm$^2$ to collide head-on with the relativistic electron beam. A schematic of the experimental setup is presented in Figure 7.2. The best overlap of the electron beam with the scattering beam was achieved by performing a raster scan of the scattering beam and measuring the signal in the gamma ray detector as shown in Figure 7.3(a).

The primary electron diagnostic in this work was a magnetic electron spectrometer. The electron beam was deflected off axis to a LANEX screen by a dipole magnet with a total magnetic length of $\int B(x)dx = 0.4$ Tm. The fluorescence of the LANEX
screen was captured using a cooled 16-bit CCD camera; an example spectrum can be seen in Figure 7.3(c). Many of the electron spectra consist of a low-charge, high-energy feature along with a high-charge, low-energy feature as shown in the spectra of Figure 7.11. The low-energy pedestal is likely the result of separate injection events stemming from fluid shocks in the gas flow. The spectrometer was used in this work to look for a decrease in the electron energy that corresponded with scattering events. The decrease in energy would indicate a decelerating effective force applied to the electron beam due to the release of radiation. A more detailed description of the radiation reaction phenomenon is found in Section 2.3.2.

The gamma rays that were produced by the strong wiggling of the electrons in the intense electric field of the counter-propagating laser were measured using a CsI(Tl) crystal array detector. The detector consists of 1551 CsI(Tl) crystals arranged in a $47 \times 33$ lattice as shown in Figure 7.3(b). The crystals are rectangular prisms that are 5 mm square in cross-section and 50 mm in length. The individual crystals are held in place by 1.0 mm thick aluminum spacers that fit together in a matrix pattern inside the housing to support the crystals.

The gamma ray beam was incident on the side of the 9 mm thick steel housing such that the 47 crystals were oriented along the propagation direction and the 33 crystals were oriented in the transverse direction. The light output of the crystal lattice was imaged with a CCD camera so that the penetration depth and vertical divergence of the beam could be measured on each shot as shown in Figure 7.5. The resulting data file is a $1024 \times 1024$ pixelated image of the light output of the CsI scintillator. Due to an aluminum faceplate on the CsI crystals holding them in place, the light output was constrained to a 4 mm diameter circle causing dark regions in between each circle of signal. To convert this to a more usable format, the number of counts in each pixel within each circular region were summed together into one data point. A comparison of the raw and reformatted data can be seen in Figure 7.5(c) &
**Figure 7.5:** (a) CsI scintillator signal without the counter-propagating scattering beam. The signal is due to bremsstrahlung and betatron radiation. (b) Signal with the scattering beam on. The increase is due to gamma ray production through inverse Compton scattering of the electron beam with the counter-propagating f/2 beam. (c) Pixelated image of the CsI detector array obtained by imaging the scintillator with a camera. (d) The processed image with the pixels in each circle averaged together for spectral analysis.

### 7.2.2 Geant4 Simulations

In order to analyze the data obtained from the crystal array, it is imperative to know how gamma rays will interact with CsI at various energies. Several 3D simulations were carried out in Geant4 with various monoenergetic photon beams irradiating a slab of CsI, as shown in Figure 7.4, to generate response curves for all levels of photon energies that would be relevant to this experiment. For this work, it was necessary to carry out 3D simulations due to significant amounts of side scatter and electron cascading that took place within the crystal array. Two dimensions were necessary to account for scattering and energy transfer between neighboring CsI crystals and the third dimension was necessary to properly capture the geometry of the array so that light yield calculated by the simulations could be accurately
Figure 7.6: (a) A “heatmap” of the dark noise background that was calculated for and subtracted from each shot. (b) The seven different bremsstrahlung background curves plotted below an example data signal. The seven curves were subtracted from each data signal in the analysis.

compared to data. The simulations ranged from 0.1 MeV to 500 MeV with finer steps at low energy and larger steps at high energy. The lower limit was chosen because photons less than 0.1 MeV would not contribute any significant signal through the 9 mm thick steel detector housing. The simulations were stopped at 500 MeV as calculations indicate that the interaction would not produce photons this high in energy.

7.2.3 Image Processing

The first step in analyzing the data files before the pixels were summed together, was to perform a background subtraction. In this data, there are two types of background that must be accounted for: the dark noise of the camera that provides a signal level of roughly 100 counts ±2% even when the laser is turned off, and the background bremsstrahlung signal that appears when the scattering beam is turned off but electrons are still produced. This signal is likely the result of stray electrons hitting the spectrometer magnets, shielding or chamber walls. These two types of background signal can be seen in Figure 7.6.
To account for the dark noise of the camera, the data section of each image was cropped out and a background map was created from the remaining part of the image. This was done by performing a linear fit across each column of the remaining image and smoothing the result to generate a $1024 \times 1024$ “heat map” of the background so that areas of slightly higher or lower dark noise were properly accounted for. This dark noise background subtraction was done for each of the data shots and background bremsstrahlung shots.

Despite significant shielding with lead bricks, the capture camera was still susceptible to stray electron and gamma hits causing occasional high peaks of signal across the image. Taking the Laplacian (i.e. taking the second derivative in $x$ and $y$) of the image after background subtraction forced these hard events to stand out due to the high contrast between the events and the neighboring points. The events were selected by identifying outlier pixels that had counts significantly higher or lower than the standard deviation of the row. The outliers were removed and replaced with an average of the surrounding points such that hard events occurring in the signal region were blended with the signal and background events were blended with the background.

Lastly, it was important to account for the bremsstrahlung signal that was blended in with the gamma ray photons produced through the inverse Compton scattering interaction. The level of this background varied from shot to shot as shown in Figure 7.5(a). Since it was impossible to know the exact level of non-inverse Compton gamma photons that contributed to the signal on every shot, this variable background was a source of error in the spectrum calculations. Further discussion of how the background was used in the calculations is covered in Section 7.2.5.
Figure 7.7: (a) Comparison of theoretical Bremsstrahlung signal to actual Bremsstrahlung data obtained during the experiment. The Brem data is normalized to the maximum value and Brem simulated is normalized to the point of first overlap with Brem data. (b) Camera correction factor generated by dividing the theoretical signal by the actual signal.

7.2.4 Correction Factor

During the experiment, calibration shots were performed by colliding the electron beam with a 9 mm thick piece of lead to measure the CsI signal resulting from a bremsstrahlung interaction. By comparing this signal to a simulated bremsstrahlung signal, it was possible to confirm the reliability of the detector and determine if the CsI detector system responded properly to the gamma beam. The signal comparison between the data and simulation can be seen in Figure 7.7(a).

The theoretical bremsstrahlung signal was generated in Geant4 by colliding a typical electron spectrum from the experiment with a 9 mm thick piece of lead as was done in the experiment. The resulting gamma ray spectrum interacted with a simulated detector to generate CsI signal that could be compared with experimental signal. This comparison required the assumption that the detector response to bremsstrahlung radiation was very similar in shape for slightly varying electron spectra; this was verified through simulation and the experimental data. The difference
between the actual signal and the theoretical signal indicated a problem with the detector, most notably in the first few and last few crystals. The discrepancy between the measured bremsstrahlung signal and the theoretical signal along with verification of consistent CsI signal resulted in the creation of a single “camera correction” curve to be applied to all the data during the analysis; this correction is shown in Figure 7.7(b).

Along with correcting for the low light yield of the first and last few crystals, likely caused by poor crystal quality or inadequate capture of the CsI fluorescence, it was important to correct for the non-uniformity of the crystal light yield. This ensured that the calculations were performed with a more accurately represented signal curve. While the correction factor accounted for the shape of the signal along with some of the inherent non-uniformity, it was not complete and applying a smooth fit to the noisy data to obtain an ideal signal level could fully account for the residual noise. The simulated signal in Figure 7.7(a) and the detector response curves in Figure 7.4 indicate that the CsI signal should not vary as significantly as the experimental data shows. The non-uniformity of the crystal light output was a source of error in the spectral calculations as the noise allowed for several different best-fit lines to the same data. For this work, it was useful to define a few different parameters to aid in describing the quality of the fit. The “experimental error” was defined as the standard deviation of the difference between the data curve and the best fit approximation, normalized by the integrated signal level of the ideal curve (Equation 7.1). Figure 7.8(a) shows an example of the experimental error calculation with the red arrows indicating the amplitude difference between the data and the best fit. The best fit to the data was calculated by performing 6th, 7th, and 8th degree polynomial fits to the data and choosing the curve with the highest fit quality.

\[
\text{Experimental Error} = \frac{\sigma(\text{Raw Data} - \text{Ideal Signal})}{\int \text{Ideal Signal}}
\]  

(7.1)
7.2.5 Iterative Calculations

An algorithm was written to calculate the input spectrum that would match the data after simulating its interaction with the Geant4 CsI array. The algorithm is based on the YOGI code [146, 147] in which the spectrum is calculated by introducing perturbations to an assumed exponential shape and checking the result of those perturbations against the data curve. The form of the exponential spectrum used in this algorithm is shown in Equation 7.2; this form was the best fit equation to a spectrum produced by a simulated inverse Compton scattering interaction. In this equation, \( A \) is the amplitude of the spectrum, \( E \) is the energy range, and \( E_{\text{crit}} \) is a characteristic energy of the spectrum with 49% of the photon energy radiated below \( E_{\text{crit}} \) and the mean photon energy is \( E_{\text{crit}}/3 \).

\[
\frac{dN}{dE} = A \times E^{-2/3} \times e^{-\frac{E}{E_{\text{crit}}}} \tag{7.2}
\]

Perturbations were made to the amplitude and/or the critical energy of the equation depending on which option improved the fit most. Iterating through these perturbations continued until the fit could not longer be improved by adjusting the equation of the incident spectrum. The quality of the fit between the data and calculated signal was measured by calculating the \( R^2 \) value as described by Equation 7.3; where \( s \) is the data, \( f \) is the calculated signal, and \( \bar{s} \) is the average of the data. Similar to the experimental error, the calculated error is defined in Equation 7.4 and can be seen in Figure 7.8(b). An example of the calculated signal with the raw data and the calculated spectrum can be seen in Figure 7.9(a) – (f).

\[
R^2 = 1 - \frac{\sum (s - f)^2}{\sum (s - \bar{s})^2} \tag{7.3}
\]
Figure 7.8: (a) Comparison between the data signal and the ideal signal obtained by a best fit. The red arrows show some of the amplitude differences that were measured in calculating the experimental error. (b) Comparison between the calculated signal and the ideal signal for the same shot as (a). The difference between the two signals is used to calculate the relative error. (c) Relationship between the critical energy and the ratio of the first data point to the first point of the calculated signal. (d) Relationship between the critical energy and the ratio of the last data point and last point of the calculated signal. Each color ‘×’ represents the 7 background subtractions of a single shot.
<table>
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<th>Parameter</th>
<th>BG 1</th>
<th>BG 2</th>
<th>BG 3</th>
<th>BG 4</th>
<th>BG 5</th>
<th>BG 6</th>
<th>BG 7</th>
</tr>
</thead>
<tbody>
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<td>0.9919</td>
<td>0.9855</td>
<td>0.9886</td>
<td>0.9882</td>
<td>0.9927</td>
<td>0.9689</td>
</tr>
<tr>
<td>FPR</td>
<td>0.9913</td>
<td>1.0025</td>
<td>1.0354</td>
<td>1.0299</td>
<td>0.9770</td>
<td>0.9912</td>
<td>1.0151</td>
</tr>
<tr>
<td>$E_c$ [MeV]</td>
<td>21.22</td>
<td>52.26</td>
<td>98.63</td>
<td>44.71</td>
<td>60.12</td>
<td>50.61</td>
<td>33.84</td>
</tr>
<tr>
<td>$A$</td>
<td>2.11</td>
<td>1.43</td>
<td>0.98</td>
<td>1.68</td>
<td>1.33</td>
<td>1.47</td>
<td>1.62</td>
</tr>
</tbody>
</table>

Table 7.1: Table showing the values returned by running the fitting algorithm to the seven background subtracted data curves. The results show that the fit with the highest $R^2$ does not always correspond to the fit with the highest first point ratio (FPR), which is why both needed to be considered in choosing the background that resulted in the best fit.

\[
\text{Calculated Error} = \frac{\sigma |\text{Calculated Signal} - \text{Ideal Signal}|}{\int \text{Ideal Signal}} \tag{7.4}
\]

As mentioned previously, determining the proper background subtraction was challenging in this work as it had a significant effect on the resulting fit and spectrum. Starting with a $[1 \times 33]$ vector of data having the correction factor applied, $s$, seven different background shots, $b$, were subtracted such that $s'$ is a matrix derived from the concatenation of the vectors $s - b$ of size $[7 \times 33]$. Now there are seven different data curves for each shot, each with its own potential background subtraction. The fitting algorithm was run to calculate a critical energy ($E_{\text{crit}}$), amplitude ($A$), fit quality ($R^2$), and ratio between the first point of the data and first point of the fit for each of the seven options. An example of these numbers for a data shot is shown in Table 7.1. The relationship between the first point of the data and the fit is important as the fit quality of the first few points has a much higher effect on the calculated critical energy than the last few points. Figure 7.8(c) shows that when the ratio of the first point of the data to the first point of the calculated signal is greater than 1, the algorithm over-estimates the critical energy and when it is less than 1, the critical energy is under-estimated. Figure 7.8(d) shows that the fit quality of the last point does not affect the critical energy.
To define which background subtraction was the “correct” subtraction, the $R^2$ value was multiplied by the first point ratio for each of the seven calculated signals and the result closest to unity was chosen as the proper background subtraction. This method took into account both the quality of the fit and accuracy of the first point fit as it was possible to have background subtraction with a high fit quality overall but poor fit on the first point resulting in a heavily over- or under-estimated critical energy.

Upon determining the correct background subtraction for each shot, a noise analysis was performed to determine the sensitivity of the critical energy calculations to noisy data. For each shot, the experimental error was used as the amplitude metric and random noise was added to each bin of the ideal signal ranging from $\frac{1}{2}$ to $2 \times$ the experimental error. Once new signal was generated with the synthetic noise, an ideal signal was found for each of the new signal curves and the fitting algorithm was re-run. Since the correct background was already chosen for each shot, the error due to the synthetic noise indicated the error of the original critical energy calculation as this depended on the ideal signal that was generated by performing a best-fit to the data. The results of the noise analysis can be seen in Figure 7.10.

7.3 Results

Comparing the calculated error to the experimental error is a good way to quantify the quality of the fit with respect to the original data. To determine how well the calculated signal fits the data, the “relative error” is defined in Equation 7.5; the integration of the ideal signal cancels when these two errors are divided by one another.

$$\text{Relative Error} = \frac{\sigma(\text{Calculated Signal} - \text{Ideal Signal})}{\sigma(\text{Raw Signal} - \text{Ideal Signal})} \quad (7.5)$$
Figure 7.9: Data and calculated fit plotted with corresponding spectra. The data obtained from the CsI detector is marked with the blue dots and the calculated fit is plotted as the solid line among the dots. (a) + (b) Example of a low critical energy fit and spectrum of $E_{\text{crit}} = 7.08$ MeV. (c) + (d) Example of a moderate critical energy fit and spectrum with $E_{\text{crit}} = 40.63$ MeV. (e) + (f) Example of a high critical energy fit and spectrum with $E_{\text{crit}} = 109.58$ MeV.
Figure 7.10 compares the critical energy, detector signal level, integrated spectrum level and fit quality parameters to one another. In (a) and (b) of the figure, are plots that describe the physics of the inverse Compton scattering interaction and in (c) - (f) are plots that describe the success of the iterative algorithm. Figure 7.10(a) shows the critical energy plotted against the signal level of the detector. There appears to be a very weak positive trend between the CsI signal level and the calculated critical energy if a few of the higher, outlier signals are ignored. This is somewhat to be expected as a more successful overlap of the electron beam with the scattering laser should produce more photons and higher energy photons. If the trend were stronger, it would indicate that the number of electrons interacting with the laser pulse on each shot remained constant. The rather weak trend, indicates that the number of electrons scattered by the laser on each shot was not constant, likely due to the charge variability in the electron beam and alignment difficulties. In Figure 7.10(b), the relationship between the detector signal level and integrated spectrum is compared. If the critical energy was the same for all spectra, all the points on this plot would lie in a linearly increasing line. The increasing trend is expected, however, because higher energy photons deposit more signal into the detector than lower energy photons, (see Figure 7.4) there is not a direct, linear relationship between number of photons and CsI signal level that can be extracted.

The trends shown in Figure 7.10(c) - (f) are more telling of the fitting algorithm than the physics of the interaction. In (c), the critical energy is plotted against the first point ratio for the noise analysis data as opposed to the background subtracted data as is the case with Figure 7.8(a). This now shows that the relationship between the first point and the critical energy is much weaker and that most of the points are very close to unity. This indicates that the critical energy measurement is more trustworthy and the over- or under-estimation is within a reasonable error. The plot in (d) shows that there is no relationship between critical energy and relative error.
Figure 7.10: Measurements of various trends across 20 shots. (a) The integrated signal level on the CsI detector plotted against the measured critical energy. (b) Shows the relationship between the CsI detector signal level and the area under the curve of the calculated spectrum. (c) The relationship between the first point ratio and critical energy after the correct background was subtracted. (d) Shows there is not a relationship between the relative error (defined in Section 7.2.5) and the critical energy. (e) Relative error plotted against the detector signal level showing a slight positive trend. (f) Relative error plotted against the added noise. The minimum and maximum added noise are highlighted with red circles and blue squares respectively.
Figure 7.11: (a) Background-corrected CsI signal. (b) Consecutive electron spectra. Each spectrum has been normalized to its own maximum. The red and blue colored bands represent the ±1 standard deviation region for the energy of the spectral features in the ‘Bright’ and ‘Beam off’ shots respectively [145].

This is desirable as a relationship between the energy of the photons and the fit quality would imply that there were errors with the simulations. The plot indicates that over 75% of the points have a relative error of less than 2.0, indicating a strong fit quality among most of the shots. Plot (e) shows that there is a slightly positive relationship between the relative error and detector signal level which is unexpected because the relative error is independent of signal level as seen in Equation 7.5. Further analysis revealed that the result of this slight positive trend is due to the positive trend that exists between signal level and the numerator of the calculated error and the lack of a positive trend between the numerator of the experimental error and the signal level. Lastly, plot (f) shows the relationship between the relative error and the added noise. As expected, the relative error decreases with increasing noise because higher noise means higher experimental error, which is the denominator of Equation 7.5. The highlighted red and blue points show the minimum and maximum added noise respectively. As expected, the minimum noise resulted in higher discrepancies between the calculated fit quality and the ideal signal.
7.4 Discussion

The data set that was used in this analysis for measuring a signature of radiation reaction consists of 18 consecutive shots, the first eight of which were performed with the f/2 scattering beam on and for the subsequent ten shots, it was turned off. Figure 7.11 shows the electron spectra and corresponding CsI signal for this shot series. The figure highlights four of the shots when the scattering beam was on as having particularly high CsI signal and a particularly low edge on the electron spectrum. The simultaneous decrease in the electron spectrum feature and increase in the CsI signal is indicative of a radiation reaction event. By measuring the decrease in the electron spectrum edge for the four shots where a radiation reaction may have occurred, it was possible to infer the $a_0$ of the laser that interacted with the beam. This was done by using both a quantum model and a classical model based on the Landau-Lifshitz equation [148, 149]. Measuring the average energy of the feature between the bright shots and the beam-off shots indicated that there was a shift of the edge from $E_{\text{initial}} = (550 \pm 20)$ MeV to $E_{\text{final}} = (470 \pm 10)$ MeV. Assuming the electron beam interacted with the laser pulse during the entire 45 fs pulse duration, the resulting $a_0$ required to cause this shift is: $a_0 = 10 \pm 2$ for a quantum model and $9 \pm 1$ for a classical model based on the Landau-Lifshitz equation [145, 148, 149]. Although the $a_0$ of the f/2 scattering beam was measured to have a peak value of about 20, it is expected that timing errors and finite beam overlap between the electrons and laser would result in a much lower $a_0$ of the interaction.

The measured electron spectra edge shift is only half of the picture when confirming a radiation reaction. From the electron energy and inferred $a_0$, it was possible to calculate the critical energy of the gamma ray spectrum produced from the interaction and compare it to the measured spectra from the CsI crystal array. To better understand the different models describing the interaction between a relativistic electron beam and ultra-intense laser, Figure 7.12(a) shows simulated relationships between
Figure 7.12: (a) Simulated measurements of $E_{\text{crit}}$ assuming the collision of a plane wave of given peak $a_0$ with the experimentally measured electron spectra. The shaded regions represent ±1σ variations arising from the measured electron beam spectral fluctuations. (b) Experimentally measured $E_{\text{crit}}$ as a function of $\Delta E/E$ measured at the electron spectral feature (points). The shaded areas correspond to the results a hypothetical ensemble of identical experiments would measure 68% of the time under different assumed radiation reaction models for a uniform distribution of $a_0$ between 4 and 20 [145].

the laser $a_0$ and the critical energy of the resulting spectrum. As expected, with increasing $a_0$, the critical energy deviates further from a linear trend due to existence of the $\chi_0$ correction parameter as described in Section 2.3.2.

To make a comparison of the four bright shots to the model, Figure 7.12(b) shows the critical energy of the data points as a function of fractional energy loss of the electron spectrum edge feature, $\Delta E = (E_{\text{initial}} - E_{\text{final}})/E_{\text{initial}}$, plotted alongside the model. The shaded regions show the $(\Delta E/E, E_{\text{crit}})$ space that is 68% consistent with the models from Figure 7.12(a) to account for experimental noise. The range of fractional energy loss plotted in Figure 7.12(b) was calculated using experimental electron spectra and $a_0$ ranging from 4 to 20. The overlay of the points on the figure indicate that the data agrees better with the quantum model of radiation reaction better than the classical model. This agreement with the models, the observed correlation between electron beam energy, gamma signal, and $E_{\text{crit}}$ are strong indications that some form of radiation reaction occurred on these shots.
In conclusion, a CsI gamma ray spectrometer was developed by placing an array of CsI crystals parallel to the gamma beam propagation direction in order to measure penetration depth. The measurements of CsI scintillation were made with the counter-propagating beam turned on and off as shown in Figure 7.5. The figure shows that on average, the CsI bricks produced a higher light yield when the scattering beam was turned on, indicating that the counter-propagating laser pulse caused the creation of the highest energy gamma rays yet reported from an all-optical source.

The CsI detector was used in conjunction with an electron spectrometer and a quantum model of electron dynamics to confirm the presence of a radiation reaction on four of the data shots. The measurement of hard gamma rays made possible by this detector allowed for this small data set of radiation reaction shots to be further confirmed by comparing the $a_0$ value independently inferred from the critical energy and the fractional energy loss. The ratio between the two calculated laser potentials: $R = a_0(E_{\text{crit}})/a_0(\Delta E)$ serves to confirm the validity of the measurements and consistency of the model. For the quantum and classical models, the measured $R$ values are: $R = 0.8^{+0.7}_{-0.3}$ and $R = 0.6^{+0.3}_{-0.2}$ respectively. The quoted $R$ value is the median among the four shots with the limits as the 16th and 84th percentiles [145]. This shows that the quantum model of radiation reaction does a much better job at bringing the inferred $a_0$ into agreement whereas the classical model under-predicts the $a_0$ experienced by the electron beam from the critical energy calculation. Therefore the data supports the quantum model of radiation reaction being necessary to model the interaction under these conditions.

The detector was used as a spectrometer by perturbing an assumed exponential spectrum and using Geant4 simulations to match the detector response to the data. With Geant4 simulations performed in advance, this detector setup along with the algorithm could be implemented in future experiments as a gamma ray spectrometer.
capable of producing a spectrum on a shot-by-shot basis. To improve the detector design, it would be beneficial to change the 9 mm thick steel side plate to a thinner, lower-Z material to mitigate the absorption of gamma rays in the housing. It would also be helpful to remove the faceplate of the detector that restricts the fluorescence to circular holes so that each crystal's fluorescence can be captured entirely.
8.1 Summary

This thesis describes experimental work centered around measurements of high-energy radiation production from a laser-plasma accelerator. The experiments and methods presented here are intended to show the diverse research capabilities of laser-driven accelerators. While the current technology for laser-plasma accelerators cannot yet match conventional technology, LPAs are highly versatile research tools with much better temporal resolution and on their way to becoming portable and highly tunable.

Chapter IV described preliminary work in measuring un-trapped electrons. The ring feature presented in this chapter was first noticed after placing an imaging screen on-axis near the gas cell to view the beam profile. Further work revealed the ring to be a consistent feature that is the result of un-trapped electrons receiving a forward momentum kick as they pass behind the wakefield buckets. Analysis of the data suggests that the ring structures may be density dependent, potentially related to the bubble size. Simulations were performed that confirmed the source of the electrons forming the ring but could not strongly affirm any relationship with the plasma density. Recent simulation work also does not show a relationship between ring divergence and density [101], however theoretical models predict some variation [103].
and additional work for this experiment would be beneficial.

The goal of the experiment described in Chapter V was to develop a better understanding of the x-ray photon source generated through betatron oscillations during the laser wakefield acceleration process on HERCULES. This experiment also served to develop the single photon counting algorithm that is used to generate x-ray spectra up to 30 keV on an ANDOR CCD camera. The spectra generated by the betatron source were characterized by comparing the critical energy across various gas cells and gas types. The results of this experiment showed that by forcing electrons to undergo self-trapping as opposed to injection via ionization, it was possible to consistently produce, on average, x-ray spectra with higher energy photons.

Perhaps the more important result of this work was an understanding of the single photon counting process, when it is viable, and its limitations and drawbacks. Section 3.6 detailed the single photon counting process, much of which was worked out during the course of analyzing spectra for this first experiment. The Monte Carlo simulations and double hit analysis were both developed to be used in this experiment so that an accurate representation of the betatron x-rays was generated under various experimental conditions.

In Chapter VI, an experiment was performed to demonstrate the application of betatron radiation as a laboratory tool. Betatron radiation was used with an ultrafast laser to probe the K-edge of aluminum on the femtosecond timescale. The experiment was set up in a pump-probe arrangement with the probe beam generating the wakefields and the pump beam heating the aluminum target. Both the pump and probe laser beamlines originated from the same seed in order to ensure a high degree of temporal accuracy.

The results of this work showed a blue-shift of the energy of the aluminum K-edge when the metal was heated. This blue shift is consistent with previous work and was possibly due to ultrafast ionization events creating a stronger Coulomb attraction.
between the K-shell electrons and the nucleus. The ionization events are carried out primarily by a hot electron population created by the incident pump laser. Because this happened on the femtosecond timescale, the bulk of the aluminum was still cold so ionization potential depression, which would act to red-shift the K-edge, had not yet developed. This work demonstrated the benefit and capability of using all-optical sources and betatron radiation in pump-probe style experiments.

The inverse Compton scattering experiment that was described in Chapter VII that was performed on the Astra-Gemini laser demonstrated a new regime of laser-electron interaction. The evidence of radiation reaction in this work shows successful interaction of the laser and electron bunch in a regime where the force of the radiation emitted by the electron bunch became relevant. This was the first time that radiation reaction has been measured in the laboratory setting. The analysis of the ICS data was made possible by the development of a novel gamma ray spectrometer and an electron spectrometer. The electron spectrometer is a very simple and well-developed diagnostic and was used to measured a decrease in the electron energy due to the scattering interaction.

The gamma ray spectrometer was a more complicated diagnostic developed by placing a CsI crystal array parallel to the gamma ray beam. This allowed for penetration and vertical divergence measurements of the gamma photons. Using Monte Carlo simulations to model gamma photons entering the detector array, it was possible to determine the incident gamma spectrum by matching a simulated signal to the data. Background analysis and noise analysis were performed for each of the shots to ensure the accuracy of the fit to the data. The $a_0$ of the laser was independently confirmed by the critical energy of the spectrometer and the decrease in the electron energy.
8.2 Future Work

The work presented and cited in this thesis displays a very impressive yet immature technology. Hardly a decade has passed since lasers were first used to accelerate a monoenergetic bunch of electrons, and already significant progress has been made in the field to that point that multi-GeV beams can be produced on a scale 10,000 times smaller than conventional accelerators. These electron beams can be used to generate photon sources ranging from a few keV to over 100 MeV without the use of large-scale undulators. Despite the impressive steps that have been made in understanding and controlling the wakefield process, this technology is still not as tunable, robust, or as stable as large scale accelerator technology and for this reason there is still significant work to be performed on laser wakefield acceleration. In this section, additional work related to the results presented in this thesis is suggested.

As mentioned in the Chapter IV introduction, there are several interesting experiments that can be continued with the on-axis rings. While they are likely not much use for driving any additional physics, detecting the ring on a shot-to-shot basis may prove to be a useful diagnostic. Under constant conditions, the number of electrons interacting with the laser pulse should be relatively equal across shots. Under this assumption, measuring the ratio of charge in the ring (un-trapped electrons) to the charge in the beam (trapped electrons) could provide a metric for trapping success. Using this diagnostic for real-time feedback could help optimize the experimental conditions and develop a more efficient LWFA system. The difficulty in implementing this type of diagnostic, however, is that both the charge in the ring and in the beam need to be measured simultaneously. The charge in the beam is too dense to be measured on-axis as shown by the saturation in Figure 4.2, and the ring must be measured before the magnet of the electron spectrometer. Geometry limitations would make it very difficult to implement this on HERCULES as there is not sufficient space between the gas target and the magnet to fit a permanent imaging screen.
Another interesting piece of future work would be to continue simulations and experimental measurements to discover any relationship between the ring size and LWFA conditions. If the data revealed that the divergence of the ring is related to either the bubble size, laser potential, plasma density, or some other parameter, it could provide useful feedback about the wakefield environment on each shot. On HERCULES, there are currently very few diagnostics that provide information about the wakefield structure or trapping success other than the existence of an electron beam. Implementing this on-axis detector and continuing work to understand its relationship to plasma conditions could greatly improve understanding of the experimental environment.

The results of the first experiment described in this thesis in Chapter V showed a difference in the x-ray spectrum for ionization trapping compared to self-injection for a single-stage gas cell, it would be interesting to continue the measurements with varying levels of nitrogen mixed into the helium. Theoretically, adding greater amounts of N$_2$ beyond the 2.5% from the experiment would further lower the critical energy of the spectrum, but finding out how much further may be interesting. Additional work for this data set might also include varying the spot size of the laser while maintaining a constant $a_0$. Increasing the focal spot of the laser should increase the size of the bubble which might lead to larger amplitude betatron oscillations and thus a higher critical energy.

The pump-probe experiment proved to be a very challenging undertaking with the HERCULES laser. There are several things that can be done to improve this experiment, some of which will require a different facility. The first improvement to be made when moving forward with this experiment is related to the target. Aluminum with thickness of 4 $\mu$m was chosen for several reasons. Aluminum is cheap and ubiquitous and the K-edge is below 2 keV, which means there are plenty of x-rays for absorption measurements, and a thickness of 4 $\mu$m provided enough contrast to
be able to see a difference in the x-ray flux on the camera between the reference and absorption halves. Although 4 \( \mu m \) made it easy to see an “edge” appear in the data, it is optically thick and therefore does not receive significant or uniform heating from an infrared laser like Hercules. To obtain more uniform heating of the target, a thinner (< 1\( \mu m \)) and perhaps different metal should be used. There are good reasons to use aluminum, but it may be more beneficial to choose a metal with a higher K-edge absorption feature such as titanium or iron. The higher absorption feature means that it will be easier to shield the camera from the substantial flux of low energy photons and keep the hits in the single photon regime.

Another useful change to implement moving forward with this experiment pertains to the geometry. In order to have a stable and uniform x-ray flux hitting the CCD detector on every shot, it is necessary to put the camera very far away from the source (\( \sim \) a few meters). Doing this, however, will substantially increase the magnification of the solid target on the camera and make measurements of any absorption feature very difficult to distinguish. When using a setup with a half-aluminum, half-reference target (or something similar) as described in Chapter VI, it is ideal to have the edge appear on the camera in the exact same place on every shot so that even if the target is too thin to distinguish a significant absorption feature, calculations of the spectrum can be made. Moving both the detector and solid target far away from the source will produce much higher stability and consistency between shots.

The inverse Compton scattering experiment performed on the Gemini laser was largely successful. The primary aim of the experiment was to measure a radiation reaction in the laboratory and the work on the gamma ray spectrometer allowed that to be confirmed. With the methodology described in this thesis, it would be very interesting to implement this algorithm with the CsI detector in real-time during another inverse Compton scattering experiment. One of the greatest challenges of colliding an electron beam with a counter-propagating laser is the spatial alignment
and timing. Real-time results of the generated gamma ray spectrum may greatly increase the ability to find the best overlap position for the two beams. Additionally, it would be interesting to extend this detector to different photon energies. It may be possible to use this detector with betatron radiation by decreasing the size of the crystals or using a different material that is more sensitive to lower energy photons. Using this style of detector with betatron radiation would allow the results to be confirmed with the single photon counting techniques described in this thesis.

Compared with conventional accelerator technology, demonstration of laser wake-field acceleration is a very new research tool. The progress and diversity of the experiments have been significant since the first quasi-monoenergetic beams were generated in 2004. The technology will continue to grow as lasers become more powerful, more stable, and easier to operate. There is no doubt that laser-plasma accelerators will continue to develop and someday rival the research performed at large, radio-frequency accelerator facilities. The development of LWFA as a light source will open the door for new research to be carried out anywhere. The work in this thesis demonstrates new physics in the laboratory and lays out a foundation for pump-probe experiments and detection of high energy photons.
BIBLIOGRAPHY


