An adjustable short-focal length, high-gradient PMQ electron-beam final-focus system for the PLEIADES ultra-fast x-ray Thomson source

A dissertation submitted in partial satisfaction of the requirements for the degree Doctor of Philosophy in Physics

by

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Dedicated to my family
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G. Travish, P. Frigola, J. Lim, J. Rosenzweig, A PMQ-based, ultra-short focal length, final focus system for next generation beam-radiation and beam-plasma


quadrupole based electron beam final focus system, Physical Review Special Topics - Accelerators and Beams 8, 072401 (2005)
In the span of a 100 year since the discovery of first x-rays by Roentgen that won him the first Nobel prize in physics, several types of radiation sources have been developed. Currently, radiations at extremely short wavelengths have only been accessed at synchrotron radiation sources. However, the current 3rd generation synchrotron sources can only produce x-rays of energy up to 60 keV and pulse lengths of several picoseconds long. But needs for shorter wavelength and shorter pulse duration radiations demanded by scientists to understand the nature of matter at atomic/molecular scale initiated the new scientific research for the production of sub-picosecond, hard x-rays. At the Lawrence Livermore National Laboratory, a Thomson x-ray source in the backscattering mode – a head-on collision between a high intensity Ti:Sapphire Chirped Pulse Amplification laser and a relativistic electron beam – called the PLEIADES (Picosecond Laser-Electron Inter-Action for the Dynamical Evaluation of Structures) laboratory has been developed. Early works demonstrated the production of quasi-monochromatic, femto-second long, hard x-rays. Initially reported x-ray flux was in the low range
of $10^5$ photons per shot.

During the early stage of PLEIADES experiments, 15 T/m electromagnet final focusing quadrupoles (in a triplet lattice configuration) were employed to focus the beam to a 40-50 $\mu$m spot-size. A larger focal spot-size beam has a low-density of electron particles available at the interaction with incident photons, which leads to a low scattering probability. The current dissertation shows that by employing a 560 T/m PMQ (Permanent-Magnet Quadrupole) final focus system, an electron beam as small as 10-20 $\mu$m can be achieved. The implementation of this final focus system demonstrated the improvement of the total x-ray flux by two orders of magnitude. The PMQ final focus system also produced small electron beams consistently over 30-100 MeV electron beam energy, which enabled the production of x-ray energy over 40-140 keV.

In this dissertation, the PLEIADES Thomson x-ray facility will be described in detail includes the 100 MeV linac and the FALCON laser system. Later, we will discuss the design, construction and implementation of the PMQ final focus system in the beamline. The measurement of electron beam parameters before and after the final focus system will be presented. The beam measurements at the interaction region were accomplished with the use of both OTR (Optical Transition Radiation) imaged by a CCD camera and the fast streak camera for respective spatial and temporal alignments. The theoretical analysis in “real beam” effects and spacetime beam jitter effects will be given to help understand the observations. A 3D simulation tool developed for x-ray data analysis was used to provide direct comparisons with the x-ray flux, spectrum distribution and transverse x-ray profile.
CHAPTER 1

Introduction

Since the discovery of x-rays by Wilhelm Roentgen in 1895 [2], we have seen the steady development of high-brightness radiation sources for biological and material sciences applications. Steady progress in the theoretical understanding of radiation emission processes has enabled experimentalists to solve challenging problems using powerful mathematical models.

In the span of just over 100 years since the first observation of x-rays, the effort to develop high brightness radiation sources has grown to a level where vastly different disciplines have come together to invent a new generation of sources which is expected to produce unprecedented brightness on the order of $10^{35}$ photons/s/0.1% bandwidth/mm$^2$/mrad$^2$. This radiation will be contained in pulse durations approaching the atto ($10^{-18}$) to femtosecond ($10^{-15}$) range, thereby allowing the observation of structural changes (e.g. in a protein molecule) occurring on an atomic timescale; other possible applications for this new generation fs sources include biological imaging and femtosecond pump-probe dynamics [3, 4, 5, 6, 7, 8].

A strong candidate for the next generation radiation source is the x-ray free electron laser (such as the Linac Coherent Light Source (LCLS) [9, 10], which is currently under construction at Stanford Linear Accelerator Collider (SLAC) and the EUROFEL [11] at DESY).
An alternative approach to producing quasi-monochromatic, ultra-short x-rays has recently been shown successfully at Lawrence Berkeley National Laboratory (LBNL) [12]. This experiment employed Thomson scattering geometry where the electron beam and the photon beam collide at a 90 degree angle (“Short pulse Thomson scattering”). The initial result in short pulse Thomson scattering reported production of fs pulses of soft x-rays but generated a photon flux which was too small to serve as probes of complex molecule systems.

The thesis proposed here is that the short-focal length final focusing permanent-magnet quadrupoles lattice can increase the Thomson x-ray flux via production of higher electron density at the interaction point. Such improvement over the number of x-rays scattered can benefit the Thomson sources as low cost, laboratory scale x-ray sources for a wide range of researches. We will provide a detailed
review of still an alternative approach under investigation at the Lawrence Livermore National Laboratory: a high-flux, 180 degrees scattering geometry Thomson x-ray source called the “Picosecond Laser-Electron InterAction for the Dynamic Evaluation of Structures” (PLEIADIES) [1, 13, 14].

In the current chapter, we will review various types of x-ray sources and proposed future radiation facilities. We finish the chapter with a general theoretical treatment on the Thomson scattering theory in the linear and non-linear regimes.

1.1 A General Survey of High-brightness X-ray Sources: Past, Present & Future

In many applications, we require a nearly parallel and “small spot” radiation beam to accurately probe an experimental sample of a very small dimensions. To characterize such a radiation beam, we introduce the concept of brightness. The brightness of an x-ray source is defined as the number of photons emitted per second, per unit source area, per unit solid angle, per unit spectral bandwidth. Consider, for example, a rectangular source emitting $N$ photons per second, per unit bandwidth from a spatial area bounded by $L_x$ and $L_y$ and emission angles bounded by $\theta_x$ and $\theta_y$, the source brightness (no correlation between $x$ and $y$ terms) then is defined by [15]

$$B = \frac{N}{(L_x\theta_x)(L_y\theta_y)}. \quad (1.1)$$

We note that this definition of radiation source brightness is frequently used by researchers to measure the quality of radiation beams.

High brightness sources are desired by many scientists in order to gain an understanding of the microscopic structure of physical samples. This is one of the main motivations driving the advances in the accelerator-light source community.
By examining Eq. (1.1) in detail, we can see that there are a number of different ways the parameters can be optimized to obtain a peak spectral brightness. For instance, use of smaller source area and smaller divergences in the radiation beam can yield a significant increase in x-ray brightness.

We now will review various radiation sources such as synchrotron, FEL and Thomson scattering radiation sources.

1.1.1 The Synchrotron Radiation (SR) Source

Synchrotron radiation (SR) produced from relativistic charged particle motion in curved paths through static magnetic fields was first observed in the mid-20th century. In the last 50 years, a broad spectrum of synchrotron radiation derived from this physical mechanism has brought new insights into how matter is structured at a very small scale. Synchrotron radiation sources now are found in countries all over the world as centers for new science and technology. A brief history of synchrotron radiation source from its beginning to present-day is presented first in this section. The interested reader is recommended to read A. Robinson’s history of synchrotron radiation in [16] for more detail.

In 1944, a theoretical model of energy loss by relativistic particles due to radiating particles in a betatron was first developed by Ivanenko and Pomeranchuk [17]. By 1945, the classical theory of radiation from accelerated relativistic charged particles was worked out in large part by Schwinger [18]. In 1947, synchrotron radiation in the visible spectrum was observed at General Electric for the first time; this development became the stepping stone to the field of relativistic radiation source physics. Followed by the success of the first synchrotron light at GE, accelerator physicists began to construct and operate synchrotron sources parasitically at high-energy storage rings from 1960s to early 1970s. Dur-
ing this period, the first x-rays were observed at synchrotron radiation facilities such as the 6 GeV DORIS ring at Deutsches Elektronen-Synchrotron (DESY) in Hamburg and 2.5 GeV SPEAR ring at SLAC. Because the design of the storage ring was driven by the high energy experiments, the radiation output was severely limited. So the idea of a storage ring facility dedicated only for the operation of a synchrotron source became favored by many researchers. The existence of a dedicated synchrotron source would enable users to build spacious experimental stations and obtain high quality experimental data using increased output of radiation. By the 1980s, dedicated synchrotron sources, now known as “second generation” sources, were constructed, allowing more uses of the SR light. We note that today generation numbers are labeled on radiation sources from a specific era of major technological advancements to categorize differently.

The successful operation of a synchrotron source is dependent on the quality of the electron beam circulating in the storage ring. Specifically, the high stability of the electron beam orbit allows for the production of high quality, high brightness, and stable synchrotron radiation. During the first and the second generations, many storage ring technological problems relating to beam stability were identified by researchers at synchrotron facilities. Most notable problems were related to thermal variations, magnet power supply regulation, ground vibrations, and magnetic field coupling. A feedback regulation system that controls the beam-orbital errors at both a local and a global scale was developed and employed in “third generation” sources. Such an orbit feedback system was initially developed at SPEAR [19] providing local corrections and subsequently, followed by the development of a global feedback system at NSLS [20, 21] to stabilize the entire storage ring.

Another noteworthy system is the Chasman-Green achromatic magnet lattice
Figure 1.2: NSLS synchrotron light source at BNL.

[22]. This zero momentum-dispersion insertion system was first employed in “third generation” synchrotron sources to produce a bright synchrotron radiation from a low emittance electron beam. The minimization of the beam emittance by the achromatic magnet lattice produces a smaller beam. According to Eq. (1.1) when the radiating source transverse area goes down then the radiation brightness increases. The result of using this system was that the source brightness increased dramatically from the previous “second generation.” Initial successes with the Chasman-Green system at the National Synchrotron Light Source (NSLS) led to the adoption of this technology by other major “third generation” facilities such as European Synchrotron Radiation Facility (ESRF), Advanced Photon Source (APS), and Super Photon Ring-8 GeV (SPring-8). Presently, these facilities are capable of delivering to users high brightness radiation and high photon counts in angstrom-wavelengths produced from an undulator.

The ever increasing demands for coherent, high-brightness and ultra-short x-rays needed for the observation of physical dynamics occurring on atomic scales – i.e. in both spatial and temporal coordinates – launched a global basic research effort. “Fourth generation” sources will be based on the x-ray self-amplified
spontaneous emission (SASE) FEL technology [23, 24]. While the x-ray FEL is acknowledged as the principal “fourth generation” technology, a small community of accelerator physicists seek an x-ray source based on the different physical concept. The Thomson x-ray source is an example of an alternative technology that is currently under research and development: the idea is to generate a quasi-monochromatic, high-flux and ultra-short duration x-ray pulse on a tabletop setting. The Thomson x-ray source can be constructed at a fraction of the cost of the “fourth generation” x-FEL, but with the price of a much reduced brightness. Prior to devoting the rest of this thesis on the Thomson x-ray source experiment, we would like to cover the basics of the x-ray FEL in the next section.

1.1.2 The X-ray Free Electron Laser (FEL)

SR beams obtained from an undulator at “third generation” sources are limited in transverse coherence and brightness. In the next generation radiation source, this problem will be resolved and the performance is expected to exceed the “third generation” performance by orders of magnitude in coherence, brilliance/brightness, and pulse time-duration (femtoseconds). This “fourth generation” will be based on an x-ray free electron laser to be constructed at a GeV-linear accelerator.

An x-ray FEL naturally produces coherent light radiation of a narrow bandwidth in spectrum. This coherence is obtained from electrons traversing the undulator field grouped in micro-bunches with a separation equal to the radiation wavelength. The FEL process starts from a randomly distributed electrons incoherently emitting spontaneous radiation field proportional to the number of electrons in a bunch. As the beam propagates downstream of the long undulator, micro-bunches start to form due to the resonant coupling between the electron
beam and the co-propagating radiation field. Towards the end of the undulator, well separated micro-bunches are formed with each micro-bunch acting as a point-like charged particle producing undulation radiation. The radiation intensity is expected to be proportional to the square of the number of electrons in a micro-bunch acting as coherent emitters. The modulation of an electron beam into micro-bunches from the resonant interaction with EM wave is called the SASE process [25, 26, 27]. Since this undulation radiation is a single-pass process, it eliminates the need to use FEL oscillators in the x-ray regime. Hence, problems associated with the quality factor of the optical cavity and the radiation damage to cavity mirrors are evaded altogether in x-ray FEL.

However, an x-ray FEL is a large scale light facility requiring a GeV energy accelerator with highly uniform undulator system of a 100 meter long at least. The current R&D trend in the accelerator community is seeking to develop an affordable table-top scale radiation source. The new light-source experiment based on the Thomson scattering process or the light-electron interaction [28] is currently underway as an alternative source for a high-flux, energy-tunable radiation source. The core technologies involved are the chirped-pulse amplification laser
system and the photo-cathode injector linac. So what separate Thomson sources from x-ray FELs are the radiation extraction mechanism in Thomson sources involves an extremely short interaction length (a few millimeters) determined by the Rayleigh range, and the nature of Thomson process allows the production of hard x-rays with a low-energy electron beams (MeV range). On the other hand, x-ray FELs are fully coherent light source whereas Thomson x-ray sources are not.

1.1.3 The Thomson X-ray Source

The Compton/Thomson scattering process as the explanation for an energy loss mechanism in the interstellar electrons via interaction with “thermal” photons (e.g. cosmic microwave background) was first postulated by Feenberg and Primakoff in 1948 [29]. The advent of coherent laser technology in 1960s enabled the recreation of the Thomson scattering process in a well controlled experimental setting [30, 31]. After the initial demonstration, the program declined and eventually became inactive altogether due to the low yield of Thomson x-rays. Inefficient scattering rate due to the small Thomson cross-section, \( \sigma_T = 6.65 \times 10^{-25} \text{ cm}^2 \) naturally resulted in low radiation brightness. But there has been a renewal of interest in the Thomson program with the development of new and more powerful laser and photo-injector electron beam source. In 1995, energy and momentum shifts of charged particles in a powerful laser field was observed by Meyerhofer, et.al. [32]. Their work inspired others like Hartemann [33] and Salamin [34] to reinvestigate the relativistic Thomson scattering experiment. In 1996 at Lawrence Berkeley National Laboratory, the first subpicosecond hard x-rays, \( 5 \times 10^4 \) photons at 30 keV on average, were produced employing the 90 degrees scattering geometry [12]. The success of this proof-of-principle experiment initiated new
experiments [35, 36]. Recently, a Thomson x-ray experiment based on an “all laser-based” scheme has been shown to produce Thomson radiation [37]. The electron beam is obtained from a laser wakefield plasma accelerator where the energy spread is small. This proof-of-principle experiment has shown the possibility for a true tabletop accelerator and ultra-short radiation source in the near future.

One of several variations in the experimental production of Thomson radiation is the PLEIADES Thomson backscattering experiment based on the linac system at LLNL. The basic working principle of PLEIADES is founded on the interaction between a high-intensity laser and a relativistic electron beam during a head-on collision which produces a tunable Doppler photon frequency upshift proportional to $4\gamma^2$ ($\gamma$ is the relativistic Lorentz factor). We will devote the rest of the thesis to the development and the results of the PLEIADES experiment. In the next section, we review the classical theory of Thomson scattering.

1.2 The Classical Thomson Scattering Theory

The interaction of charged particles and radiation can be categorized by the energy exchange or “recoil.” In the Thompson limit, the incident laser energy is negligible compared to the mass energy of the electron. The converse limit (involving a momentum transfer from incident laser to electron) is known as Compton scattering. In the Thomson limit, the final number of scattered photon is limited by the total Thomson scattering cross-section. The derivation of the cross-section will be based on the well-known early work on scattering between a free-standing/non-relativistic electron charge and a plane EM wave.

In the non-relativistic regime, a plane EM wave interacting with a free charge
causes the particle to execute an oscillatory motion thereby producing a dipole-pattern radiation. The frequency of the scattered wave is taken identical to that of the incident wave as long as the process is occurring in the non-relativistic limit. The Thomson cross-section in the case of an unpolarized incident radiation follows [38]:

$$\frac{d\sigma}{d\Omega} = \frac{r_e^2}{2} (1 + \cos^2 \theta).$$  \hspace{1cm} (1.2)

where $r_e = \frac{e^2}{4\pi \epsilon_0 mc^2}$ is the classical electron radius, and $\theta$ is an angle between the incident and the scattered radiation direction. From this formula, an integration over all solid angles gives the well-known total Thomson cross section

$$\sigma_T = \frac{8\pi}{3} r_e^2.$$  \hspace{1cm} (1.3)

The Thomson cross section for the electron is $6.65 \times 10^{-25} \text{cm}^2$.

The above results are only valid in the non-relativistic limit. At PLEIADES experiment, a high relativistic electron beam with a $\beta$ approaching one is used in backscattering mode with an incident high intensity laser to generate x-rays. Real beam effects also complicates the scattering process and the resulting x-ray patterns. Thus, a relativistic Thomson scattering theory is summarized in the next section.

1.2.1 The Relativistic Linear Thomson Scattering

We first define the coordinate system to be used in the relativistic linear ($a_0 < 1$) Thomson scattering theory. We consider the case of an incident plane-wave with a frequency $\omega_{\text{inc}}$ colliding with an electron at an angle $\theta_{\text{inc}}$ with respect to the $+z$ axis. The electron is also traveling with a velocity $\mathbf{v} = v \hat{z}$ in the $+z$ axis. Figure 1.4 shows the scattering orientation of the incident radiation and the electron. We also define the angle between the observation point and the initial
electron beam’s propagation direction to be \( \theta_{\text{obs}} \). The wavenumber 4-vectors in the lab frame for the incident and scattered radiation in terms of these angles are [39, 40]

\[
\mathbf{k}_{\text{inc}}^\mu = \left( \frac{\omega_{\text{inc}}}{c}, k_{\text{inc}} \sin \theta_{\text{inc}} \cos \phi_{\text{inc}}, k_{\text{inc}} \sin \theta_{\text{inc}} \sin \phi_{\text{inc}}, -k_{\text{inc}} \cos \theta_{\text{inc}} \right)
\]  
(1.4)

\[
\mathbf{k}_{\text{obs}}^\mu = \left( \frac{\omega_{\text{obs}}}{c}, k_{\text{obs}} \sin \theta_{\text{obs}} \cos \phi_{\text{obs}}, k_{\text{obs}} \sin \theta_{\text{obs}} \sin \phi_{\text{obs}}, k_{\text{obs}} \cos \theta_{\text{obs}} \right)
\]  
(1.5)

Lorentz transform the 4-vectors to the electron’s own rest frame, we obtain

\[
\mathbf{k}_{\text{inc}}' = \gamma \left( \frac{\omega_{\text{inc}}}{c} + \beta k_{\text{inc}} \cos \theta_{\text{inc}}, k_{\text{inc}} \sin \theta_{\text{inc}} \cos \phi_{\text{inc}}, k_{\text{inc}} \sin \theta_{\text{inc}} \sin \phi_{\text{inc}}, -k_{\text{inc}} \cos \theta_{\text{inc}} - \beta \frac{\omega_{\text{inc}}}{c} \right)
\]  
(1.6)

\[
\mathbf{k}_{\text{obs}}' = \gamma \left( \frac{\omega_{\text{obs}}}{c} - \beta k_{\text{obs}} \cos \theta_{\text{obs}}, k_{\text{obs}} \sin \theta_{\text{obs}} \cos \phi_{\text{obs}}, k_{\text{obs}} \sin \theta_{\text{obs}} \sin \phi_{\text{obs}}, k_{\text{obs}} \cos \theta_{\text{obs}} - \beta \frac{\omega_{\text{obs}}}{c} \right)
\]  
(1.7)

We assume that the Thomson limit is still valid in the electron’s rest frame such that recoil is negligible, \( \omega'_{\text{inc}} = \omega'_{\text{obs}} \), then

\[
\frac{\omega'_{\text{inc}}}{c} = \frac{\omega'_{\text{obs}}}{c}, \quad \gamma \left( \frac{\omega_{\text{inc}}}{c} + \beta k_{\text{inc}} \cos \theta_{\text{inc}} \right) = \gamma \left( \frac{\omega_{\text{obs}}}{c} - \beta k_{\text{obs}} \cos \theta_{\text{obs}} \right)
\]

After rearrangement of the previous equation and using the assumption that the process is occurring in vacuum, we obtain the frequency of a scattered photon to be

\[
\omega_{\text{obs}} = \omega_{\text{inc}} \left( \frac{1 + \beta \cos \theta_{\text{inc}}}{1 - \beta \cos \theta_{\text{obs}}} \right)
\]  
(1.8)

In the highly relativistic limit (\( \beta \approx 1 \)) and the peak energy backscattered at \( \theta_{\text{obs}} = 0 \), \( \gamma^2 \approx 1/2(1 - \beta) \). The peak scattered photon frequency observed in +z-direction is

\[
\omega_{\text{peak}} = 2\gamma^2 \omega_{\text{inc}} \left( 1 + \beta \cos \theta_{\text{inc}} \right)
\]  
(1.9)
Figure 1.4: The Thomson scattering geometry of an electron and an electromagnetic wave.

This expression for the frequency of scattered photon is true for any laser polarization at any chosen incident direction.

The Thomson differential cross section in the lab frame can be obtained from the electron's rest frame by Lorentz transformation. In the electron's rest frame, the differential cross section is given by [38]

$$\frac{d\sigma}{d\Omega} = r_e^2 |\vec{e}_{\text{obs}}' \cdot \vec{e}_{\text{inc}}'|^2,$$

(1.10)

where $\vec{e}_{\text{inc}}'$ and $\vec{e}_{\text{obs}}'$ are respectively the incident laser and the scattered photon polarization vectors. First, the scattered photon polarization $\vec{e}_{\text{obs}}'$ can be decomposed into two perpendicular components: $\vec{e}_{1,\text{obs}}'$ confined in the plane spanned
by $k'_{\text{obs}}$ and $k'_{\text{inc}}$ and the other, $\vec{\epsilon}'_{2,\text{obs}}$ contained in the plane perpendicular to $k_{\text{inc}}$ and $\vec{\epsilon}'_{1,\text{obs}}$. The definition of scattered photon polarization components are given by
\[
\begin{align*}
\vec{\epsilon}'_{1,\text{obs}} &= \cos \theta'(e'_x \cos \phi' + e'_y \sin \phi') - e'_z \sin \theta', \\
\vec{\epsilon}'_{2,\text{obs}} &= -e'_x \sin \phi' + e'_y \cos \phi'.
\end{align*}
\] (1.11)
where angles (in the electron rest frame) are shown in Figure 1.4. And the polarization state of the incident laser is $\vec{\epsilon}'_{\text{inc}} = \epsilon'_x e'_x + \epsilon'_y e'_y + \epsilon'_z e'_z$. After performing the summation over all polarization states, the Thomson cross section in the non-relativistic limit in the electron’s frame is given by [39]
\[
\frac{1}{r^2} \frac{d\sigma}{d\Omega'} = \epsilon'^2_x (1 - \cos^2 \phi' \sin^2 \theta') + \epsilon'^2_y (1 - \sin^2 \phi' \sin^2 \theta') + \epsilon'^2_z (1 - \cos^2 \theta') - 2 \epsilon' x \epsilon' y \cos \phi' \sin \phi' \sin \theta' \gamma - 2 \epsilon' x \epsilon' z \cos \theta' \sin \theta' - 2 \epsilon' y \epsilon' z \sin \phi' \sin \theta' \sin \phi' \sin \theta'.
\] (1.12)
The differential cross section in the rest frame (primed variables) will be Lorentz transformed into the electron’s lab frame (unprimed variables) using the following angular Lorentz transformations
\[
\begin{align*}
\cos \theta' &= \gamma \frac{c(k_\parallel - \beta \omega)}{\gamma (\omega - \beta c k_\parallel)} = \frac{\cos \theta - \beta}{1 - \beta \cos \theta}, \\
\sin \theta' &= \gamma \frac{k_\parallel c}{\gamma (\omega - \beta c k_\parallel)} = \frac{\sin \theta}{\gamma (1 - \beta \cos \theta)}.\end{align*}
\] (1.13) (1.14)
The azimuthal angle $\phi$ about $+z$-direction, is a Lorentz invariant. After inserting these trigonometric functions in Eq. (1.12), we find
\[
\begin{align*}
\frac{1}{r^2} \frac{d\sigma}{d\Omega'} &= \epsilon'^2_x \left(1 - \frac{\cos^2 \phi \sin^2 \theta}{\gamma^2 (1 - \beta \cos \theta)^2}\right) + \epsilon'^2_y \left(1 - \frac{\sin^2 \phi \sin^2 \theta}{(1 - \beta \cos \theta)^2}\right) + \epsilon'^2_z \left[1 - \left(\frac{\cos^2 \theta - \beta}{1 - \beta \cos \theta}\right)^2\right] - 2 \epsilon' x \epsilon' y \cos \phi \sin \phi \sin \theta \gamma - 2 \epsilon' x \epsilon' z \cos \phi \sin \theta \sin \theta \gamma - 2 \epsilon' y \epsilon' z \sin \phi \sin \theta \sin \phi \sin \theta \gamma.
\end{align*}
\] (1.15)
Figure 1.5: The linear Thomson x-rays have a dipole radiation pattern. The theoretical curves in the plane $x$-$z$ ($\phi = 0$) are generated for different electron relativistic velocities.

where $\epsilon_{x,y,z}$ are the polarization state in the lab frame. In this generalized coordinate system, the Lorentz invariance does not hold for the transformation of laser polarizations ($\vec{\epsilon} \neq \vec{\epsilon}'$), because of $\vec{k}_{inc} \neq \vec{k}'_{inc}$ [41]. Nevertheless, the differential cross section in the electron’s lab frame is

$$\frac{d\sigma}{d\Omega} = \frac{d\sigma}{d\Omega'} \frac{d\Omega'}{d\Omega} = \frac{d\sigma}{d\Omega'} \frac{d\cos \theta'}{d\cos \theta} \frac{1 - \beta^2}{(1 - \beta \cos \theta)^2}.$$ (1.16)

We consider the simplest case of the incident photon polarization and the direction of dipole oscillation both along the $x$-axis and the initial electron trav-
eling in the $z$ direction. In this case, we can assume that $\vec{\epsilon} = \vec{\epsilon}'$, because $\vec{k}_{\text{inc}}$ is aligned in $z$. In the differential cross-section expression, the only surviving term is the one with $\epsilon_x$. The total energy scattered per unit solid angle (obtained by multiplying $\hbar \omega$) for this special case is given by

$$
\frac{dU}{d\Omega} = \hbar \omega_s r^2 \left( \frac{1 - \beta^2}{(1 - \beta \cos \theta)^2} \right) \left( 1 - \frac{\sin^2 \theta \cos^2 \phi}{\gamma^2(1 - \beta \cos \theta)^2} \right),
$$

(1.17)

where the scattered frequency $\omega_s$ can be given in terms of the incident frequency $\omega_{\text{inc}}$ in Eq. (1.8). The most notable effect from the Thomson scattering is that in the ultra-relativistic limit, the radiation is tightly focused in a forward cone. This behavior is determined significantly by the $(1 - \beta \cos \theta)^{-3}$ term. The relativistic effect of a narrowing angular radiation distribution in the direction of the electron’s main motion is shown in Figure 1.5 where we go from the free-standing electron case to the relativistic electron case. The other notable result due to the presence of $\phi$ in the distribution is that the radiation lobe is flatter in $y$-axis than in $x$-axis resulting in a slightly oblong radiation profile in $y$ direction.

### 1.2.2 The Spectral Broadening from 3-D Effects

The theory of Thomson scattering between an electromagnetic plane-wave and a single electron particle has been covered in the previous section. In a real experimental situation, idealized scattering sources need to be replaced by an accurate complex physical model that accounts for a focusing in the incident wave near the interaction point and for “beam-bunch” with a typical particle number of $> 10^9$ with associated beam parameters. Real beams have inherent energy spread, angular divergence, and other small bandwidth effects, which all add up to produce a significant spectral broadening effect in the observed Thomson radiation [39, 42, 43]. The various spectral broadening effects will be covered in this section starting with the laser source and then the electron source.
1.2.2.1 The Laser Effects

The chirped-pulse amplification (CPA) [44, 45] system is capable of producing an ultra-short and high-intensity laser pulse. The CPA laser system plays a crucial part in the PLEIADES Thomson backscattering experiment. Since real lasers deviate substantially from the ideal plane-wave model, we need a better theory of Thomson scattering. In this new model, “Bandwidth Contributions” from real lasers are included. “Bandwidth Contributions” refer to the laser frequency bandwidth and the laser focusing effect [39]. The magnitude of this focusing effect is determined by the laser angular spread $\Delta\theta$.

The relationship between the frequency bandwidth and the time-duration of a laser is given by the condition

$$\Delta\omega \Delta t \geq 2,$$

(1.18)

where both widths are the rms sizes of the intensity distributions. The peak frequency observed about the forward direction in the high $\gamma$ limit is

$$\omega_{\text{obs}} = \omega_{\text{inc}} \frac{2 \gamma^2}{1 + \gamma^2 \theta_{\text{obs}}^2} (1 + \beta \cos \theta_{\text{inc}}).$$

(1.19)

So, the Thomson scattered frequency bandwidth is the same as that of the incident laser

$$\frac{1}{\omega_{\text{obs}}} \frac{d\omega_{\text{obs}}}{d\omega_{\text{inc}}} = \frac{1}{\omega_{\text{inc}}} \Rightarrow \frac{\Delta\omega_{\text{obs}}}{\omega_{\text{obs}}} = \frac{\Delta\omega_{L}}{\omega_{L}}.$$

(1.20)

where $\omega_{L}$ is a laser frequency. In Eq. (1.20), we substitute $\Delta\omega = 2/\Delta t$ from Eq. (1.18). The laser frequency-bandwidth contribution found in the PLEIADES is around 2% with $\lambda = 820$ nm and $2\Delta t = 100$ fs.

The remaining bandwidth contribution comes from the laser’s focusing property. “Laser focusing” introduces a perpendicular $k$-vector components within
the laser pulse contributing to the laser’s non-plane-wave characteristic. The divergence behavior in a laser is represented by an effective laser beam emittance:

$$\varepsilon_L = \frac{\lambda_L}{\pi} = w_L w_L', \quad (1.21)$$

where $w_L$ is the transverse laser size and $w_L' = \Delta \theta_L = \Delta k_L / k$ is the width of the laser divergence. From Eq. (1.19), we obtain the spectrum broadening relation as

$$\frac{1}{\omega_{obs}} \frac{d\omega_{obs}}{d\theta_{inc}} = \frac{\beta \sin \theta_{inc}}{1 + \beta \cos \theta_{inc}} \Rightarrow \frac{\Delta \omega_{obs}}{\omega_{obs}} = \frac{\beta \sin \theta_{inc}}{1 + \beta \cos \theta_{inc}} \Delta \theta_{inc} \quad (1.22)$$

The equation shows that the bandwidth equation varies as a function of an incident angle $\theta_{inc}$. Now we need to derive the bandwidth equation in the special cases of the head-on and the side-on scattering geometry. First, in the relativistic limit, we take the approximation of $\beta \sim 1$, and expand the laser incident angle as $\theta_{inc} \rightarrow \theta_{inc} + \Delta \theta_L$.

In a head-on scattering geometry ($\theta_{inc} = 0$), the frequency bandwidth predicted is

$$\omega_{obs} = \omega_{inc} \frac{2\gamma^2}{1 + \gamma^2 \theta^2_{obs}} (1 + \cos \Delta \theta_L) \approx \omega_{inc} \frac{4\gamma^2}{1 + \gamma^2 \theta^2_{obs}} \left(1 - \frac{\Delta \theta^2_L}{4}\right)$$

$$\frac{\Delta \omega_{obs}}{\omega_{obs}} \approx \frac{\Delta \theta^2_L}{4} = \frac{1}{4} \left(\frac{\lambda_L}{\pi w_L}\right)^2 \quad (1.23)$$

A much more pronounced effect due to the perpendicular $k$-vector is expected from a side-on scattering geometry where the frequency bandwidth derived is

$$\omega_{obs} = \omega_{inc} \frac{2\gamma^2}{1 + \gamma^2 \theta^2_{obs}} [1 + \cos(\pi/2 + \Delta \theta_L)]$$

$$\approx \omega_{inc} \frac{4\gamma^2}{1 + \gamma^2 \theta^2_{obs}} (1 - \Delta \theta_L)$$

$$\frac{\Delta \omega_{obs}}{\omega_{obs}} \approx \Delta \theta_L = \frac{\lambda_L}{\pi w_L} \quad (1.24)$$
In any case, the spectral bandwidth due to laser focusing is negligibly small, and a plane-wave is generally a good approximation for our discussion.

1.2.2.2 The Electron-Beam Effects

The “collective electron effect” is very pronounced in x-ray bandwidth broadening as we shall see in this section. These bandwidth contributions are produced by the energy spread and the emittance of the electron beam [42, 43]. We will start with the discussion of the energy spread contribution to the total x-ray bandwidth.

From Eq. (1.20), the scattered frequency bandwidth observed in the ultra-relativistic limit is

\[ \frac{1}{\omega_{\text{obs}}} \frac{d\omega_{\text{obs}}}{d\gamma} = \frac{\Delta\omega_{\text{obs}}}{\omega_{\text{obs}}} = \frac{2}{1 + \gamma^2 \theta_{\text{obs}}^2} \Delta \gamma. \] (1.25)

The bandwidth equation varies with respect to the observation angle $\theta_{\text{obs}}$. The energy spread value of $\Delta \gamma / \gamma = 0.2\%$ is typically measured at PLEIADES and at $\theta_{\text{obs}} = 0$, a $0.4\%$ spectral bandwidth is obtained which is generally less than that of the laser bandwidth effect.

While the beam energy spread is found to be small, the beam emittance is potentially a significant contribution to spectral broadening. A rms geometric/stochastic emittance is defined by

\[ \sigma'_e \sigma_e = \varepsilon_g, \quad \sigma'_e = \frac{\Delta \theta_e}{2}, \] (1.26)

where $\sigma'_e$ is the rms divergence of the electron beam and $\sigma_e$ is the rms beam size at the beam waist. The angular spread introduces variations in the angle of laser incidence and the observation angle relative to the electron beam direction. Again from Eq. (1.20), the bandwidth contribution from the electron emittance is given by

\[ \frac{1}{\omega_{\text{obs}}} \frac{d\omega_{\text{obs}}}{d\theta_{\text{obs}}} = \frac{\Delta\omega_{\text{obs}}}{\omega_{\text{obs}}} = \frac{\gamma^2 \theta_{\text{obs}}}{1 + \gamma^2 \theta_{\text{obs}}^2} \Delta \theta_e. \] (1.27)
for small $\theta_{obs}$, the emittance bandwidth grows proportional to $\gamma^2$. For the special case of a head-on collision ($\theta_{obs} \rightarrow \Delta \theta_e / 2$), the on-axis broadening is given by

$$\frac{\Delta \omega_{obs}}{\omega_{obs}} \approx \frac{\gamma^2 (\Delta \theta_{ze}^2 + \Delta \theta_{ge}^2)}{2}.$$  \hspace{1cm} (1.28)

The electron emittance contribution with a normalized beam emittance (i.e. $\varepsilon/\beta \gamma$) of 5 mm-mrad, $\gamma=100$ and $\sigma=20\mu m$ is $\Delta \omega / \omega \simeq 12\%$, which is significantly larger than any other contributions to spectral broadening.

Real beams’ effects on the scattered photon spectrum has been calculated using the 3D Thomson simulation code [42] and is shown in Figure 1.6. The parameters of the incident laser and the electron beam used in the simulation are listed in the Tables 1.1. As shown in the figure, the spectrum in the case of the cold beam (no emittance) is sharply peaked and symmetrically distributed about
Figure 1.6: Inclusion of real beam effects such as energy dispersion & finite emittance have spectral broadening and reduction effects.

the Doppler-shifted frequency of 12.84 keV. In the case of the hot beam (a finite emittance), the spectrum has a gradually decreasing tail towards the low-energy range and a reduction in peak brightness. Also, we note that the peak frequency appears at a slightly lower value of 12.78 keV.

We conclude the discussion on real beam effects by giving the expression of the total contribution to spectral broadening. The total contribution from both laser and electron effects for the head-on collision can be summed up nicely in a
single equation as

\[
\frac{\Delta \omega}{\omega} = \sqrt{\frac{\Delta \omega^2_{L}}{\omega^2_{L}} + \left( \frac{\Delta \theta^2_{xL} + \Delta \theta^2_{yL}}{4} \right)^2 + \left( \frac{2 \Delta \gamma}{\gamma} \right)^2 + \left[ \frac{\gamma^2(\Delta \theta^2_{xe} + \Delta \theta^2_{ye})}{2} \right]^2}. \tag{1.29}
\]

From the real beam effect analysis, the main contribution to bandwidth broadening is from electron beam emittance.

### 1.2.3 The Classical Non-linear Thomson Scattering Theory

The Thomson scattering theory presented in the previous section was only valid in the linear regime defined by the laser potential parameter \( a_0 = eA_0/mc^2 \) being less than unity.

In order to optimize the rate of high x-ray production, we need to supply the highest number of scattering particles to a scattering point or to increase the scattering probability. The laser intensity expression, defined by \( I_0 [W/cm^2] = 2P_0/\pi w_0 \), shows that the incident photon density can be increased either through greater laser power or a smaller laser spot size. We can define the laser potential as a parameter dependent function of \( I_0 \): \( a_0 = 0.85 \times 10^{-9} \lambda_0 [\mu m] I_0^{1/2} [W/cm^2] \).

From this relation, a non-linear condition \((a_0 > 1)\) can be achieved at sufficiently high laser intensity \([46, 47]\). In this strong potential regime, interesting harmonic radiation features are produced from the non-linear electron motion \([48]\). In this section, We present the physics of non-linear Thomson scattering.

#### 1.2.3.1 The Non-linear Equation of Motion

The Thomson radiation is strictly described within the classical theory because the motion of electron is due to recoilless interaction with classical fields. The radiation energy distribution per unit solid angle per unit frequency can be obtained by solving the Lienard-Wiechert potentials. First, we need to determine
the electron’s equation of motion in the case of the linearly polarized laser field.

The electron motion in intense laser fields is governed by the laser vector potential. The laser vector potential can be parameterized by

$$a = a_0 \frac{(1 + p)^{1/2}}{\sqrt{2}} \cos \phi(t) \hat{x} + (1 - p)^{1/2} \sin \phi(t) \hat{y},$$

where \( \phi(t) = \omega [t - z(t)/c] \), and \( p \) is the polarization state. Linear polarization can be obtained from this equation by setting \( p = 1 \) (circular polarization is given by \( p = 0 \)). The rest of the theory presented in the section will concentrate on the linear polarization case.

The equation of motion in the laser is described by the relativistic Lorentz equation given in the form

$$\frac{1}{c} \frac{du}{dt} = \nabla \hat{\Phi} + \frac{1}{c} \frac{\partial a}{\partial t} - \vec{\beta} \times (\nabla \times a),$$

where \( u = \gamma \vec{\beta} \) is the normalized electron momentum. Note that from now on, we simplify the analysis by dropping the space-charge field (\( \hat{\Phi} \)) of the electron beam from the equation of motion. In terms of the actual field definitions (\( E \) and \( B \)), the relativistic equation of motions can be given by

$$\frac{dP}{dt} = q[E(r, t) + v \times B(r, t)],$$

$$\frac{dE}{dt} = qv \cdot E(r, t),$$

where \( P = \gamma m v \) and \( E = \gamma mc^2 \). The electromagnetic fields of a plane-wave traveling in \( z \) direction polarized along \( \hat{x} \) are given by \( E(r, t) = \hat{x} E_0 \sin \phi(t) \) and \( B(r, t) = \hat{y} E_0 \sin \phi(t)/c \). These EM fields can be elegantly described by the laser vector potential, \( a = a_0 \cos \phi \hat{x} \). We also will not treat the radiation damping or the ponderomotive effect.

We note that Thomson scattering is exactly equivalent of the undulator radiation (see Figure 1.7), except for the fact that the scale of the field wavelength
Figure 1.7: The comparisons between FEL (top) & Thomson scattering (bottom). [Courtesy of D. J. Gibson]

is vastly different. For instance, $\sim 1\mu m$ incident plane-wave wavelength in the Thomson scattering and the undulator period of typically around a centimeter long produces a difference of a factor $10^3$. As shown in Figure 1.7 (top), the undulator field in the lab frame resembles an alternating magnetic field. After the Lorentz transformation of the undulator field from the lab to the electron’s rest frame, an additional electric field component arises. This can be demonstrated mathematically as follows. The Lorentz transformation undulator fields in the electron rest frame are given by

$$E' = \gamma(E + \vec{\beta} \times B) - \frac{\gamma^2}{\gamma + 1} \vec{\beta} (\vec{\beta} \cdot E),$$
$$B' = \gamma(B - \vec{\beta} \times E) - \frac{\gamma^2}{\gamma + 1} \vec{\beta} (\vec{\beta} \cdot B),$$

(1.33)

where the unprimed variables are the lab frame undulator fields given by $E = 0$ and $B = \hat{y}B_u \cos(2\pi\gamma z/\lambda_u)$. After the substitution of these lab frame fields into
Eq. (1.33), we have

\[
E' = \gamma \beta (\hat{n} \times B),
\]

\[
B' = \gamma B, \tag{1.34}
\]

where \( \hat{n} \equiv \hat{z} \). As \( \beta \to 1 \), the undulator field appears like a propagating EM plane-wave with a wavelength of \( \lambda_u = 2\pi c\gamma / \omega_u \) and an amplitude of \( E'_0 = \gamma B_0 \).

In this new frame, an electron is subjected to both magnetic and electric forces.

We use an unorthodox approach by applying the undulator equation of motion to derive the Thomson theory of non-linear electron motion. Then, we attempt to draw an equivalence between the two theories with the assumption of \( K \to a_0 \) [50].

The equations of electron motion, as defined in Eq. (1.32), in the lab frame of an undulator are given by

\[
\ddot{x} = -\frac{eB}{\gamma m} \cos(k_u z) \dot{z}, \tag{1.35}
\]

\[
\ddot{z} = \frac{eB}{\gamma m} \cos(k_u z) \dot{x}. \tag{1.36}
\]

By solving these differential equations, we obtained the time-dependent coordinate functions in the lab frame as:

\[
x(t) = \frac{K}{\beta \gamma k_u} \cos(\omega_u t), \tag{1.37}
\]

\[
z(t) = \beta^* c t + \frac{K^2}{8\beta^2 \gamma^2 k_u} \sin(2\omega_u t), \tag{1.38}
\]

where the average velocities are defined by \( \langle \beta_x \rangle = 0 \) and \( \langle \beta_z \rangle = \beta^* c \), and \( K = eB/mck_u \). We can then obtain the equation of motion in the electron rest frame moving at a drift velocity of \( \beta^* c \). Using the Lorentz transformation

\[
x' = x \quad z' = \gamma^*(z - \beta^* c t), \quad t' = \gamma^* \left( t - \frac{\beta^* z}{c} \right) \tag{1.39}
\]
Figure 1.8: The non-linear “figure-8” motions for different laser vector potentials are plotted.

where $\gamma^* = 1/\sqrt{1 - \beta^2} \sim \gamma/\sqrt{1 + K^2/2}$. Finally, the electron trajectories in the rest frame are given by

$$x'(t) = \frac{K}{\beta \gamma k_u} \cos(\omega_u t),$$  \hspace{1cm} (1.40)

$$z'(t) = \frac{K^2}{8 \beta^2 \gamma \sqrt{1 + K^2/2}} \sin(2\omega_u t)$$  \hspace{1cm} (1.41)

Combining the two equations, a hallmark “figure-8” trajectory (see Figure 1.8) is produced from the expression given by

$$z'^2(x') = \frac{K^2}{16 \beta^2 (1 + K^2/2)} x'^2 \left( 1 - \frac{x'^2}{(K/\beta \gamma k_u)^2} \right)$$  \hspace{1cm} (1.42)
Substitution of \( E = \gamma B \) and \( \omega = 2\pi c\gamma/\lambda_u \) in the limit \( v \to c \) into the definition of \( a_0 \), we obtain the important relation

\[
a_0 = \frac{e\gamma B}{mc(2\pi c\gamma/\lambda_u)} = \frac{eB}{mck_u} = K
\]  

(1.43)

Exactly equal to the undulator strength parameter \( K \). Therefore, using the undulator results (Eq. (1.42)) we can show that the electron’s motion in a strong laser field limit \( (a_0 > 1) \) is a “figure-8” pattern. The electron’s motion in its rest frame is calculated for various values of \( 0.1 \leq a_0 \leq 1.0 \) (see Figure 1.8). Note that for \( a_0 = 0.1 \) there is a small dipole motion along the direction of the \( E \)-field. On the other hand, for the \( a_0 = 1.0 \) case, the \( B \)-field has taken the full non-linear effect on the motion.

A small remark regarding Thomson scattering involving circular laser polarization. Scattering from the circularly polarized laser vector field induces circular helix motions where the circular motion is contained in the \( x-y \) plane and the average velocity is in the \( z \)-direction. This circular motion about the axial direction produces a symmetric radiation profile. Again, an equivalent undulator radiation description exists in the helical undulator.

### 1.2.3.2 The Harmonic Generation & Radiation Energy Depression

In the last section, the non-linear motion of an electron in a strong laser field was examined in the context of undulator radiation theory. The non-linear “figure-8” motion of an electron is produced when the laser vector potential satisfies the condition, \( a_0 > 1 \). This condition is generated either by sufficiently long laser wavelength or by high laser intensity \( (a_0 = 0.85 \times 10^{-9}\lambda_0 I_0^{1/2}) \). In this condition, trademark high-intensity laser effects such as the Doppler-frequency shift and the generation of harmonic radiations in the high-frequency domain are observed; the effects are not present in linear Thomson scattering.
In the previous section, we only examined the non-linear electron motion by borrowing from the motion in an undulator. To continue with the analysis of the non-linear spectrum distribution, we will need to solve the Lienard-Wiechert potential of the Thomson scattering process. The analysis of the non-linear effects are largely drawn from the paper by Ride [49] and Esarey [47]. The general theory developed by Esarey includes arbitrary scattering angle and incident polarization state effects. We will later show in the present section that non-linear electron motion produces a very interesting energy distribution feature over a range of an observation angle and a scattered photon spectrum. In the present thesis, we examine a special case of Esarey’s complete theory: a linearly polarized incident laser scattering a head-on electron beam.

We now derive the energy spectrum distribution of Thomson scattering of a plane-wave. The Lienard-Wiechert intensity distribution is

\[
\left. \frac{d^2 I}{d\omega d\Omega} \right| \frac{c^2 \omega^2}{4\pi^2 c} \left| \int_{-T}^{T} n \times (n \times \vec{\beta}) e^{i\omega(t-n \cdot r(t)/c)} dt \right|^2 , \quad (1.44)
\]

where \( T \) is the interaction time. In the intensity distribution, we used the spherical coordinate system \((r, \theta, \phi)\) and the unit vectors \((e_r, e_\theta, e_\phi)\), where \( x = r \sin \theta \cos \phi \), \( y = r \sin \theta \sin \phi \), and \( z = r \cos \theta \). We define the observation direction of the scattered radiation to be \( n = e_r \). The vector product and the scalar multiplication found in the integrand of the Eq. (1.44) are given by

\[
n \times (n \times \vec{\beta}) = - (\beta_x \cos \theta \cos \phi + \beta_y \cos \theta \sin \phi - \beta_z \sin \theta) e_\theta \\
+ (\beta_z \sin \phi - \beta_y \cos \phi) e_\phi,
\]

\[
n \cdot r = x \sin \theta \cos \phi + y \sin \theta \sin \phi + z \cos \theta . \quad (1.45)
\]

The vector product term dictates the polarization state of the scattered light. The product divides into two polarization terms, \( e_\theta \) and \( e_\phi \). A scattered photon could assume one or the other state depending on the initial incident polarization. The
Lienard-Wiechert intensity distribution then decomposes in these polarization states according to \( I = I_\theta + I_\phi \), where the respective expressions are

\[
\frac{d^2 I_\theta}{d\omega d\Omega} = \frac{e^2 k^2}{4\pi^2 c} \left| \int_{-\eta_0}^{\eta_0} d\eta \left[ \frac{dx}{d\eta} \cos \theta \cos \phi + \frac{dy}{d\eta} \cos \theta \sin \phi - \frac{dz}{d\eta} \sin \theta \right] \exp(i\psi) \right|^2,
\]

\[
\frac{d^2 I_\phi}{d\omega d\Omega} = \frac{e^2 k^2}{4\pi^2 c} \left| \int_{-\eta_0}^{\eta_0} d\eta \left[ \frac{dx}{d\eta} \sin \phi - \frac{dy}{d\eta} \cos \phi \right] \exp(i\psi) \right|^2,
\]

where \( k = \omega/c \), \( \eta_0 = L_0/2 \), and \( L_0 = c(1 + \beta^*)T \) is the laser interaction length and

\[
\psi = k[\eta - z(1 + \cos \theta) - x \sin \theta \cos \phi - y \sin \theta \sin \phi].
\]

The variable \( t \) in the Lienard-Wiechert integral has been replaced with a general definition of \( c\beta dt = (d\mathbf{r}/d\eta)d\eta \). The non-linear motion of an electron in the case of a linear polarization in the \( x \)-direction is given by

\[
\begin{align*}
x(\eta) &= x_0 + r_1 \sin k_0 \eta, \\
y(\eta) &= y_0, \\
z(\eta) &= z_0 + \beta_1 \eta + z_1 \sin 2k_0 \eta,
\end{align*}
\]

where we grouped coefficient terms in the new variables by defining

\[
\begin{align*}
r_1 &= \frac{a_0}{\gamma_0(1 + \beta_0)k_0}, \\
z_1 &= -\frac{a_0^2}{8\gamma_0^2(1 + \beta_0)^2k_0}, \\
\beta_1 &= \frac{1 - (1 + a_0^2)/\gamma_0^2(1 + \beta_0)^2}{2}
\end{align*}
\]

The analytical solution of the integral is obtained by first converting the phase factor using the Bessel identity of

\[
\exp(ib\sin a) = \sum_{n=-\infty}^{\infty} J_n(b) \exp(ina)
\]

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With the identity, the phase factor is

$$\exp[i(\psi + l k_0 \eta)] = \sum_{m,n=-\infty}^{\infty} J_m(a_z) J_{n-2m+1}(a_x) \exp[i(\psi_0 + k_n \eta)],$$  \hspace{1cm} (1.50)$$

where

$$a_z = k z_1 (1 + \cos \theta),$$
$$a_x = k r_1 \sin \theta \cos \phi,$$
$$k_n = k [1 - \beta_1 (1 + \cos \theta)] - n k_0.$$  \hspace{1cm} (1.51)$$

The integrals in Eq. (1.46) are evaluated with the identities in Eq. (1.49) & (1.50), along with the electron equation of motion, Eq. (1.49):

$$I_x = k_0 r_1 e^{i\psi_0} \sum_{n,m=-\infty}^{\infty} \frac{\sin k_n \eta_0}{k_n} J_m(a_z) [J_{n-2m-1}(a_x) + J_{n-2m+1}(a_x)],$$
$$I_y = 0,$$
$$I_z = 2 e^{i\psi_0} \sum_{n,m=-\infty}^{\infty} \frac{\sin k_n \eta_0}{k_n \eta_0} \times$$
$$k_n J_m(a_z) \{ \beta_1 J_{n-2m}(a_z) + k_0 z_1 [J_{n-2m-2}(a_x) + J_{n-2m+2}(a_x)] \}.$$  \hspace{1cm} (1.52)$$

Note that integral results are grouped accordingly with respect to \(x, y\) and \(z\) components in Eq. (1.46).

We note that Bessel functions of higher orders have a negligible contribution to the total summation. This is because the argument of the function is small such that the summation converges quickly to a stable value at lower orders. In evaluating the summation, it is sufficient to sum up to a small finite \(m\) value.

Combining all evaluated integrals in the Lienard-Wiechert equation, the final result is given by

$$\frac{d^2 I}{d\omega d\Omega} = \frac{e^2 \omega^2}{4 \pi^2 e^3} \left( |I_x \cos \theta \cos \phi - I_z \sin \theta|^2 + |I_x \sin \phi|^2 \right).$$  \hspace{1cm} (1.53)$$

In the equation, the higher harmonics are generated as resonant frequencies \(\omega_n\).
Figure 1.9: The resonant function showing $n = 1, 2, 3$ resonance.

(a “square” plane-wave) of the resonance function $R(k, nk_0)$, where

$$R(k, nk_0) = \left(\frac{\sin k_n\eta_0}{k_n\eta_0}\right)^2,$$  \hspace{1cm} (1.54)

and

$$\omega_n = \frac{n\omega_0}{1 - \beta_1 (1 + \cos \theta)}.$$  \hspace{1cm} (1.55)

The resonant frequency is downshifted by the term $\beta_1$ in the denominator.

We notice that the denominator of a resonance has a dependence on $\theta$. In the plot of the resonance function in the $\phi = 0$ plane, resonant peaks start at $n\omega_0/(1 - 2\beta_1)$ and trails off towards the lower spectrum limit of $n\omega_0/(1 - \beta_1)$. Also, the amplitude of the resonant function is independent of all angular range
Figure 1.10: The linear Thomson shows only the fundamental frequency (blue). The non-linear harmonics \( n = 1, 2, 3, 4 \) are shown (red). The first harmonic is frequency down shifted.

of \( \theta \). The plot of the resonant function in Eq. (1.54) is shown in Figure 1.9.

At \( \theta = 0 \), the scattered intensity in “weak” \( (a_0 = 0.1) \) and “strong” \( (a_0 = 1.0) \) cases are shown in Figure 1.10. The energy depression or the Doppler down-shift in harmonic frequencies is noted in the resonant energy for case \( a_0 = 1.0 \). As for the characteristic of the amplitude of the resonant spectrum at off-axis-angle \( (\theta \neq 0) \) is governed by the Bessel functions in the Eq. (1.53). For \( n = 1 \) in Eq. (1.53), the first harmonic has a peak on-axis and reduces in amplitude as the angle \( \theta \) increases. Similar traits are observed in the even harmonics, but showing more separate peaks off center axis. Also, for the odd harmonics, the peaks of the resonance are located away from the \( \theta = 0 \). We will examine two special cases with \( a_0 = 0.1 \) and \( a_0 = 1 \) in the next discussion.
Figure 1.11: The intensity distribution plotted over normalized parameters $\omega/4\gamma_0^2 \omega_0$ and $\gamma_0 \theta$. (top) for $a_0 = 0.1$ shows the fundamental frequency only & (bottom) for $a_0 = 1.0$ shows high harmonics.
The intensity distribution as functions of normalized frequency $\omega/4\gamma_0^2\omega_0$ and normalized observation angle $\gamma_0\theta$ is plotted for the case of the “weak” laser field ($a_0 = 0.1$) and the “strong” laser field ($a_0 = 1$) in Figure 1.11. The scattering geometry is the head-on collision between a relativistic electron beam and a linearly polarized laser. In the “weak” case, the fundamental harmonic radiation ($n = 1$) is at a maximum near $\omega/4\gamma_0^2\omega_0 = 1$ and trails off towards lower frequencies as the observation angle increases. In this linear Thomson scattering case, no high harmonics nor energy down-shift is observed. On the other hand, for the “strong” laser case, harmonics up to $n = 3$ are shown in the plot. The fundamental scattered frequency exhibits the frequency down-shift behavior as predicted from the theory when $a_0 > 1$. Also, the second harmonic, which is an even harmonic, is peaked away from $\theta = 0$. The third harmonic is peaked at $\theta = 0$. Again, the intensity of the fundamental frequency is the dominant over all harmonics.

The frequency width observed in the harmonic radiation in the case of the relativistic electron beam can be estimated with the resonant harmonic frequency $\omega_n$ of Eq. (1.54). The frequency width $\Delta\omega$ with respect to an angular width $\Delta\theta$ is written as

$$\frac{|\Delta\omega|}{\omega_n} \approx \frac{|M_0(\theta\Delta\theta + \Delta\theta^2/2)|}{2 + M_0\theta^2/2},$$

where in a relativistic regime, $M_0 \approx 4\gamma_0^2/\gamma_\perp^2$ and $\gamma_\perp = (1 + a_0^2/2)^{-1/2}$. From Eq. (1.56), the angular width with a center at $\theta = 0$ for the linearly polarized laser is given by

$$\Delta\theta \approx \frac{\gamma_\perp}{\gamma_0} \sqrt{\frac{\Delta\omega}{\omega_n}}.$$  \hspace{1cm} (1.57)

We let frequency line-width in Eq. (1.57) to be the intrinsic frequency line-width $\Delta\omega/\omega = 1/nN_0$, where $N_0 = L_0/\lambda_0$. So, the angular width of a single harmonic $n$, using the intrinsic width in Eq. (1.56), is written as

$$\Delta\theta \approx \frac{\gamma_\perp}{\gamma_0} \frac{1}{\sqrt{nN_0}}.$$  \hspace{1cm} (1.58)
We note that the angular width narrows as the harmonic number \( n \) increases and as the electron beam energy \( \gamma_0 \) increases.

In an ultra-intense laser beam, \( a_0 \gg 1 \), the numerous harmonic radiations that are generated by the beam start to overlap together essentially resembles the synchrotron spectrum. This synchrotron spectrum-like behavior is further enhanced when the spectrum broadening effect from the finite electron energy spread is included. So, in the ultra-intense limit, the non-linear Thomson spectrum appears like a broadband spectrum which extends out to a critical frequency \( \omega_c = n_c M_0 \omega_0 \) where the critical harmonic number is given by \( n_c \simeq 3a_0^3/4 \). The spectrum intensity beyond the critical frequency falls at an exponential rate. The spectrum distribution characteristics for \( \phi = \pi/2, a_0 \gg 1, n \gg 1 \) and \( \theta^2 \ll 1 \) can be described by [47]

\[
\frac{d^2 I}{d\omega d\Omega} \simeq N_0 \frac{12c^2}{\pi^2c} \frac{\hat{\gamma}^2 \xi^2}{1 + \hat{\gamma}^2 \theta^2} \left[ \frac{\hat{\gamma}^2 \xi^2}{1 + \hat{\gamma}^2 \theta^2} K_{1/3}^2(\xi) + K_{2/3}^2(\xi) \right],
\]

where

\[
\xi = \frac{\omega}{\omega_c} (1 + \hat{\gamma}^2 \theta^2)^{3/2},
\]
\[
\hat{\gamma} = h_0/2.
\]

We note that emission in the vertical direction is confined to the angle \( \theta_v \sim 1/h_0 \). In the horizontal direction, the angle is \( \theta_h \sim a_0/h_0 \). The opening angle of scattered emission in this limit resembles synchrotron light strongly polarized in the horizontal direction.

We will end the current section by giving a summary of Thomson scattering in the linear and non-linear regimes.

In the weak laser field limit, \( a_0 < 1 \), the electron motion in the lab frame will simply be that of a sinusoidal motion with a frequency of \( \omega = 2\pi c/\lambda \) where \( \lambda \) is the laser wavelength. In the moving frame (i.e. in the electron’s rest frame), this
motion is a simple harmonic oscillation with upshift frequency of $\omega' = \gamma^* \omega$ and a Lorentz contracted period, $\lambda' = \lambda / \gamma$. A simple harmonic motion in this frame creates the well known dipole radiation emitted in both $z$-directions. Upon Lorentz transformation back to the lab frame, the emitted radiation is tightly concentrated in a forward direction of electron’s average velocity $\beta^* c$ with a radiation opening angle of $\theta \sim 1 / \gamma$. The back-scattered radiation frequency in a weak field regime is given by $\omega = 4 \gamma^2 \omega_0 / (1 + \gamma^2 \theta^2)$. We note that the $\gamma \theta$ contributes the spectrum bandwidth up to the maximum angle of $\sim 1 / \gamma$ in the lab frame.

In the strong laser field limit, $a_0 > 1$, the electron trajectory no longer follows a simple harmonic motion; instead, in the case of a linearly polarized incident laser field, a “figure-8” motion is produced. This motion can be decomposed into harmonic oscillations along the $x$-direction producing odd harmonics of frequency $\omega'$, and longitudinal harmonic motions along the $z$-direction with even harmonics of $2\omega'$. Again, relativistic effects confines odd harmonic radiation to a narrow angle of $1 / \gamma$ and concentrate even harmonics on a ring around the odd harmonics. The scattered frequency observed at arbitrary angle $\theta$ has a dependency on $a_0$ according to: $\omega = 4 \gamma^2 \omega_0 / (1 + a_0^2 / 2 + \gamma^2 \theta^2)$. With increasing laser strength, the energy spectrum depression is expected to become more pronounced with a broader spectrum containing many peaks.
CHAPTER 2

PLEIADES Experiment

The PLEIADES experiment - a head-on collision Thomson x-ray source - was built at Lawrence Livermore National Laboratory with the intent of producing a tunable angstrom scale wavelength and a sub-picosecond range time duration from a head-on collision of a TW FALCON laser [51, 52] with a high-relativistic electron beam. The tunable femto-second x-rays produced at the facility promises to be fully capable of probing matters at the atomic scale. One of the initial uses of this x-ray source is the study of the behavior of various ultra-fast dynamic phenomena in high-Z materials. The hard x-ray energy spectrum in the 20-140 keV range will be specifically used to investigate K-edge energy features in heavy metals. These heavy metals are tantalum (67.5 keV), gold (80.7 keV), bismuth (90.54 keV) and uranium (115.0 keV) [14, 53].

The success of this experiment is largely dependent on the production of a high-quality, high-brightness electron beam in 20-120 MeV energy range. The beamline system was designed to provide all necessary beam diagnostics systems, and the optimal operation necessary to minimize the charge-loss during the beam transportation. Time synchronization was also factored into the beamline design in order to guarantee the absolute temporal beam alignment at the interaction region. In the PLEIADES system, an existing 100 MeV linear accelerator [54], with the rf photoinjector as a high-brightness beam source [55], and the dual purpose laser system [1] are integrated to make up a highly efficient hard
femto-second x-ray source. A spatial-temporal overlap in this x-ray machine has
achieved better than 2 ps separation and the 20 µm necessary to produce 10^6-10^8
Thomson x-ray photons.

In the first half of this chapter we will review the low-energy electron beam
linear accelerator system. Then, we will discuss the PLEIADES laser system of
the PLS (rf photoinjector) and the FALCON (high-intensity laser) lasers.

2.1 Electron beam source

A high-quality, high-brightness electron beam plays a very important role in the
PLEIADES experiment. This is because the number of x-rays scattered from the
electron beam depends on the phase-space distribution that an electron beam is in
at the moment of the interaction event. In general, the phase space distribution is
represented by the beam emittance, which plays a significant role in the spectral
broadening and the intensity reduction effect as we discussed in the last chapter.
A low emittance beam is therefore highly desirable in this application for the
higher x-ray flux generation. This is straightforward, as the denser region of the
descriibed region of the
particle beam creates a higher probability that scattering events will occur. If a
low emittance is hard to obtain in the beam from the electron source, an ultra-
short focal length final focus system can be employed to produce a small beam
(independent of beam emittance issue) to maximize the electron density.

A low emittance electron beam is desirable for an improved brightness of
the x-ray source. In many respects, the electron beam emittance is found to be
the cornerstone of the high-brightness x-ray beam generation. Before detailed
review of the complete PLEIADES system is possible, an understanding of this
emittance is needed. An in depth cover of the accelerator physics is found in
2.1.1 Collective Charged Particle Phase Space Distribution

Successful production of a high-brightness x-ray beam is measured by the electron beam property at the Thomson interaction region. Transverse size, angular divergence and longitudinal length are parameters of the electron beam that determines this output of the Thomson x-ray brightness. These parameters, which characterize the quality of an electron beam, are described by phase space distribution. To understand the property of an electron beam and the dynamics in the accelerator, we will review the basic concept of motion in phase space.

A collective characteristic of an ensemble of \( N \) charged particles is represented by a six-dimensional phase space distribution in \( x \) and \( p \), where each particle has three spatial coordinates \( x_i \) and three momentum components \( p_i \). An \( N \) charged particle distribution in the phase space is given by

\[
    f(x, p) = \sum_{i=1}^{N} \delta^{(3)}(x - x_i)\delta^{(3)}(p - p_i),
\]

where a sharp localization of a particle \( i \) at \( (x_i, p_i) \) coordinates is represented by the product of six Dirac delta functions. We will approximate this distribution with a smooth, continuous distribution. In the smooth distribution regime, an assumption is that many particles are found in a small volume of \( dx dp \). The number of particles within a small volume is \( dN = f(x, p)dx dp \) and the total number of particles occupied in a phase space is

\[
    N = \int f(x, p)dx dp,
\]

where a typical charged particle number of \( N \sim 10^9 \) or higher are found in a beam.
To understand the phase space distribution’s dynamic evolution in the beam-line, we need the Vlasov equation. To derive this Vlasov equation, we start from the fact that the total time derivative of distribution function \( f(x(t), p(t), t) \) is equal to zero, \( df/dt = 0 \). This is Liouville’s theorem, which states that while the particle distribution evolves with time, the density of particles in a volume of a phase space remains invariant of time \( t \). Then, from the Liouville’s theorem, we obtained the Vlasov equation:

\[
\frac{df}{dt} = \left( \frac{\partial}{\partial t} + \dot{x} \cdot \nabla_x + \dot{p} \cdot \nabla_p \right) f = 0.
\] (2.3)

where the divergence operators are differentiations with respect to \( x \) and \( p \).

To simplify the description of a phase space function, we restrict to a case of decoupled particle motions such that a separate distribution function is obtained \( f = f_x f_y f_z \). This decoupled distribution model closely approximates the real beam model. For this mathematical simplification to work, we must assume that there is no correlation between the dimensions. To consider the simplest case, we assume the special case of an equilibrium state, i.e., the partial time derivative term is dropped from Eq. (2.3):

\[
x \frac{df_x}{dx} + p_x \frac{df_x}{dp_x} = 0,
\] (2.4)

where \( p_x = F_x \) is a linear force in \( x \). If we assume that \( f_x \) is separable so that we may write \( f_x(x, p_x) = X(x)P(p_x) \), and furthermore that the bi-Gaussian distribution function is self-consistent with Eq. (2.4), we obtained the following expression

\[
f_x(x, p_x) = f_0 \exp \left( -\frac{x^2}{2\sigma_x^2} \right) \exp \left( -\frac{p_x^2}{2\sigma_{p_x}^2} \right),
\] (2.5)

where \( f_0 \) is a normalization factor equal to the total charged particle number, \( N \). The \( \sigma_x \) is the distance covered by the restoring force and the \( \sigma_{p_x} \) is the momentum range of the kinetic energy in the paraxial approximation.
\[ \sigma_x = 0 \quad \sigma_x' > 0 \quad \sigma_x > 0 \quad \sigma_x' = 0 \]

Figure 2.1: Idealized electron beam sources with zero emittance. [courtesy of M. C. Thompson]

Since the total time derivative of the distribution is zero, the phase space area remains independent of time. The boundary curve of a distribution function is generally expressed by an ellipsoidal equation:

\[ \frac{x^2}{\sigma_x^2} + \frac{p_x^2}{\sigma_{p_x}^2} = \text{constant.} \]  \hspace{1cm} (2.6)

We note that this ellipsoidal equation is equivalent to a Hamiltonian for a harmonic oscillator, because by recognitions of \( x^2 \propto V \) and \( p_x^2 \propto T \).

The total statistical information of all particle momenta and positions at any position along the accelerator beamline is not static, then in that case we reinstate the partial time derivative term in Eq. (2.3). We can then describe the distribution changes in orientation in response to the different electromagnetic forces it encounters in the beamline. But the total size of the phase space distribution remains constant in linear optics. This invariant size can be identified
with the emittance. There are two situations where an emittance is exactly zero: a dimensionless point beam with a finite divergence and a finite beam size with a parallel beam. These extreme cases are shown in Figure 2.1. However, in any real beams, the emittance is always non-zero ($\varepsilon > 0$).

2.1.2 The Electron Beam Parameters

A mathematical method of evaluating the collective behavior of $N$ particles in a beamline (of either a circular or a linear type) has satisfactorily produced solutions to most standard problems. The Vlasov equation for a phase space distribution is the first step to understanding the beam-dynamics in the accelerator. Other parameters derived as solutions of this method provide important information regarding the properties of a beam in any beamline configuration. These parameters are well known as Twiss parameters. Most important of all is the area enclosed by the outer boundary of a phase space distribution. The phase space area known as the emittance in general measures the quality of the charged particle beam.

The starting point is redefining the variables of phase-space. The variables in the trace space $(x, x')$ will be used instead of $(x, p)$, in accordance with the standard practice. In a paraxial approximation, we replace the $p_x$ with the slope of the particle trajectory $dx/dz$. Likewise, a phase space distribution $f(x, p_x)$ is replaced by a trace space distribution $f(x, x')$.

The second moments of the beam’s trace space distribution function can be formally calculated from the moment integrals

$$
\sigma_x^2 = \langle x^2 \rangle = \int \int x^2 f(x, x') dx dx',
$$

$$
\sigma'_x^2 = \langle x'^2 \rangle = \int \int x'^2 f(x, x') dx dx',
$$

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Figure 2.2: The phase space of charged particle distribution.

\[ \sigma_{xx'}^2 = \langle xx' \rangle = \int \int xx' f(x, x') dx dx'. \]  

(2.7)

The first integral describes the square of an rms transverse beam dimension. The second integral describes the square of an rms beam divergence. The last integral represents the correlation between \( x \) and \( x' \). The elliptical area of the trace space distribution with arbitrary alignment in \( x-x' \) plane is:

\[ \gamma_x(s)x^2 + 2\alpha_x(s)xx' + \beta_x(s)x'^2 = \varepsilon_x, \]  

(2.8)

where \( \beta_x, \alpha_x \) and \( \gamma_x \) are Twiss parameters and \( \varepsilon_x \) is a trace space area (see Figure 2.2). Further simplification in the description of orientation of an ellipse in trace space can be accomplished by considering that the invariant \( \varepsilon_x \) is redundant. Then, an ellipse can be parameterized by three parameters - \( \beta_x, \gamma_x \) and \( \alpha_x \) - and related by \( \beta_x \gamma_x - \alpha_x^2 = 1 \). These parameters expressed in terms of the second
moments are given by
\[
\begin{align*}
\gamma_x &= \frac{\langle x'^2 \rangle}{\varepsilon_x}, \\
\beta_x &= \frac{\langle x^2 \rangle}{\varepsilon_x}, \\
\alpha_x &= -\frac{\langle xx' \rangle}{\varepsilon_x}.
\end{align*}
\] (2.9)

And the area of ellipse or the rms emittance is given by
\[
\varepsilon_x = \sqrt{\langle x^2 \rangle \langle x'^2 \rangle - \langle xx' \rangle^2}.
\] (2.10)

We note that $\varepsilon_x$ is not truly an invariant of a beam energy, because the angle with respect to the $z$-axis changes with changes in the relativistic factor term in the denominator $p_x/\beta \gamma mc$. Thus, the rms emittance varies with respect to the beam energy. A normalized emittance, which is invariant with respect to the beam energy, is given by
\[
\varepsilon_{x,n} = \beta \gamma \varepsilon_x.
\] (2.11)

We will always adopt this terminology unless specify otherwise in this dissertation. Phase space parameter values at any given location in an accelerator beamline can be obtained by different mathematical formalisms. The result is fundamentally equivalent, independent of which method used. A solution can be obtained by either a transport matrix or a rms envelope equation method.

The differential equations of Twiss parameters will be used to describe the dynamics of a beam. The solutions are obtained with a given initial condition of Twiss parameters. The equations describing the evolution of Twiss parameters are given by
\[
\begin{align*}
\beta_x' &= -2\alpha_x, \\
\alpha_x' &= -\gamma_x + \beta_x k_x^2, \\
\gamma_x' &= 2k_x^2 \alpha_x,
\end{align*}
\] (2.12)
where $k_x$ is a constant of lattice element (like “spring constant”). A well known set of solutions of these equations can be studied. In the force-free drift ($k_x = 0$), a transverse beam as a function of a longitudinal parameter $s$ is obtained from these differential Twiss parameter equations.

The transverse beam size as a function of a beam trajectory is expressed by

$$
\sigma_x^2(s) = \varepsilon_x \beta_x^* \left[ 1 + \left( \frac{s - s_0}{\beta_x^*} \right)^2 \right].
$$

(2.13)

where $\beta_x^*$ is the minimum $\beta$-function at the waist of the funnel as shown in Figure 2.3. We note the trace space distribution evolves in a clockwise direction as the beam is being focused to a small waist and later expands outward after the focal point. In Chapter 4 we will revisit this expression associated with the ultra-high gradient final focusing system and a small electron beam.

Figure 2.3: The phase space dynamics near the focal region.
2.1.3 Electron Beam Injectors

An electron beam source is obtained when bound-state electrons in a metal gain enough energy to escape the trapping potential well. This energy is gained from either a thermal heating or a photo-emission process. Based on these principles, there are two different types of electron beam source technology: DC thermionic guns (thermal emission process) and rf photo-injector guns (photo-electric emission process). The thermionic gun can practically produce more electrons, while the photoinjector gun produces a higher beam brightness which is of a great interest to modern advanced accelerator research. In particular the rf photoinjector system is found to be a high quality, high brightness electron beam source for the production of high-flux, short pulse Thomson x-rays. In this section, we first examine the working principle of the thermionic gun and its limitations [58]. Then, how the photoinjector overcomes those limitations.

2.1.3.1 DC Thermionic Guns

The thermionic DC gun is known as a “thermal-emission” electron source. A main component of a thermionic injector is a cathode, which is heated to a high temperature $T$. From the heated cathode surface, free electrons are released. These free electrons are then instantly accelerated by a high electric field established between the cathode and the anode. This high electric field is generated with a high-voltage power supply producing a high-voltage drop of $V$ between a cathode and an anode plate. When the electron beam reaches its peak acceleration near the end of the gun structure, it exits through a hole placed in the anode. A design of electrode plates is used which are slightly curved, so that a radial focusing field is produced which counteracts the space-charge expansion of the beam.
But the thermionic gun is known for low-brightness beams inapplicable for experiments requiring high-brightness beams such as LCLS x-ray FEL. The limitation of a thermionic gun is that the amount of current produced is limited by the intrinsic thermal properties of the emitter. The other limiting effect relates to the Child-Langmuir effect [59]. In the thermionic gun, an electron gas just liberated from the hot metal surface generates an opposing electric field strong enough to partially cancel the accelerating electric field in the gun. The beam current then grows until a balance between the beam electric field and the cathode electric field is reached. Current growth is then completely stopped beyond this point. The point of this balance is given by

\[ I = KV^{3/2}, \]  

(2.14)

where \( K \) is the beam perveance and \( V \) is the voltage. The maximum operational
voltage is limited by the high-voltage breakdown of the cathode that normally occurs around 400-500 keV [58].

The beam brightness achieved in a thermionic gun already operating at the Child-Langmuir limit is still orders of magnitude lower than the brightness required by most advanced accelerator applications. In the thermionic gun, the emittance produced is governed by the relation

$$\varepsilon_n = \frac{r_c}{2} \sqrt{\frac{k_BT}{mc^2}},$$  \hspace{1cm} (2.15)

where $r_c$ is the radius of the source beam at the cathode surface and $k_BT$ is the thermal energy used in electron emissions. We assume that the radius of the source beam of a constant current density across the emitting surface is given by $r_c = \sqrt{I/\pi J}$. If so then, the transverse beam brightness is given by

$$B_\perp = \frac{I}{\varepsilon_n,x \varepsilon_n,y}.$$ \hspace{1cm} (2.16)

With the substitution of the emittance defined, the thermionic beam brightness in terms of the measurable parameters is given by

$$B_\perp = 4\pi J \frac{mc^2}{k_BT}.$$ \hspace{1cm} (2.17)

The maximum achievable transverse brightness from a thermionic gun is typically about $10^{13}$ amps/m², which is two orders of magnitude lower than the required LCLS beam brightness.

Several schemes to overcome or replace the low-brightness thermionic gun source have been proposed. There is the method of using beam bunch compression to increase the current density in order to drive up the brightness [60]. However there is an introduction of emittance growth and energy spread increase from this approach. A sounder solution is proposed where the rf photoinjector is used as a replacement for the high-brightness beam source. The basic working
principle of an rf photoinjector evades the use of thermal process. In the next section, the rf photoinjector will be discussed.

### 2.1.3.2 The Photo-injector Gun

The photoinjector gun [61, 62] was developed to replace the DC thermionic gun as the next high-brightness electron beam source for the advanced accelerator applications. The photoinjector combines both the photoemission process and the high-peak rf electric field gradient (~100 MV/m) to produce a high-current density (a couple thousand Amp/cm$^2$), low-emittance, relativistic beam.

The central components of an rf photoinjector gun are a metal cathode serving as an electron source by photoemission and rf cavities serving as a high electric gradient standing wave accelerating structure (The schematic is shown in Figure 2.5). An electron gas is generated when a laser beam energy exceeding the work-function of a metal impacts on the photocathode surface. The UV laser works well with copper metal used as a photocathode material. We note that a driving laser duration is proportional to the electron beam length so that tuning a laser system enables an initial control over the electron beam length. In this way, picosecond electron beams can be generated. Additional benefits of using a short driving laser are found in the production of a small energy spread and a low emittance beam. An energy spread relation to the driving laser duration is shown by $\sigma_\gamma \propto (k\sigma_z)^2$. Similarly an emittance has a relation of $\varepsilon \propto k^3\sigma_z^2\sigma_x^2$. These relations show that since $\sigma_z$ is proportional to the laser duration, a high-quality electron beam can be produced.

Electrons obtained from a photocathode are accelerated immediately to relativistic velocities in order to minimize the growth in the transverse space-charge force. This is evident from the transverse space-charge force varying as $1/\gamma^2,$
where $\gamma$ is a beam’s relativistic energy. A huge gain in beam acceleration is provided by the rf cavity structure, which consists of electromagnetically coupled cells with axial openings between them to allow a beam propagation from cavity to cavity. The acceleration is accomplished either in “$\pi$-mode”, where the phase of the rf fields in cells are exactly 180 degrees out of phase or in the “0-mode” where the fields in cells are in phase. The “$\pi$-mode” on axis has the time-dependent sinusoidally varying field

$$E_{z,rf}(z,t) = E_0 \sum_{n=1}^{\infty} a_n \cos(nk_{rf}z) \sin(\omega_{rf}t + \phi_0),$$

(2.18)

where $\phi_0$ is the rf phase at the cathode $z=0$ at time $t=0$, $k_{rf}$ is the rf wave number and $\omega_{rf}$ is the rf frequency. The acceleration of a beam in rf cavities starts from the photocathode where free electrons are produced. Immediately these free
electrons see longitudinal gradient-field accelerating them towards the aperture leading into the second cell. By the time a beam arrives at the aperture, the rf field in cells drops to zero in order to minimize the fringe field destabilizing this beam. As this beam emerges into the second cell, the rf field reverses its gradient direction forward to further accelerate the beam down the rf cavity. The process of this beam acceleration mode in rf cavities is illustrated in Figure 2.6.

The photoinjector system has been proven to be the best high-brightness electron beam source. But the copper metal used as a photocathode can only take a certain amount of UV laser energy before irreversible surface damage may occur. A driving laser dedicated to maximize the quantum efficiency of the cathode has been built at PLEIADES. The details of the design of this laser system will be discussed later.

2.1.3.3 PLEIADES RF Photo-injector

A high-brightness beam is central to the high production of Thomson x-rays. This feat is accomplished with the employment of the BNL/SLAC/UCLA rf photoinjector [63, 55] at the LLNL PLEIADES facility. The rf photoinjector gun installed at the facility is a 1.6 cell-type, based on a high isostatic pressure annealed oxygen free high conductive copper metal that was built at UCLA in
The photocathode was carefully polished with a diamond-turn polishing process to a perfect flat surface plane. The photocathode surface has a measured flatness of 79 nm peak-to-valley over a 2 inch diameter. The BNL/SLAC/UCLA 1.6-cell photoinjector gun was installed and produced the first electron beam in 2000.

The rf photo-injector was designed to produce a 100 MV/m standing wave electric field operating in an S-band frequency (2.8545 GHz) with 7 MW peak power and 3 µs pulse length. The S-band frequency is originally generated and multiplied to 2.8545 GHz from a 81.557 MHz laser master oscillator source. The standing wave field created in the rf cavities induces a rapid relativistic energy gain in electrons to 5 MeV in less than a 10 cm distance. A focusing solenoid integrated with the photoinjector provides the preservation of the transverse emittance and the easy transportation of an electron beam into the SLAC-type linac section.

An average peak photoelectron charge measures 0.4 nC per pulse with a pulse length of about 10 ps. Theoretically, a full nC charge extraction is possible by focusing a UV laser harder, but serious cathode damage can result if the cathode is persistently put under extreme conditions. The quantum efficiency of a photo-
injector measures $2.2 \times 10^{-5}$ with an integrating current transformer positioned directly at the photo-injector output and a calibrated UV energy meter monitoring a split laser beam. The measured transverse emittance is found anywhere in the range 1-10 mm mrad, depending on the condition of the driving laser system. A beam energy spread is measured consistently below 1%. From these measured electron beam parameters, we can deduce that the normalized beam brightness is well over $>10^{30}$.

### 2.1.4 100 MeV LINAC

The high-density current beam at the output of the rf photoinjector is injected into the LLNL rf linac [54] in order to gain more beam energy. The rf linac consists of five SLAC-type traveling waveguides used as a beam energy booster to produce an output beam energy anywhere between 20-140 MeV. Each section is an iris-loaded waveguide of 2.6 m in length. The rf linac can operate in either a high-repetition rate, which is the long macropulse thermionic electron source
used for the high-average positron current experiment or a low emittance beam at a slower repetition rate of 10 Hz, using the photoinjector for the PLEIADES experiment. A magnetic switchyard was built around the rf photoinjector, so that the thermionic gun is placed upstream of the photoinjector. When the positron-production experiment is in operation by other experimentalists, the magnetic switchyard is turned on to steer the thermionic beam around the photoinjector.

### 2.2 The PLEIADES Laser System

The state-of-the-art laser system dedicated for the Thomson x-ray scattering experiment was constructed at LLNL PLEIADES. A single laser source located in the laser-room above ground provides the photo-injector UV laser and the high-intensity scattering IR laser source. The PLEIADES laser system also serves as a master oscillator maintaining a sub-ps time-synchronization of the Thomson scattering operation. The PLEIADES laser system is divided into two subsystems: Chirped-Pulse Amplification (CPA) FALCON laser and Photoinjector Laser System (PLS) (see Figure 2.9) [1]. In the remaining half of this chapter, the PLEIADES laser facility will be described.

#### 2.2.1 The Master Oscillator Pulse System

The head of the laser system is a mirror-dispersion-controlled Kerr-lens mode-locked, Ti:Sapphire oscillator from Femtosource, which generates a 818 nm and 37 nm bandwidth IR laser with a pulse duration of 30 fs (oscillator output measurement shown in Figure 2.10). The laser power output generated is 400-450 mW at the repetition frequency of 81.557 MHz (the 35th subharmonic of the S-Band). The oscillator outputs a laser energy of 5.5 nJ per pulse. Overall, the
Figure 2.9: The overall schematic of PLEIADES. [Ref. [1]]
oscillator is a critical component for providing a seed laser for both FALCON CPA laser and photoinjector UV laser. The main oscillator also acts as a master triggering signal for all of the timing-sensitive equipment.

The latter feat begins by monitoring the oscillator output signal with two photodiode detectors. The first photodiode signal is split into two halves, where the first half is sent over to an amplifier where frequency-filters produce a sinusoidal waveform that is frequency multiplied in a phase-locked dielectric resonant oscillator (PDRO) to 2.8545 GHz. The output PDRO signal is then used to drive the rf linac amplifier. The second-half is sent to a Time-Bandwidth CLX-100 timing stabilizer where an 81.557 MHz crystal oscillator signal and the measured oscillator pulse are phase locked. The phase lock is achieved with a feedback-control system composed of a piezo-actuator. The piezo-actuator controls the high reflector at the end of the oscillator cavity in order to keep high-speed variations.
under control. The longer term drift is corrected with the adjustment of the cavity length, which is controlled by a motorized translation stage. Any phase lock stabilization not resolved by the CLX-1000 can still be tracked by the rf system, minimizing timing jitter problems. The second photodiode signal detected is sent to a SpectraPhysics SM-1 synchronization box to be compared with a 10 Hz pulse train. The synchronization box then generates a 10 Hz signal phase locked with the oscillator signal, which then is used as a trigger signal for the rest of the timing-sensitive systems. The relative jitter of 1 ns is measured between signals.

The phase-lock oscillator pulse is stretched to 680 ps by an aberration free all-reflective parabolic mirror based expander as shown in Figure 2.11 [64]. The 1480 line pair/mm grating stretcher is used as an expander to diffract pulses towards the parabolic mirror, which is located at an angle of 55°. The focused pulse reflecting off the flat mirror is recollimated at a roof mirror. Finally, a pulse reflected off the roof mirror is propagated back along the same path just above
the input pulse. The emerging beam then splits into two beams with a dielectric beam splitter where the 30% of the beam is sent to drive the photoinjector and the remaining 70% of the beam is used for the FALCON laser.

2.2.2 FALCON Laser: Chirped Pulse Amplification (CPA) System

The FALCON laser is a titanium-doped sapphire (Ti:Sapphire) based chirped pulse amplification (CPA) [44, 45] laser system, which produces a compressed 540 mJ, infrared (820 nm) laser beam suitable for the Thomson scattering experiment. At the final stage, a parabolic final focusing mirror is used to produce a $1/e^2$ laser beam-waist of 36 $\mu$m. The parabolic focusing lens currently installed produces a laser vector potential of $0.2 < a_0 < 0.5$, evading the non-linear Thomson effect for $a_0 > 1$.

The 70% of the oscillator pulse from a dielectric beam splitter is amplified to an energy of 7.3 mJ in a standard linear regenerative amplifier cavity. This regenerative amplifier consists of a 2.4 m long linear cavity with two 1.5 m focal...
Figure 2.13: Schematic of FALCON laser. [Courtesy of D. J. Gibson]
length concave end mirrors and a 10 mm (diameter) × 10 mm (long) Ti:Sapphire rod, which is pumped by 45 mJ of 532 nm light from a frequency doubled SpectraPhysics GCR-190 Nd:YAG laser. The injection of the input pulse into the amplifier-cavity is controlled by the polarization switching pockels cell located between two polarizers. The s-polarized pulse is changed to p-polarization state when the pockels cell is at the half-wave voltage. Then the p-polarized pulse filters through the second polarizer. Immediately after the passage of the p-polarized pulse, the pockels cell voltage is removed before the pulse returns to it, effectively trapping the pulse in the cavity. After the pulse energy is amplified by the pump laser, it is ejected from the cavity by switching back to the initial s-polarization in the same pockels cell. A second pockels cell, a waveplate and a Glan-Taylor polarizing cube located immediately after the regenerative amplifier are employed to filter out all other noise generated in the amplifier cavity and the oscillator.

The amplified 7.3 mJ pulse emerging from the regenerative amplifier is injected into two power amplifiers in sequence. The first power amplifier is an $\alpha$-4-pass consists of a four-pass bow tie geometry with a 25 mm (diameter) × 16 mm (long) Brewster cut Ti:Sapphire crystal used as a energy gain medium. The measured spectrum from the $\alpha$-4-pass is shown in Figure 2.14. This Ti:Sapphire is pumped with the 200 mJ of 532 nm laser from the GCR-190. The $\alpha$-4-pass amplifier is also equipped with an automated pointing centering system, which helps to minimize long-term drift in the laser transportation to the experimental chamber. This system monitors the laser leakage through the two mirrors with two cameras, and if a leakage is detected the beam position is adjusted correctly to the center. The final output beam energy produced is 63 mJ (the total gain of 9), and $1/e^2$-beam dimension of $1.48 \times 1.28$ mm.
Figure 2.14: a) Regenerative amplifier output spectrum and beam profile. b) α-4-pass output spectrum and beam profile.

The laser pulse goes through a final energy amplification process in the β-4-pass amplifier with a four pass bow tie configuration similar to the α-4-pass. A 25 mm (diameter) × 25 mm (long) Ti:Sapphire crystal is pumped simultaneously by a SpectraPhysics PRO-350 frequency-doubled Nd:YAG laser with an output energy up to 1.5 J of 532 nm and a Continuum frequency-doubled Nd:YAG laser with 1.2 J of 532 nm. The β-4-pass is also equipped with an automated pointing and centering system. The total gain from the β-4-pass is 20, producing an uncompressed laser energy of 1.2 J. From the measured beam profile, the upcollimated $1/e^2$-beam radius is 42 mm.

The amplified beam is transported 46.5 m to the pulse compressor chamber.
where it is compressed in a double-pass grating compressor. The final compressed pulse length measured with the GRENOUILLE [65] system at a low power is 54 fs long with a phase retrieval error of 0.6%. The compressed beam then propagates 20 m to the final focusing optics, which focuses a laser down to a spot-size of 42 \( \mu m \) radius FWHM and an \( M^2 \) value of 1.64. 45% of the beam is lost during the compression and the transportation stages, making the output an effective beam energy of 540 mJ that is available for the Thomson scattering.

### 2.2.3 Photo-injector Laser System (PLS)

The photoinjector for the production of a high-brightness electron beam is driven by a dedicated high quality UV laser system. The second 30% oscillator pulse of a 680 ps duration from the dielectric beam splitter is transported 50 meters through a single-mode fiber to a PLS hutch located in an underground accelerator area. The PLS hutch is designed to maintain a stable laser operation temperature, and minimizes dust with the high efficiency particulate air (HEPA) filters. The oscillator pulse, which has an average energy of 90 pJ, is then coupled to a linear regenerative amplifier to produce 5.9 mJ pulses. At the heart of the amplifier system is a Ti:Sapphire, which is pumped by a 50 mJ of 532 nm light from a SpectraPhysics DCR-2 laser. The regenerative amplifier is equipped with a fast photodiode light-leak detector for adjustments. The IR laser pulse from this regenerative cavity then enters a bow-tie configuration 4-pass power amplifier, where the energy per pulse is upscaled to 90 mJ. Again, the system employs a Ti:Sapphire crystal and a 280 mJ laser from the SpectraPhysics DCR-2.

The amplified laser is then transported to the beam compressor where the laser is compressed to 10 ps long. The laser is intentionally not fully compressed to its Fourier transform limit in order to produce the best quality electron beam.
Otherwise, a fully compressed laser produces the Coulomb repulsion of the electrons. An additional advantage of having a slightly longer pulse is a reduction of the residual cubic phase distortions introduced by the 50-meter fiber being used for transporting the laser to the PLS.

After the pulse compression stage, the pulse is transported to the next system where the frequency of the light is doubled in a Type I BBO crystal and immediately tripled in a second Type I BBO crystal to produce a 269 nm in wavelength. After the compression the UV laser is left with 1.2 mJ in energy, but the system is turned down to 500 µJ to avoid the damage to the photo-cathode.
material. An aperture is employed to provide a hard-edged round beam of 2 mm in diameter, which improves the quality of electron beam emittance. Finally, a 7 ps long UV laser is transported a distance of 50 meters to the photoinjector Cu cathode where high-quality electron beams are generated via photo-electric effect.
CHAPTER 3

The Ultra-strong Halbach Permanent Magnet Quadrupoles

3.1 Introduction

In the past decade, the introduction of the rf photoinjector has revolutionized the field of accelerator physics. The technology allowed the production of high brightness electron, defined as simultaneous high peak current and low transverse emittance, such that making possible to explore the innovative accelerator concepts like a plasma wakefield accelerator. But, a little effort has been paid attention to the quadrupole magnet, a component of focusing optics necessary for the production of ultra-small electron beam immensely important for the production of high Thomson x-ray flux. To produce the small transverse focal spot, the length scale of the focal system has to have the minimum “β-function,” which is related to the minimum transverse beam size by $\sigma_x^* = \sqrt{\beta_x^* \varepsilon_x}$.

On the other hand, when focusing the beam, one must simultaneously avoid chromatic aberration due to energy spread, and emittance growth due to residual beam space-charge effects. These effects impose serious limitation in achieving 10-20 μm beam with conventional electromagnet quads whose focusing strength is weak and has a long focal length. The reason is that in order to compensate for the weak focusing strength of the EM quad, a large input beam into the
focusing optics is used. A larger input beam would subject to a greater bending angle needed to produce a desire small beam, but at the cost of introducing the chromatic aberration effect. This issue is simply illustrated by the following relation governing the ratio of final spot size $\sigma^*$ to initial $\sigma_0$, [66],

$$\frac{\sigma^*}{\sigma_0} = \sqrt{\frac{1 + \left(\frac{\beta_0}{f}\right)^2 \left(\frac{\sigma_{\delta p}}{p}\right)^2}{1 + \left(\frac{\beta_0}{f}\right)^2 \left[1 + \left(\frac{\sigma_{\delta p}}{p}\right)^2\right]} \approx \frac{\sigma_{\delta p}}{p},}$$

(3.1)

where the final term obtained in the limit $\beta f \gg \frac{p}{\sigma_{\delta p}}$. When the initial $\beta$-function $\beta_0 = \sigma_0^2/\varepsilon$ is larger than the effective focal length $f$ of the final focus system, strong compression of the beam size is possible, up until the chromatic aberrations begin to dominate. There is a minimum in the beam size, for a fixed focal length system, when $\beta_0/f \cong p/\sigma_{\delta p}$. In this condition, the demagnification is $\sqrt{2}\sigma_{\delta p}/p$. With larger initial beam size, the chromatic contribution to the emittance grows rapidly and because of the demagnification limit, the final beam size actually grows with larger initial beam size.

We illustrate the argument with an example of EM quad triplet system at PLEIADES. With the present generation of large aperture quads that UCLA employs, which have a maximum gradient of 15 T/m over $\sim$10 cm length, the minimum value of the effective focal length of the final focusing triplet is $f \approx 50$ cm, and thus optimum demagnification for $\sigma_{\delta p}/p=0.5\%$ is only possible for $\beta_0=100$ m, or (with $\varepsilon_x = \varepsilon_n/\gamma=0.14$ mm mrad) $\sigma_0=3.2$ mm. We thus anticipate a minimum achievable spot size (Eq. (3.1)) of $\sigma^* \approx 32 \mu$m. This spot size estimate, which has been verified in simulation, is unacceptably large for the Thomson source, which ultimately demands a factor of 4 smaller spot size for nonlinear Thomson interaction studies, as well as for increased x-ray photon production.

In addition, by allowing the beam to expand before final focus, space-charge forces (for beams of energy $< 100$ MeV) will increase the emittance [67], giving
an even larger final spot size. This limitation also precludes more elaborate approaches, based on sextupole correction, to eliminate chromatic aberrations. At low energy therefore, transport and final focusing systems must be compact; short focal lengths are needed.

From this introduction, it is clear that one must look into the use of much shorter focal length lenses in order to create minimum $\beta$-functions small enough for many modern electron beam applications. One might consider superconducting quadrupoles or solenoids, but these devices are neither inexpensive, nor easy to build with the necessary small (cm-scale) dimensions. On the other hand, permanent magnet-based quadrupoles (PMQs) can access high-field gradients, and have been under study recently in two arenas: linear collider (LC) main accelerator focusing – to mitigate the cost of the needed power supplies – and final focus magnets in both circular accelerators [68] and LCs [69, 70, 71, 72]. Also, an optimum, static intra-magnet configuration of magnet has been studied in the context of the Compact Linear Collider (CLIC) program [70, 71], and has been implemented at the Cornell Electron Storage Ring (CESR) in the CLEO detector mini-$\beta$ system [73].

In the present chapter, we describe the design parameter studies and error-simulation studies of PMQ. The discussion then will continue with designing, construction and measurements of the PMQ triplet system. In the next chapter, we will continue our discussion on the beam dynamic studies of the optimum quad triplet system, magnet error effects on the beam dynamics, and the system implementation and the beam measurements.
3.2 Magnet Theory and Simulations

The starting point of the design study of PMQ is determining the right physical size such that it is compact, but also be able to produce a high magnetic field gradient. In fact, the focal length of the triplet system as a whole should be less than 10 cm, the individual PMQ length is chosen to be 1 cm. If one restricts, as is appropriate for a triplet configuration, the focusing phase advance (for the initial operating energy of 70 MeV) to at most one-half ($k_q l_q = (B'/BR)l_q \leq 0.5$) in a single 1 cm lens, then over a 600 T/m field gradient is deduced. We shall see in the next chapter that this estimate on appropriate field gradient indeed yields effective guidelines for optimizing the beam optics in the PLEIADES final focus system.

In this section, a 2-D analytical model originally approached by Halbach and a 3-D simulation model are presented to illustrate the dependence of the PMQ performance on the geometric and magnetic design parameters in the Halbach pure PM design.

3.2.1 2-D Halbach PMQ Model

We present first the two dimensional Halbach-type PMQ model, reviewing results that were originally established by Halbach [74]. These results are then used to guide initial magnetic and beam optics design decisions.

A PMQ structure can be composed of multiple segments – in the design considered here, 16 wedges – where each wedge is a trapezoid. A complete assembly of all 16 segments gives a cylindrical piece with a small bore about the center of its axis. The magnetization easy-axis in each trapezoidal segment rotates 45° from one segment to the next. A schematic cross section of a 16-piece quadrupole
Figure 3.1: A cross section of the 16 segment PMQ and the magnetization orientations shown as arrows.

is illustrated in Figure 3.1.

To estimate the achievable gradients one obtains with a PMQ, we first assume a permanent magnet geometry in which the magnetization varies continuously with an azimuthal angle $\phi$. The total magnetic field gradient for this pure quadrupole consists of Amperian current contributions from the boundaries at inner and outer surfaces, as well as from the bulk of the magnetic material, and is given by

$$B' = 2B_r \left( \frac{1}{r_i} - \frac{1}{r_o} \right).$$

(3.2)

With $B_r=1.22$ T, $r_i=2.5$ mm, and $r_o=7.5$ mm, we have $B'=650$ T/m. In practice, the field gradient is most strongly controlled by the inner radius $r_i$.

It is not possible to make a PMQ with a continuously varying remanent field, but the use of uniformly magnetized segments can well approximate the pure quadrupole case. This was recognized by Halbach, who studied the approximation to the ideal geometry that can be achieved by segmenting the magnet into $M$
geometrically identical pieces, which have magnetization rotated appropriately in each segment. In this case, the multipole expansion of the \( M \)-segmented magnet can be expressed by \([74]\)

\[
B = B_r \sum_{j=0}^{\infty} a_n \left( \frac{r}{r_i} \right)^n \frac{n}{n-1} \left[ 1 - \left( \frac{r_i}{r_o} \right)^{n-1} \right],
\]

\[
a_n = \cos \left( \frac{n\pi}{M} \right) \sin \left( \frac{n\pi}{M} \right) \frac{M}{n\pi}.
\] (3.3)

The segmentation of the magnet into \( M \) uniformly magnetized sections is accounted for by the coefficient \( a_n \). This term is slightly less than unity, so the field predicted in Eq. (3.3) is slightly diminished compared to the field obtained from the ideal case. For the Halbach geometry considered, where the magnet has \( M=16 \) segments, the coefficient of quadrupole term, \( a_2 \), is found to be 0.94 – the field gradient is diminished by only 6% from the ideal geometry.

The model analysis above omits the effect of finite length in the longitudinal dimension of the PMQ. This analysis requires a 3-D model, which is presented in the next section.

### 3.2.2 3-D Radia Simulation Model

The magnetostatic simulation code Radia [75] was used in 3-D modeling and numerical analysis of a 16-piece Halbach PMQ. Radia, in contrast to most magnetostatic field solvers based on Finite Element Method (FEM), is a calculation tool utilizing the Boundary Integral Method (BIM). The code is written in Object Oriented C++ and possesses interface capability with Mathematica for pre- and post- simulation processing. The magnetostatic object of interest can be formulated as a composition of a set of smaller geometric components bounded by planar polygons which then are linked to make up the desired magnet model. Furthermore, each component can be segmented into smaller subdivisions for
In this section, we present the simulation results of the ideal PMQ model and analyze the various causes contributing to demagnetizations in the ideal PMQ. The section finishes with the discussion of the field property errors arising in the imperfect geometrical PMQ model.

The desired goal from 3-D simulation analysis is to obtain a set of optimized object dimensions which, consistent with a given remanent field, yield an appropriately high field gradient. Our application demands that the overall dimension of the PMQ be consistent with the beam optics dimensions (few cm focal lengths for each lens), and that the beam be easily passed through the small bore where the strong field gradient is produced. The design analysis also must factor in technical difficulties in the manufacturing process. NdFeB, a rare-earth PM material, was chosen for our simulations over other rare-earth types, e.g., SmCo, because of our demand for a higher remanent field range.
Figure 3.3: 3D code Radia PMQ field: the magnetic field (left), the field gradient (middle), and the longitudinal field taken at \( r=2.0 \) mm (right).
Each component wedge block is bounded by two parallel trapezoidal planes separated by a length of $l_z$ and a given remanent field of $B_r$. By implementing a simple loop algorithm routine in the Radia, 16 wedge blocks are sequentially stacked side-by-side where the final model design is created. The final optimum PMQ dimensions determined from the simulations in considerations of a design compactness and a peak field gradient are: $r_i=2.5$ mm, $r_o=7.5$ mm and $l_z=10.0$ mm. The field gradient that can be produced with this set of parameters is $B'=573$ T/m. Note that the gradient predicted in Radia simulation is lower than that deduced from Eq. (3.2) mainly due three-dimensional (edge) effects. The field gradient and the magnetic field profiles plotted in the bore region are shown in Figure 3.3. The effective magnetic length determined from the simulation is $l_{eff}=10.4$ mm (see Figure 3.3).

3.2.2.1 Demagnetization

In this section, two of the most relevant factors contributing to changes in the performance of the NdFeB based PMQ are discussed: relaxation in PMQ due to intra-magnetic interactions, and a partial demagnetization due to temperature variations. The demagnetization studies in Radia are thus used for determining whether the coercivity and remanence in the NdFeB material are adequate.

The close proximity and high remanent fields of the segments can influence the overall magnetic characteristics of the assembled PMQs through mutual magnetic interactions. Relaxation modeling was performed by iterative calculations of the field strength and the magnetization of each wedge block under the influence of the other 15 wedge blocks, until the values reach the specified precision in the stabilized magnetization.

With the NdFeB magnet parameters from the material datasheet and the
Figure 3.4: The magnetic field lines plotted showing decrease in slope as temperature value increases (top) and the percentage loss of field gradient as a function of temperature (bottom).
calculation precision set in Radia, the simulation showed a negligible relaxation in the chosen geometry of the PMQ. However, a permanent magnet material with a high coercivity is still desirable for its high resistance against opposing external fields and its resistance to temperature variation induced demagnetization.

Understanding the effect of temperature variation on the permanent magnet is important in the situation where an accidental exposure of a NdFeB PMQ to a temperature elevation can cause demagnetization as, for instance, during a vacuum bake-out. Demagnetization effects in the simulation studies are introduced through a second order temperature dependence in the magnetization and the coercivity equations as [76]:

\[
M(T) = M(T_0) \left[ 1 + a_1(T - T_0) + a_2(T - T_0)^2 \right],
\]

\[
H_{cJ}(T) = H_{cJ}(T_0) \left[ 1 + b_1(T - T_0) + b_2(T - T_0)^2 \right],
\]

where \(T_0\) is a reference temperature, and coefficients \(a_i\) and \(b_i\) are obtained from the material data sheet.

The relaxation of a PMQ magnetization is performed for 5 different chosen temperatures: 0, 20, 60, 90 and 120 °C. The temperature increase has a negligible effect in the magnetization loss in the PMQ at low temperature elevations, but the effect becomes noticeable with higher temperature. In Figure 3.4, the irreversible loss of a PMQ magnetization is shown as a function of a temperature change. This loss is small under 100 °C, but it increases dramatically above 120 °C. This behavior is similar to that obtained from simulations reported in Ref. [76].

### 3.2.2.2 Geometric and Magnetic Errors

Geometric imperfections in PMQ due to fabrication and assembly errors are investigated in simulations and the evaluation of the magnet performance are discussed
Figure 3.5: The field gradient plots with magnetization orientation errors. As the angular orientation error increases, a less symmetrical magnetic profile results with increasing center offset.

here. We concentrated our simulation efforts in three main cases of PMQ imperfections: angular orientation errors in the magnetic easy-axis; sector-to-sector geometric errors; and PMQ bore aperture radius errors. From these simulations, good approximations in allowable tolerances in mechanical errors for the individual modules are obtained. A complete error budget, including alignment errors, is discussed in later sections.

In the first case, the following set of errors are introduced into the magnetization easy-axes: 0.225°, 0.450°, 0.675° and 1.125°; corresponding to relative errors of 1%, 2%, 3%, and 5%, respectively. The magnetic field gradient is solved and presented for each angular error cases as shown in the Figure 3.5; as the magnitude of the angular error increases, the field profile suffers from a shift in the magnetic center and introduction of variation in the field gradient (higher
Figure 3.6: The field gradient plots with magnet wedge shape errors. Each curve represents wedge shape (mechanical) deviation from the ideal. Again the plots show increasing uneven characteristics with increasing percent error.

The wedge geometry is different from that of an ideal trapezoid sector in two different ways: an angular deviation from 22.5° and a wedge length deviation from \(|r_o - r_i|\). The following errors in wedge angles (in this particular exercise, the side length error is excluded, as we shall discuss in the magnet fabrication, that final boring out of inner pole surface yields a highly uniform surface) are chosen as the simulation input parameters: 0.225°, 0.450°, 0.675° and 1.125°; again corresponding to relative errors of 1%, 2%, 3%, and 5%, respectively, as shown in Figure 3.6. The error parameters are assigned to an arbitrary number of wedges in different positions which consequently create varying gaps. The field gradient profile shows a significant sloping effect and a reduction in amplitude when the errors exceed 2%. 

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As we shall discuss below, one of the final steps in the manufacturing process is boring out the PMQ hole to a specified value, in our case 2.5 mm. In doing so, it is prudent to understand the effect of an undercut or overcut. As shown in the theoretical analysis, a variation in the PMQ bore radius is inversely proportional to the field. In the final case of our imperfect PMQ model, the following set of PMQ bore radii are considered: 2.45 mm, 2.50 mm and 2.55 mm. As expected, the field gradient predicted in different radius parameters gives 591 T/m and 556 T/m, corresponding to inner radius of 2.45 mm and 2.55 mm, respectively. This variation is $\Delta B' \approx \pm 18$ T/m. In comparing with the 2-D Halbach PMQ model, the variation is deduced from Eq. (3.2) to be $|2B_r \frac{\Delta r}{r}| \approx 19.5$ T/m.

Three notable error studies are done providing the following results that have been implemented in our magnet design: (1) a magnetization easy-axis angular orientation error is allowable up to 2% before a notable change develops between
the magnetic center and the mechanical center; (2) the wedge shape error (the wedge angle) is limited and allowed up to again 2%; and finally (3) bore radius variations of ±0.05 mm from the nominal 2.5 mm radius produced the field gradient variation of ±3% from the ideal \( B' \). In Chapter 4, magnet errors that are incorporated into the final focus system have un-desired effects on the beam dynamic, therefore the beam simulation results will be discussed.

### 3.3 System Design and Properties

#### 3.3.1 Magnet Design and Fabrication

Fabrication of an ultra-high field gradient Halbach-type PMQ with good field quality has proven to be very difficult from an engineering perspective; assembling highly magnetized, small trapezoidal segments into a precisely aligned unit with a small bore poses mechanical challenges. In addition, fabrication of multiple PMQs with both identical magnetic and geometric characteristics increases the constraints.

Nevertheless, a simplified fabrication procedure for producing a high quality, and ultra-high gradient PMQ may be given as follows: a long trapezoidal wedge which is pre-magnetized to a \( B_r = 1.22 \) T is machined to a specified design shape from a larger magnet block through an high precision cutting process, wire electrical discharge machining (EDM). The 16 magnetized wedges are then coupled together carefully into an equally long aluminum keeper tube. In addition to magnet wedges being held inside the keeper, a thin layer of cohesion material is applied, strongly joining neighboring magnets. The tube keeper and a bonding glue together ensure structural stability of the PMQ under extreme magnetic repulsions. The long assembly is then cut into six identical PMQs, with
Figure 3.8: An illustration of the PMQ manufacturing process showing an individual wedge, the assembled 16 wedge rod, and a single PMQ cut from the rod.

mechanical length of 1 cm, to ensure consistent performance between all PMQs. In particular, construction of a single, long magnet assembly assured that the relative azimuthal (roll) errors of individual units may be essentially eliminated. Finally, the inner diameter of bore is slowly ground to the specified value of 5 mm with a diamond cutter. An illustration which schematically summarizes the PMQ fabrication procedure is shown in Figure 3.8.

3.3.2 Magnet Material Properties

High performance permanent magnet materials are available in large quantities and a variety of formulations such as neodymium iron boride (NdFeB), samarium cobalt (SmCo) and ferrite (Fe$_2$O$_3$). NdFeB is particularly suitable for our application due to its low-cost, high remanent field, high intrinsic coercivity ($H_{ci}=21$
kOe), and good mechanical characteristics. In contrast, SmCo is generally known to be weaker in $B_r$ by 25%, extremely brittle and more expensive than NdFeB. On the other hand, the major disadvantage of NdFeB is that its Curie temperature is much lower than that of SmCo. Consequently, a quadrupole fabricated from NdFeB is more sensitive to temperature fluctuations.

For the application at hand, the benefits of NdFeB far outweighed the drawbacks. A particular grade of material was selected based on various parameters, availability and homogeneity. Key parameters of the NdFeB (grade 35SH) used in the PMQs are listed in Table 3.1. It is worth noting that higher remanent fields are available in different grades; however, it was felt that grade 35SH provided sufficient field without compromising coercivity and homogeneity. In applications requiring larger bores and higher field gradients, one might consider other grades.

The manufacturing process required uniform magnet material properties because no sorting process was relied upon when selecting for the individual magnet wedges. Instead, the long magnet wedges were all machined from a single pre-magnetized block, with the field orientation assured by the cut angle rather than

Figure 3.9: A PMQ side-by-side size comparison with a US quarter.
Table 3.1: NdFeB 35SH properties

<table>
<thead>
<tr>
<th>Property</th>
<th>Value</th>
</tr>
</thead>
<tbody>
<tr>
<td>Remanence $B_r$ [T]</td>
<td>$1.19 \sim 1.25$</td>
</tr>
<tr>
<td>Coercivity $H_{cB}, H_{cJ}$ [Oe]</td>
<td>$1.13 \sim 1.21 \times 10^4, 2.10 \times 10^4$</td>
</tr>
<tr>
<td>Maximum Energy Product $(BH)_{\max}$ [MGOe]</td>
<td>$34 \sim 39$</td>
</tr>
<tr>
<td>Recoil Permeability $\mu_r$ [-]</td>
<td>$1.05$</td>
</tr>
<tr>
<td>Thermal Coefficient of $B_r, H_{cJ}$ [%/K]</td>
<td>$-0.09, -0.53$</td>
</tr>
<tr>
<td>Coefficient of Thermal Expansion $C_\parallel, C_\perp$ [1/K]</td>
<td>$7.2 \times 10^{-6}, -1.1 \times 10^{-6}$</td>
</tr>
<tr>
<td>Curie Point $[^\circ\text{C}]$</td>
<td>$360$</td>
</tr>
<tr>
<td>Density $[\text{kg/m}^3]$</td>
<td>$7.55 \times 10^3$</td>
</tr>
<tr>
<td>Thermal Conductivity $C_\parallel, C_\perp$ [W/m $\cdot$ K]</td>
<td>$6.2, 6.7$</td>
</tr>
<tr>
<td>Specific Heat $[\text{J/kg} \cdot \text{K}]$</td>
<td>$4.6 \times 10^2$</td>
</tr>
</tbody>
</table>

a post-machining magnetization. This recipe transferred the possible sources of error and precision to the machining process, and spared us from having to measure very small magnets and sort hundreds of wedges.

### 3.3.3 Mover System

Before implementing the PMQ system on PLEIADES, general problems associated with the application of this focusing system in the experimental environment had to be solved. Addressing the mechanical tolerances and space constraints of the PMQ system, while maintaining an environment consistent with ultra-high vacuum (UHV) and magnetic fields, required careful integration.

Precision machining of all parts eliminated the need for adjustable alignment in the movers that the PMQs were mounted upon. Computer controlled custom-fabricated vacuum linear motion feedthroughs were utilized to control
each quadrupole’s longitudinal position, which was precisely constrained by a rail system. Motion of the PMQs was set by use of stepper motors and LabVIEW control software. The entire magnet axis of the system was aligned by use of a pulsed wire system and optical techniques. Other features of the assembly include a ceramic ablation disc at the front of the assembly to protect the magnets from misaligned high-intensity FALCON laser damage (during the Thomson experiment), and a scintillating disc at the (beam entrance) side of the assembly, for locating mis-steered electron beams.

From a mechanical design approach, the system had to meet the following requirements: independent (bench top, pre-installation) correction of magnetic center to mechanical center errors for each PMQ assembly, independent position control of PMQs to within 100 $\mu$m in the beamline direction, magnetic center-to-center transverse placement with 10 $\mu$m precision for each PMQ to a center line along the range of motion, insertion of assembly into a standard 4″ inner diameter (ID) vacuum chamber with 6″ outer diameter (OD) flanges, and compatibility with vacuum at a nominal pressure of $<10^{-6}$ Torr.

A mechanically simple design was chosen to comply with the above stringent requirements. A precision computer numerical control (CNC) machined stainless steel conflat flange was fabricated to serve as a reference monument, mechanical support and vacuum seal. Three precision ground rods were fitted into the flange by means of expansion and shrinkage fits.

Pairs of magnets were supported by CNC machined aluminum disks, each with three precision linear bearings press fit into the outer diameter. The disks and bearings move along the rods, and are controlled by custom fabricated vacuum linear motion feedthroughs. The entire mechanical assembly is supported by the main flange using a massive precision ground stand. The stand affords five axis
alignment using the three shimmed or machined legs, Figure 3.10.

An initial set of disks were machined to place each PMQ into the mechanical center of the assembly (±0.0005”). Using optical techniques (theodolite and autolevel) along with a pulsed wire system, the magnetic center of each PMQ was found relative to the assembled mechanical center – this measurement included offsets due to the sag of the rods, and machining errors. New discs, if needed, would then be machined with the offsets found, thus zeroing out any errors. Fiducials – clear cylinders with etched cross hairs – are placed inside the front and back PMQ for optical alignment.
3.4 PMQ System Measurements

Various measurements on the PMQs were performed using pulsed wire and Hall probe. In comparison with the 3-D magnetic simulations and tolerance studies, the quadrupole magnet properties indeed agree well with predicted values. More details of measurement techniques, as well as their usefulness for strength and alignment testing, are discussed in the following sections.

3.4.1 Pulsed-Wire Technique

Measurements of the PMQs were made with the pulsed wire technique, which was developed for tuning and characterizing undulator magnets [77]. The test set-up included a single meter-long, 50 \( \mu \text{m} \) diameter BeCu wire which was stretched out through the PMQ magnetic field region and anchored at end mounts. One end was fixed, while the other end had a pulley over which the wire was tensioned by a hanging weight. A pulse of 20-50 V, with a duration of a few msec to sec, and a repetition rate of 1 Hz was produced from a high voltage pulser triggered by a function generator. When the current pulse is passed along the wire through the PMQ magnetic region, it interacts with the PMQ \( B \)-field and deflects the wire. A detector consisting of a photo-diode sensor and a diode laser – with a small focusing lens to focus the beam on to the center of the wire – mounted in a Plexiglas ring, and arranged so the transmitted light is proportional to the wire displacement, was employed to quantify the deflection. Two sets of laser-sensors were used in order to measure both transverse deflections simultaneously.

The pulsed wire technique was used to detect non-linearity in the \( B \)-field dependence on transverse displacement. The BeCu wire was translated horizontally through the magnetic center from \(-2.0 \text{ mm}\) to \(+2.0 \text{ mm}\). The deflection ampli-
Figure 3.11: Field linearity scan over accessible PMQ aperture obtained with the pulsed-wire setup.

tude gives a measure of the (longitudinally) integrated vertical component of the $B$-field. According to the data reported in Figure 3.11, the transverse $B$-field shows no apparent higher multipole content; the dependence of $B_y$ on $x$ is linear to within measurement error over the scan interval.

### 3.4.2 Alignment Measurements

The final focus system in a F-DD-FF triplet configuration requires a high degree of alignment of magnetic centers between the component PMQs. Determination of the relationship between the geometric and magnetic centerlines of each individual PMQ as well as that of the complete lattice was accomplished using the pulsed wire technique.

The PMQ magnetic center to the mechanical center alignment has been stud-
ied using the pulsed wire technique. Each PMQ was mounted on a custom design V-block mount, which in turn is bolted onto two high precision micrometer stages which provide vertical and horizontal position variation. First, the end mounts were adjusted until the wire was level with respect to the vibration isolation optical table. The PMQ was then translated in both transverse directions until the BeCu wire was aligned to the geometric center of the magnet bore as the process was observed with a theodolite (see Figure 3.12). If pulsed wire deflection was observed on the digital oscilloscope connected to the detector apparatus, the BeCu wire was relocated through translations of micrometer stages to find the PMQ magnetic center (zero field). From the pulsed wire measurement, it was determined that the magnetic center was aligned to the geometric center to within the instrument observation capability which was 25 µm.

A more difficult alignment measurement of the PMQ mover system has been carried out employing the pulsed wire setup. A crucial requirement of the final
focus system is that the PMQs riding on the rails maintain transverse alignment when moved longitudinally. A main source of motion deviation can be expected due to the sagging of the rail rods due to gravity. This sagging can introduce a dynamic steering of the electron beam at the IP, as the PMQs are moved.

The alignment of the final mover system was observed with a high precision theodolite for horizontal motion and an auto-level for vertical motion. The auto-level indicated a noticeable sagging in the vertical direction. The alignment in the horizontal direction, however, was good to within 25 µm, like that of the pulsed wire system. The maximum vertical sagging was observed when magnets were moved to the downstream end of the mover system and minimized when they were placed at the fixed end of the rails. The sagging was corrected with a set-screw located on the bottom of the ring-piece which ties together the free ends of the rails.

3.4.3 Hall Probe Measurements

Although the pulsed wire technique was well suited to the tasks described above with high spatial resolution, attempts to use it as a means to measure the field gradient were limited by the apparatus conditions: wire sag, dispersion and sensitivity to vibration noise. In particular, the wavelength dependence of the wire response prevented calibration of the 1 cm long magnet measurements with a known 10 cm long electromagnetic quadrupole. As an alternative option for quantitative measurements, a miniature magnetic Hall probe was employed to determine the field gradient and the effective magnetic length of the PMQ. With the PMQ mounted on the same translation micrometer stages, the probe scans horizontally reading $B_y$. The field gradient obtained from this measurement in the center of the PMQ bore region was measured to be 560 T/m (see Figure 3.13)
using this method. We note that a Hall probe is limited in spatial resolution, which in turn limits the characterization of non-linear field components in a quadrupole magnet.

An effective magnetic length was also determined using a Hall probe by reading the value of $B_y$ at each point along a longitudinal line located 0.5 mm diagonally from the magnet axial axis. The Radia simulation predicted an effective magnetic length of 10.4 mm, where the magnetic length is defined as the integral $\int B_y dz / B_{max}$, with $B_{max}$ indicating the peak field. The measured trace, as shown in Figure 3.14, shows a good agreement with the simulation. While the entire range of the magnet is not accessible to measurement, the reproduction of the Radia results allows us to deduce the effective length, which are employed as inputs in the beam dynamics simulations, to be essentially that given by Radia.

Figure 3.13: A comparison of PMQ field profile measurements using a Hall probe to Radia simulation results.

![Graph showing comparison of PMQ field profile measurements using a Hall probe to Radia simulation results.](image-url)
Figure 3.14: The effective magnetic length is plotted with measurements (squares) and simulations (line).
CHAPTER 4

The Beam Simulations & Measurements

In the LC main accelerator application, an effort has been underway for several years to design a moderate strength, but tunable magnet. An example of such a magnet, the Halbach ring-tuned hybrid quadrupole, is shown in the 3D magneto-static simulation displayed in Figure 4.1. Much of the effort in the NLC-oriented work has been concentrated on tunability, while simultaneously holding the effective magnetic center stable, to less than 1 micron as required for utility in the beam-based alignment process growth in the LC linac.

On the other hand, the need for extremely large field gradients makes the introduction of significant gradient tuning into a permanent magnet quadrupole very challenging. In this chapter, we discuss an alternative scheme that sidesteps these difficulties, by tuning the strength of the final focus array (a class of quadrupole triplet) only through the longitudinal (z) positioning of the individual fixed-strength PMQs. In this way, the issue of tunability is separated from both that of alignment, and that of achieving high gradient.

An equally important study is understanding what effects the magnet errors have on the electron beam dynamics. Two most concerning error effects are the skew quadrupole and the magnet-center displacement errors, which could produce a substantial emittance growth in the beam if no tuning adjustment is implemented. The beam simulation studies, that will be discussed in the current chapter, performed were used in the design tolerance of the focal system and the
Figure 4.1: International Linear Collider (ILC) prototype Halbach permanent magnet-hybrid ring-tuned quadrupole RADIA simulation model, showing geometry and magnetic induction map.

tuning adjustments.

At the end of the chapter, we will describe the implementation of the final focus optics, and the design of the beamline section downstream and upstream of the Thomson scattering chamber. A methodical beam tuning procedure devised for obtaining the desired small electron beam spot at the interaction point will be described. But first, the optimum quad triplet design theory is presented in the next section.

4.1 Beam Dynamics

The beam dynamics section will begin first with the theory of optimum triplet configuration based on thin-lens analysis. Then the first-order beam simulation with Trace3D proving the focusability of the final focus system. And finally with
the second-order beam simulations with Elegant for magnet errors and beam chromatic aberration effects on the final beam properties.

4.1.1 Optimized Triplet Design

We start with the theory of optimum configuration of a simple triplet design and the consideration motivating the choices of PMQ configuration and up-stream optics and diagnostics used in an experiment.

Consider here an idealized system composed of three thin lenses separated by drift spaces and alternating between focusing and defocusing in the horizontal \((x)\) dimension. An illustration of this model problem is shown in Figure 4.2. The assumed constraints for this design are that the focus be symmetric, and that the incoming beam is symmetric and at a waist, \(\sigma_{x0} = \sigma_{y0} \) and \(\sigma'_{x0} = \sigma'_{y0} = 0\). The usual matrix formalism [78] may be employed to describe the particle trajectories in the system. Using \(M_f = \begin{pmatrix} 1 & 0 \\ -1/f & 1 \end{pmatrix}\) for a focusing thin lens (with a change in sign of \(f\) in the case of a defocusing lens) and \(M_L = \begin{pmatrix} 1 & L \\ 0 & 1 \end{pmatrix}\) for a drift, the transport matrix of the system is simply \(M = M_{L3}M_{f3}M_{L2}M_{f2}M_{L1}M_{f1}\).

For a given maximum obtainable focusing strength, or equivalently minimum focal length, \(f_{\text{min}}\), one seeks the parameters \(f_1, f_2, f_3, L_1, L_2, L_3\) that maximize the convergence angle, \(x'\) at the focal point. Ignoring the small initial angle \((x'_0 = 0)\), the final angle is given by \(x'_f = M_{21}x_0\), where

\[
M_{21} = -\frac{1}{f_3} + \frac{1 - \frac{L_2}{f_3}}{f_2} - \frac{1 - \frac{L_2}{f_3} + L_1 \left(-\frac{1}{f_3} + \frac{1 - L_2/f_3}{f_2}\right)}{f_1}. \tag{4.1}
\]

This expression can be simplified by requiring a round beam spot at the focus. This forces the matrix element \(M_{21}\) to be equal for \(x\) and \(y\). Setting the righthand
side of Eq. (4.1) equal to its counterpart in \( y \) leads to the relation:

\[
f_1 f_2 - f_1 f_3 + f_2 f_3 - L_1 L_2 = 0. \tag{4.2}
\]

Similarly, requiring \( x = y \) at the third lens by equating the \( \tilde{M}_{11} \) elements of the \( x \) and \( y \) versions of the matrix \( \tilde{M} = M_{L_2} M_{f_2} M_{L_1} M_{f_1} \) gives the equation

\[
f_1 = f_2 \left( 1 + \frac{L_1}{L_2} \right). \tag{4.3}
\]

The expected result shown by Eq. (4.3) is that the second lens must be stronger than the first, provided \( L_1 > 0 \).

Applying Eqs. (4.2) and (4.3) to Eq. (4.1) and solving for variables \( f_2 \), \( L_1 \), and \( L_2 \) leads to

\[
M_{21} = -\frac{1}{f_2^2} \frac{L_1 L_2}{L_1 + L_2}. \tag{4.4}
\]

Now, if it is assumed that the second lens provides the shortest focal length, \( f_2 = f_{\text{min}} \), the minimum spot size is achieved by maximizing the expression \( L_1 L_2 / (L_1 + L_2) \), subject to the constraints given by Eqs. (4.2) and (4.3) and the fact that \( f_3 \geq f_{\text{min}} \). These three conditions lead to the inequality, \( L_1 L_2 \leq f_{\text{min}}^2 \).
which is sufficient to find that \( \frac{L_1L_2}{(L_1+L_2)} \) is maximized when \( L_1 = L_2 = f_{\text{min}} \).

With back substitution, the optimum triplet configuration has parameters:

\[
L_1 = L_2 = L_3 = f_2 = f_3 = f_{\text{min}}, \quad \text{and} \quad f_1 = 2f_{\text{min}}.
\] (4.5)

Note that if Eqs. (4.1)-(4.3) were solved for \( f_3 \) and the assumption \( f_3 = f_{\text{min}} \) is made, the maximum convergence angle still occurs with the parameters of Eq. (4.5), showing this (that is, a \( 2f, -f, +f \) asymmetric triplet lens configuration) to be the optimal design.

### 4.1.2 Simulations using Trace3D

The beam dynamics simulation codes Elegant [79] and Trace3D [80] were employed to study the initial lattice design, sensitivity to errors, as well as for online tuning of the focusing system during beam runs.

We employ the final electron beam focusing magnet triplet configuration suggested by the thin-lens analysis of the previous section: Focus-Defocus Defocus-Focus Focus (F-DD-FF), where the double strength (half-focal length) lenses are created by concatenating two identical magnets of the type discussed in Chapter 3. The double-length magnets are moved in tandem, to act as a movable single lens. The asymmetric triplet is first studied using an envelope code based on the first order matrix particle transport simulation code: Trace3D. The optimum lattice configuration was thus studied with this code to establish tuning conditions and performance predictions across a wide range of beam energies. Further, as is discussed below, the code is employed as a tuning guide in the experimental application of the focusing system.

As stated above, the Trace3D code was utilized to demonstrate the focus-ability of the wide beam energy range through dynamically configuring the drift
Figure 4.3: Trace3D simulations demonstrating the tunability of the final focus system. Beam energy: 72 MeV (top) and 92 MeV (bottom). The horizontal axis corresponds to the longitudinal beamline and the vertical axis for the matching $\beta$-functions. The final $\beta$-functions is in a 3-5 mm range.

spaces between PMQ lenses. Two examples of such simulation studies are illustrated in Figure 4.3 for two different beam energy settings: 72 MeV and 92 MeV. The following characteristic were simulated: rms emittances of $\varepsilon_x=7$ mm-mrad, $\varepsilon_y=15$ mm-mrad and $\beta$-functions of $\beta_x=2.0$ m, $\beta_y=0.9$ m. Trace3D as in Figure 4.3, shows that the magnet spacings $L_1, L_2$ and $L_3$ between quadrupole magnets expand with increasing electron beam energy, in order to increase the focal power of the triplet as it is applied to a more rigid beam.

Lower beam energy operation, below 60 MeV, is not possible in a system where 560 T/m PMQs are utilized. To access this lower energy range one may employ a weaker focusing lens configuration. The following system has been recently implemented for lower energy (20-35 MeV) THOMSON experiment mode:
weak Focus-strong Defocus-weak Focus (WF-SD-SF). Here the weak PMQ is achieved by boring out an existing PMQ to a larger inner radius, which produces a measured field gradient of 270 T/m. In addition to replacing the first PMQ with a weaker PMQ, the two following lenses consist of a single PMQ as opposed to a double PMQ in the original configuration. Trace3D simulations which show tunable focusability down to 30 MeV have been recently verified experimentally.

Detailed studies, with a sophisticated higher order matrix calculation code, Elegant, of effects on the performance of the PMQ final focus system in the presence of both mechanical and magnetic errors are presented in the following section.

4.1.3 Simulations using Elegant

In addition to the magnet manufacturing errors discussed above, there are several types of intra-triplet PMQ alignment errors which may have profound effects upon the performance of the PMQ final focus system. The two most serious are rotated (skewed) PMQs and transverse magnetic center offset error.

Understanding the effects of these errors requires use of a more powerful tool for analysis, the higher-order envelope and tracking code Elegant. The effects of these errors are examined here separately. Solutions for mitigating the effects of the errors in the final focus system’s overall mechanical design are discussed below which effectively correct errors to within tolerances allowed.

A systematically rotated field in a quadrupole magnet is expected when the 4-fold symmetry of the quadrupole is broken by an error in manufacturing process, for instance of un-even wedge sector shapes and incorrectly magnetized blocks. In addition, errors in angular placement of the PMQs with respect to the beamline may also produce a skewed magnet.
A simple physical model can be constructed to study the effects of magnet rotation error. The rotation angle error $\phi$ around the $z$-axis is equivalent to a rotation of the beam axes. The transverse space-coordinate rotation is applied through a matrix written as follows [56]:

$$\vec{R} = \begin{pmatrix} \cos \phi & \sin \phi \\ -\sin \phi & \cos \phi \end{pmatrix}. \quad (4.7)$$

In the simulation code, the skew rotation term is individually introduced in quadrupoles. In our study, we outline a tolerance envelope in rotation by applying a systematic error. A clockwise rotation error is applied to a focusing quadrupole and a counter-clockwise rotation error applied to a defocusing quadrupole. Based
on the simplified model constructed, Elegant simulations are generated for three
different cases of magnet rotation angle errors: 0, 0.01 and 0.02 radians. The nor-
malized emittances in both transverse dimensions display growth proportional to
the rotation error, as shown in Figure 4.4. The final rms beam spot size grows
from 20 to 34 \( \mu m \) when a 0.01 rad systematic rotation error was introduced. The
beam spot size produced of course grows progressively worse with increasing ro-
tation; in the case of 0.02 rad rotation error, a 200% growth in final beam spot
size is found.

In actuality, all of the PMQs shared the same magnet characteristics, as they
are all created by slicing each magnet from a long assembly. Thus, a rotation
error due to a segment or magnetization variability found in one particular magnet
is equally found in other magnets and therefore an azimuthal rotation error is
essentially eliminated from the focusing system. Any other rotation errors can
be also introduced through incorrect mounting of the lens assembly. Therefore,
considerable emphasis was placed on mechanical tolerances. In particular, magnet
holders for enclosing PMQs are specifically machined to comply with less than
0.01 radian angular error.

The electron beam dynamics in the presence of displacement error in PMQ
center with respect to beamline axis is also studied in the simulation code. As
noted in the magnet simulations of magnet easy-axis orientation and wedge errors,
the magnetic center is shifted from the nominal geometric center of PMQ due
to the cumulative effect of magnet errors. When such PMQs are installed into
the translation rail of the final focus system, magnet centers of the focusing
and the defocusing PMQs will be in offset positions due to 90 degrees rotation
relation. The effects of quad misalignments are to produce steering; if the vertical
midplane of a quadrupole is displaced vertically by \( \delta y \), it will introduce a dipole
field \( \delta B_z = B' \delta y \) [56]. Furthermore, this steering produces momentum dispersion, which leads to degraded performance. Determining the acceptable tolerance in magnet displacement places demands on the types of adjustments one needs in mechanical design of magnet holders. These tolerances are established through simulation. In practice, with magnetic center measurement data, corrections are applied through adjustments, after centerline measurement of each magnet, of the PMQ holder.

To simulate misalignments, the PMQs are simply specified to have transverse misalignments in the Elegant input deck. The focusing PMQs are assigned a vertical displacement error and the defocusing PMQs the same horizontal displacement. The following values in displacement error are used: 0, 0.1, 0.2 and 0.4 mm. The displacement error’s effect on the horizontal normalized emittance
is less evident when compared to the vertical emittance as the beam encounters a greater total misalignment in this direction, as shown in Figure 4.5. The emittance growth due to misalignment is less severe than that due to the skew rotation error, because it is driven mainly by momentum dispersion, which is a small effect for the design momentum spread of 0.5%. The emittance growth in vertical term is near negligible up to 0.1 mm misalignment, growing to 8% at a larger 0.4 mm displacement from the centers.

As noted in the last chapter, the use of short focal length magnets in the final focus system is in fact motivated by the need to evade chromatic effects. We now discuss the aberrations expected due to momentum errors. The focusing error from chromatic aberration arises due to the fact that the focusing strength is inversely proportional to the particle momentum. Thus, particles with more momentum will focus downstream of particles with a nominal beam energy, and this chromatic aberration causes a blurring in the beam spot size. A severe case of a beam demagnification is un-desirable at PLEIADES as the x-ray flux is inversely proportional to the (final beam size)\(^2\) at the interaction. The effects of beam momentum spread effect on the final focused beam spot have been simulated with Elegant, and are shown in Figure 4.6.

As mentioned in Chapter 3, the chromatic aberration becomes the dominant factor contributing to the final beam size as the input beam size \(\sigma_0\) is expanded. Since Eq. (3.1) was derived as a generalization to a thick triplet lens, as opposed to the original first-order, thin lens treatment found in Ref. [66], a simulation study of its validity was performed using Elegant. The results of this analysis are plotted in Figure 4.7. In this study, we used the original design emittance of the PLEIADES injector, so the overall scale of the achievable spot sizes is smaller.

As can be seen in Figure 4.7, chromatic aberration dominates the initial beam
Figure 4.6: Elegant simulation of chromatic effect: $x$ phase space diagrams at the final focal point (where the color code indicates the intensity); no energy dispersion (top) and chromatic aberrations (bottom).
Figure 4.7: Study of chromatic aberration effects in PLEIADES PMQ final focus system, using Elegant simulation (dots). Also, the line shown is the theoretical prediction.

emittance above the optimum beam input size \( \sigma_0 > \sigma_0 \vert_{opt} \) where

\[
\sigma_0 \vert_{opt} = \sqrt{\frac{\varepsilon_{x0}}{\sigma_{p}/p}} f. \tag{4.8}
\]

The corresponding minimum final beam spot size obtained in the presence of the chromatic aberrations is

\[
\sigma^* \vert_{min} = \sqrt{2f} \sqrt{\frac{\varepsilon_{x0} \sigma_{p}}{p}}. \tag{4.9}
\]

The simulations shown in Figure 4.7 assume a round input beam with initial \( \beta \)-functions in the range \( 2 \text{ m} \leq \beta \leq 250 \text{ m} \), 0.6% fractional rms momentum spread and \( 3.6 \times 10^{-2} \text{ mm-mrad} \) rms geometric emittances. The Elegant results are plotted against the prediction of Eq. (3.1). In this way, the effective focal length of the final focus system in the \( x \)-dimension was determined to be \( f_{eff}=6.2 \text{ cm} \). The predicted optimum input \( \beta \)-function for the minimum final beam focal spot

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size of 5 μm is 10 m for this case, which has a relatively small emittance. For the actual experimental parameters with degraded emittance, the initial optimum $\beta$ is smaller.

4.2 Beam Based Measurements

4.2.1 Installation at PLEIADES

As explained in the Chapter 2, the PLEIADES electron beam is generated using a BNL/SLAC/UCLA style 1.6 cell photo-cathode RF gun [81]. Following the gun, the beam is accelerated to energies of 20-100 MeV using as many as four S-band 2.5 m traveling wave sections. For the measurements described below, the gun produced 200-300 pC bunches with 3 ps rms duration, while the final beam energy was varied between 50 to 70 MeV, and the typical normalized, rms emittance was around 10 mm-mrad.

To install the PMQs as the e-beam final focus system, the high energy electron beamline was redesigned both to provide approximate beam matching into the PMQs and to retain diagnostic capabilities. A schematic of this beamline is shown in Figure 4.8.

The laser-electron IP (interaction point) is located at the center of a 6-inch cube vacuum chamber. As described above, the PMQ mover assembly is mounted directly into the interaction chamber, allowing positioning of the magnets at a range of distances directly upstream of the IP (see Figure 4.9). The minimum distance of the final PMQ to the IP is determined by the width of a polished aluminum cube positioned at the IP. This cube functions as the spatial and temporal overlap diagnostic by directing the laser light and the optical transition radiation (OTR) produced by the electron beam into a CCD camera and a streak
Quadrupoles (emittance and beam matching)
Dipole/Spectrometer (energy spread)
Compton scattering interaction point
PMQ final focus system
CCD for x-ray detection
FALCON laser 800 nm, 1 J, 50 fsec
Steering magnets

Figure 4.8: Schematic drawing of the PLEIADES interaction beamline.
camera. In addition to enabling the overlap of the two beams, the CCD images naturally allow optimization of the focused spot size. This cube is extracted from the beam path for x-ray production but cannot be fully removed from the PMQ path. As a result, the minimum distance from the edge of the final PMQ to the IP is 5 mm.

After the IP, a round pole dipole magnet is used to bend the electron beam onto a trajectory which separates it from the x-ray pulse, allowing the beam to be dumped into a shielded Faraday cup, minimizing the noise reaching the x-ray detectors. Note that a spectrometer magnet was also located just upstream of the final matching quads to measure the energy and energy-spread of the beam, which are essential steps in tuning the linac. A typical energy spread measurement performed with a dipole spectrometer is shown in Figure 4.10.
Figure 4.10: An energy spread of 0.2% is normally expected at PLEIADES. The measurement is performed with upstream dipole spectrometer

4.2.2 Beam Diagnostics

Beam imaging diagnostic is an essential part of the accelerator facility as an image measurement tool for profiling beam from the moment it is emitted from the photoinjector cathode to all the way down to the final focal spot. Also, the quad scan necessary to experimentally determine the beam Twiss parameters heavily depends on the success of obtaining clear transverse beam images. In addition, various matching conditions in beam lattices essential for the preparation of an experimental run requires observing beam status in real time. Another important aspect of beam imaging diagnostic is that it gives a linac operator with necessary means to optimally place the off-axis beam right in the center of the beamline with steering magnets such that the probability of beam hitting accidentally into the wall of beamline is avoided. Especially, this is very important for the successful operation at PLEIADES in regarding the clean detection of x-ray free of undesirable bremsstrahlung radiation as by-product of interactions between ultra-relativistic electron beam and metal.

During the reconfiguration of the beamline to accommodate the installation
of the PMQ final focus system, beam imaging diagnostics were setup along the various points at upstream of IP. Each diagnostic is equipped with a pop-in actuator with a phosphor screen attached at the tip of the actuator rod is inserted into the beam path and at a skew angle. In the process of setting up imaging diagnostics, a HeNe laser was used as a reference guide line of camera sight alignment. When electron beam hits the scintillator, light is generated and exited properly through the vacuum port window at 90 degrees angle from the beam axis. To view this scintillation light, a COHU 4910 CCD camera equipped with properly chosen focal length lens is stationed just outside the viewport window, which then captured image is relayed to the monitor screen in live.

OTR signal emitted from the aluminum cube evades the issue with the phosphor screen’s collection of dark current along with the beam itself. In some cases, due to integrating nature of the phosphor, photo-electrons are overwhelmed by such dark current effect to create a difficult situation to obtain real beam image. Moreover, OTR resolution degradation from admittance and inherent divergence as a minor effect, in contrast with the Čerenkov radiation produced from charge particle moving at velocities greater than the phase velocity of light in the medium such as quartz [82]. The success of PLEIADES is depended on the high-resolution beam image, which can be achieved with relatively simple OTR setup. The transition radiation as a broad spectrum emission in the form of highly narrow light beam is a consequence of re-organization and detachment of charged particle field in the process of crossing the boundary between two different media. The detection of optical frequencies is easily accomplished with commercially available CCD camera, hence the name optical transition radiation (OTR) is originated. For OTR to generate appreciable photon lights, phase coherence between driving fields in medium 1 and generated wave in medium 2 is desirable. This coherence condition is met if where medium 1 is vacuum such that $\varepsilon(\omega)^{1/2} \gamma \theta \leq 1$ [82]. In all

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cases, a metal is chosen so that frequencies in optical range \((10^{15} \text{ Hz})\) are emitted backward into vacuum. In optical frequency range which is an order magnitude smaller than the plasma frequency \((10^{16} \text{ Hz})\), OTR radiation penetration depth is shallow into metal, therefore OTR has 100% tendency to re-emerged into vacuum.

In general, simultaneous observation of electron beam and laser beam and most importantly carrying out spatial-temporal alignments require use of aluminum cube with ultra-polished flat smooth, planes. In addition, this cube has an edge sharpness about a few tens \(\mu\text{m}\) corner radius, which enables linac operator to overlap beams in spacetime with very high accuracy. An aluminum metal, a cube material, is chosen based on its high reflectivity to IR laser beam and high OTR intensity from electron beam. Aluminum cube position is controlled with 3-axis stepper motor with coordinate information readily available at the operation station. A complete setup of IP alignment diagnostic is shown in Figure 4.11. As shown in the diagram, cube is mounted with its planes in 45 degrees orientation with respect to the respective initial beam directions and its conjoining edge lined in with the camera perspective.

### 4.2.3 Tuning the Final Focus Optics

In accordance with the quad triplet model used to design the PMQ final focus system, the beam must be manipulated up-stream of the PMQs to produce a moderately focused waist at the entrance to the first PMQ. The first step in this process is the determination of the rms Twiss parameters of the beam exiting the accelerator, which is done using the quadrupole scan emittance measurement technique. In the quad scan measurement, the current in the first electromagnetic (EM) quad in Figure 4.8 is varied while the rms beam size is measured after a known drift length. Taking into account the finite length of the quadrupole
Figure 4.11: Configuration of beam diagnostic at IP, which enables simultaneous observation of laser and electron beam in real time.

magnet (using a “thick lens” analysis) the final beam size in each transverse direction, $\sigma_f$, follows the well-known relation [57]

$$
\begin{align*}
\sigma_f^2 &= \left[ \cos(\sqrt{K}l_q) - \sqrt{K}l_q \sin(\sqrt{K}l_q) \right]^2 \sigma_{11,i} \\
&+ 2 \left[ \frac{\sin(\sqrt{K}l_q)}{\sqrt{K}} + l_d \cos(\sqrt{K}l_q) \right] \left[ \cos(\sqrt{K}l_q) - \sqrt{K}l_d \sin(\sqrt{K}l_q) \right] \sigma_{12,i} \\
&+ \left[ \frac{\sin(\sqrt{K}l_q)}{\sqrt{K}} + l_d \cos(\sqrt{K}l_q) \right]^2 \sigma_{22,i},
\end{align*}
$$

(4.10)

where $l_{q,d}$ are the EM quad and drift lengths, respectively, $K$ is the focusing strength of the quad given by its magnetic field gradient and the beam momentum, $K = eB'/p$, and the initial Twiss parameters are $\beta_i = \sigma_{11,i}/\varepsilon$, $\alpha_i = -\sigma_{12,i}/\varepsilon$, and $\gamma_i = \sigma_{22,i}/\varepsilon$, where $\varepsilon$ is the rms emittance. Because quad scans must be performed as part of the routine operating procedure, the process was automated using LabVIEW based control system software. The immediate determination of the emittance provides the linac operator with feedback on the quality of the
Figure 4.12: Emittance measurement for a 300 pC, 60 MeV bunch using a quad scan method.

linac tune, as well as giving the Twiss parameters needed for matching the beam into the PMQs. The emittance measurements using EM quad scan technique were taken, where from data shown in Figure 4.12 for a 300 pC, 60 MeV beam, a normalized emittance of 15 mm-mrad was obtained.

Once the Twiss parameters have been determined, they are used as input to a Trace3D model of the beamline. Given the beam parameters at the first element of the model, and those desired at some latter point, Trace3D can perform a matching function, in which a user defined set of model parameters are adjusted in an algorithm to produce the desired final beam. In this case, we adjust the strengths of EM quad doublet and triplet to produce a beam waist ($\alpha = 0$) in both transverse dimensions at the entrance to the PMQs with an rms size around 350 $\mu$m. An example of the output of this matching process is shown in Figure 4.13.
Figure 4.13: Trace3D model of the beamline used to compute the focusing strength of the upstream electromagnetic quadrupoles for proper matching into the PMQs. The values encircled in a circle line are calculated quad strengths in the upstream EM doublet and triplet.

The choice of rms beam size at the final focus entrance is a compromise between emittance and aberration considerations, and will be discussed further below.

Once the Trace3D model has determined the proper settings for the upstream quadrupoles, measuring the beam size at two pop-in beam profile monitors and comparing the results with the model predictions allows a partial check of the model. With good agreement found between model and measurement, the Trace3D model is used again to determine the proper spacing of the PMQs. The result of this matching process is shown in Figure 4.14. Here the spacing between...
Figure 4.14: Typical Trace3D output specifying the spacing between PMQ blocks and their position with respect to the beam focus. The values encircled in a circle are calculated drift spacings between PMQs.

3 PMQ holders is adjusted to produce the smallest $\beta$-function possible at the IP.

The linac operator uses the Trace3D values of upstream quad strengths and PMQ positions as a starting point. One important practical consideration the operator must take into account is steering the beam onto the axis defined by the PMQ mover system. This is needed both to minimize the effects of chromatic and spherical aberrations in the PMQs and to precisely define the direction of the Thomson x-rays generated. There are two steering magnets located between the EM quad triplet and the PMQs that are used for this job. Steering is also made easier by imaging the beam on a green emitting plastic scintillator (EJ-260
Figure 4.15: Electron beam focused with 15T/m triplet quads. At optimum performance, a 50 µm beam spot was obtained at the IP.

produced by Eljen Technology) mounted on the upstream face of the first PMQ holder. This scintillator has a 5 mm diameter hole that is mounted concentric to the PMQ aperture. Although, YAG based scintillator is superior in terms of light output generated, a solid thin-plate with a specified drilled hole is easier to produce in EJ-260 and permanently put in the beamline without compromising the vacuum level. With these tools, proper beam steering is accomplished using the two steering magnets to iteratively move the beam to the center of the scintillator hole and observe the motion (or lack of motion) of the beam at the IP while moving the final PMQ holder.

The process described above has been used to focus the PLEIADIES electron beam at the IP using the PMQ system. The spot size was measured using OTR generated by the beam colliding with the interaction cube diagnostic as described above. This light is collected by a CCD camera/lens system and digitized using
a computer frame grabber. A typical and highly repeatable beam spot is shown in Figure 4.16. In this case we focused a 70 MeV, 200 pC beam to less than 20 µm rms size in both transverse dimensions. This is a significant improvement over the previous final focus system, which achieved spot sizes of roughly 50 µm, rms. As a comparison of how much has improved before and after PMQ shows in the Figure 4.15 a 50 µm transverse beam achieved with 15 T/m EM triplet quads. The rms beam emittances $\varepsilon_x = 5$ mm-mrad and $\varepsilon_y = 13$ mm-mrad and the $\beta$-function of 1 cm were measured for that particular beam measurement.

4.2.4 Emittance Measurements with the PMQs

The standard quad scan employing the variation of quad strength has been normally performed upstream to obtain electron beam emittance. We have performed the different variation of quad scan at the interaction point to test whether a beam emittance can be obtained confirming the standard quad scan technique. A variation of the quad scan technique is the so-called three-screen or moving
screen measurement, in which there is a focusing lens followed by a drift space and measurement screen. In this case the beam size is measured while varying the drift length instead of the lens strength. In the simplified thin lens analysis (where the focal length is $f = 1/Kl_q$) the equation analogous to Eq. (4.10) is

$$
\sigma^2(z) = \beta^* \varepsilon_0 \left[ 1 + \left( \frac{z-z_0}{\beta^*} \right)^2 \right] + \sigma_{res}^2,
$$

(4.11)

where $\beta^*$ is the minimum $\beta$-function at the waist, $\varepsilon_0$ is the beam emittance, and $\sigma_{res}$ is the image resolution obtained from the limitations of optical transport and camera performance. In the free drift, the square of the rms beam spot size expands quadratically away from the beam waist position, $z_0$.

In the case of the PMQ final focus system, the focal length, $f$, is fixed and the drift length, $z$, is changed most easily by moving all three of the PMQ holders in unison. This is not technically a correct application of Eq. (4.11), since by moving the quadrupoles instead of the screen, we are varying the initial Twiss parameters. However, since we know the initial $\beta$-function of the input beam to be at least 1 m, and since the scans must only cover several $\beta^*$, or less than one cm, the initial conditions remain essentially constant over the scan.

Thus, the analysis of the moving screen technique to the “moving quads” measurement is valid in the present scenario. One such dataset and analysis is shown in Figure 4.17. The effects of limited image resolution can be clearly seen in these measurements. The expected parabolic dependence of $\sigma^2$ on $z$ is observed well away from the minimum beam sizes, but the data give a nearly flat dependence in the zone near $z_0$. In the case studied in Figure 4.17, normalized initial emittances measured at 58.2 MeV beam energy using the EM quad triplet are $\varepsilon_x = 16.6$ mm-mrad and $\varepsilon_y = 9.7$ mm-mrad, while the normalized emittances obtained with the PMQ scan are $\varepsilon_x = 16.3$ mm-mrad and $\varepsilon_y = 12.9$ mm-mrad. During the initial period of running the final focus system, beam spot-sizes as low
Figure 4.17: Quadrupole scan of the final focus beam at LLNL PLEIADES, performed by moving the entire PMQ assembly.

As noted, this is an upper limit on the actual beam size, possibly due to the resolution limitations in the measurement system. We note from Figure 4.17 that the value of the $\sigma_{\text{res}}$ obtained from the best fit of the quad scan data to Eq. (4.11) is also 18 $\mu$m. With the value of $\sigma_{\text{res}}$ determined by the fit to the data, we deduce that the minimum beam size in each transverse dimension in Figure 4.17 would have been in the range of 15 $\mu$m in the absence of image resolution problems (e.g. geometric and chromatic aberrations).

One of the most critical needed capabilities of the PMQ-based final focus
system, that of tunability over a wide range of beam energies, has been demonstrated. Through longitudinal motion of the PMQs, the final focus system was capable of adjusting to consistently produce a small beam spot-size at different input beam energies. This variability, over 40% in beam energy, has produced variation of x-ray photon energies over the range between 40 to 140 keV.
CHAPTER 5

Linear Thomson X-ray Experiment

The current chapter will deal with the making of ultra-short, high-brightness, quasi-monochromatic high-energy x-rays via intense scattering interactions between a high-power, IR laser and a high-density, ultra-small electron beam. The pre-experimental work in devising the method to minimize both the spatial and temporal jitters to maximize the interaction will be discussed in conjunction with the time-domain analysis, which details the determination of the bound-limits in the tolerance of jitters. The detailed experimental set-up for the head-on collision geometry will also be given in the current chapter. Finally, the chapter will close with the discussion of the measurement results in x-ray flux, energy spectrum and brightness achieved at the PLEIADES laboratory.

5.1 Electron-Laser Beam Interactions: Linear Thomson Scattering

While the side-on (the 90 degrees) collision produces ultra-short x-rays in the sub-ps timescales, the flux generated is far below that of the 3rd generation synchrotron source. This is due to the number of particles involved in the scattering interaction, which depends on the length of the interaction time. In this case, the time of interaction is very small, because an incident laser sees only the narrow cross section of the electron waist. At PLEIADES, the head-on scattering
geometry is chosen so that the interaction time taken is over the whole length-
scales of the beam longitudinal profiles yielding interactions to be held at the
maximal level. On the other hand, for the x-ray energy spectrum, it is of interest
to extend the scattered photon spectrum above \( \geq 100 \, \text{keV} \). The energy scale in
this range with the low-energy accelerator is easily achievable by employing and
benefiting from the head-on scattering geometry where the energy spectrum is
\( E_x = 4\gamma^2 E_{\text{inc}} \), which is by a factor of two larger than the 90 degrees interaction
geometry. In this way, x-ray energies in the 10 - 100 keV range can be produced
with the existing 100 MeV linear electron beam accelerator at LLNL.

One of the main objectives of the PMQ final focus system in the experiment is
to upgrade the rate of photon production by creating a denser region of electron
particles at the point where the laser focal spot is. In addition, the spatial-
temporal jitters must be contained and controlled. The relating parameters of
interest to Thomson photon yields is studied under the model of idealized time
dependent particle distributions. In conjunction with the theoretical model pro-
posed, qualitative discussions of experimental observations are reviewed and the
prior steps necessary to perform in setting desirable experimental conditions are
outlined. Later on we will see that the model serves to provide the lower bound
timing synchronization parameter in the laser transport delay system. Another
important aspect is the x-ray photon brightness, and its relation to the spot size
and angular divergence of the electron and laser beam. This is all to be discussed
in the current chapter with some attempts made to correlate this model with
experimental observations.
5.1.1 Time-domain Analysis

In Chapter 1, we have stated that the total x-ray photon produced can be expressed as the integral of convolution of two beam particle density functions over the spacetime continuum. But the specific density functions, which closely model respective beam profile over the time-domain have not been defined. In order to characterize the production of Thomson photons, which under various instances of density behaviors would give the best descriptions of observations and expectations, we need to devise correct forms for corresponding density functions.

Prior to defining proper density functions with respective laser and electron beams, a unified coordinate system for the Thomson interaction of particle distributions is necessary. Let the electron beam density be given by $n_e(\vec{r}_e, t_e)$, and the laser density respectively by $n_\gamma(\vec{r}, t)$, where each distribution is uniquely described in its own separate coordinate systems. We will redefine the space-time coordinate system of laser beams in terms of $(\vec{r}_e, t_e)$ so that it will unify into the electron’s rest frame. We assume that the transverse beam distribution holds a perfect axial symmetry about the longitudinal axis. The unified coordinate system in the cylindrical coordinate format will be assumed as follows:

\[
\begin{align*}
x_e &= x \cos \theta + z \sin \theta, \\
y_e &= y, \\
z_e &= -x \sin \theta + z \cos \theta, \\
t_e &= t,
\end{align*}
\] (5.1)

where the $\theta$ is a collision angle between the two incident beams. The electron radial coordinate is given in terms of laser photon coordinates using above relations:

\[ r = \sqrt{x^2 + y^2}. \]
\[ r_e = \sqrt{x_e^2 + y_e^2} = \sqrt{(x \cos \theta + z \sin \theta)^2 + y^2}. \] (5.2)

In this coordinate system, we assume that the interaction occurs in \(x\)-\(z\) plane with arbitrary scattering angles defined by \(\theta\). With this unified coordinate system, various different scattering scenarios, such as a side-on, a head-on or any intermediate angle, can be studied.

As described in Chapter 2, both laser and electron beams are originated from a single oscillator pulse, so their beam distributions are nearly identical to each other. For PLEIADES, beams fairly resemble the Gaussian distribution; similarly the mathematical density model we chosen will be described by a Gaussian function. We will begin with the collision occurring in the counter-propagation mode in order to correlate closely with the real experimental situation at Livermore. Then, we will give the full generalization of the collision process of particle beams at any angular degrees.

A good paraxial approximation of laser density model is assumed to have a cylindrically symmetric Gaussian transverse profile and a temporal Gaussian distribution as follows [39, 83]

\[
n_\gamma(r, z, t) = \frac{\eta_\gamma}{1 + (z/z_0)^2} \exp \left\{ 2 \left( \frac{t - z/c}{\Delta t} \right)^2 \right\} \exp \left\{ -\frac{2r^2}{w_0^2[1 + (z/z_0)^2]} \right\}, \tag{5.3}
\]

where from integrating overall coordinate dimensions of the photon density distribution

\[
\int n_\gamma(\vec{r}, t) d^3 \vec{r} = N_\gamma,
\]

we retrieve \(N_\gamma\) as the total number of photons in a single laser pulse (i.e. \(E_L/\hbar \omega_L\)). The rest of the important parameters in Eq. (5.3) are a focal waist \(w_0\), a Rayleigh length \(z_0 = \pi w_0^2/\lambda_0\), and a pulse duration \(\Delta t\).
Under the assumption that an average laser pulse duration is very short compared to the transit time through a Rayleigh length ($\Delta t \ll z_0/c$), a time-dependent Gaussian function in Eq. (5.3) can be approximated by a Dirac delta-function. In the short-pulse duration limit, a mathematical transformation is given as follows

$$\exp \left\{ -2 \left( \frac{t - z/c}{\Delta t} \right)^2 \right\} \rightarrow \frac{c\Delta t\sqrt{2\pi}}{2} \delta(ct - z).$$

Now, with the Dirac function defined in Eq. (5.3), the normalization factor is obtained from an evaluation of the following integral

$$\eta_{\gamma} c\Delta t\sqrt{2\pi} \int_{-\infty}^{+\infty} dz \frac{\delta(ct - z)}{1 + (z/z_0)^2} \int_0^\infty 2\pi r dr \exp \left\{ -\frac{2r^2}{w_0^2[1 + (z/z_0)^2]} \right\} = \eta_{\gamma} w_0^2 c\Delta t \sqrt{\pi/2}.$$

The integration result as shown on the right hand side is equal to the total number of photons, so a proper normalization of the integral is therefore $\eta_{\gamma} = N_{\gamma}/w_0^2 c\Delta t \sqrt{\pi/2}$.

Similarly, an electron density distribution also takes the same cylindrical Gaussian distribution properties. Except that a propagation direction is taken to be in the $z_e = -z$ ($\theta = 180$), which is an opposite of the propagation direction of a laser. An electron density distribution is written [39, 83]:

$$n_e(r, z, t) = \frac{\eta_e}{1 + (z_e/\beta_e)^2} \exp \left\{ -\left( \frac{t - z_e/v}{\Delta \tau} \right)^2 \right\} \exp \left\{ -\frac{r_e^2}{\sigma_e^2(1 + (z_e/\beta_e)^2)} \right\},$$

(5.4)

and the proper normalization calculated is given by

$$\eta_e = \frac{N_e}{\sqrt{\pi}\sigma_e^2 v \Delta \tau} \approx \frac{N_e}{\sqrt{\pi}\sigma_e^2 c \Delta \tau},$$

where the last approximation is that the beam velocity is taken to be equal to $c$.

The beta-function of an electron beam, which is characterized by a convergence angle $\sigma_e'$, is defined by

$$\beta_e = \frac{1}{\tan(\sigma_e')} \sigma_e \approx \frac{\varepsilon_n}{\gamma_0 \sigma_e^2},$$
where the convergence is related to the normalized emittance, \( \varepsilon_n = \gamma_0 \sigma_e \sigma_e' \). The \( \gamma_0 \) is a familiar relativistic factor.

While the particle density functions defined above for an incident laser and an electron beam are useful for studying the rate and the total number of x-rays produced, they do not provide information regarding the charged particle dynamics in the electromagnetic field complex and or the resulting characteristics of Thomson scattering process as was done with the Lienard-Wiechert potential described in Chapter 1. Rather, a primary goal here is to show spatial and temporal jitter effects on the Thomson x-ray flux produced.

The rate of photons scattered is the time differentiation of \( N_x(t) \) \[84\], which is expressed by

\[
\frac{dN_x}{dt} = c \sigma_T \int n_e(r,t)n_e(\gamma, t)d^3r.
\] (5.5)

And by the substitution of beam density functions, we have the integral as follows

\[
\frac{dN_x}{dt} \approx c \sigma_T \eta_e c \Delta t \sqrt{2\pi} \int_{-\infty}^{+\infty} dz \int_0^{\infty} 2\pi r dr \frac{\delta (ct - z)}{1 + (z/z_0)^2} \exp \left\{ -\frac{2r^2}{w_0^2[1 + (z/z_0)^2]} \right\} \\
\times \left\{ -\frac{1}{1 + (z/\beta_e)^2} \exp \left\{ -\left( \frac{t + z/v}{\Delta \tau} \right)^2 \right\} \exp \left\{ -\frac{r^2}{\sigma_e^2(1 + (z/\beta_e)^2)} \right\} \right\},
\] (5.6)

The radial integral is first computed as follows

\[
\int_0^{\infty} r dr \exp \left\{ -r^2 \left[ \frac{2}{w_0^2(1 + (z/z_0)^2)} + \frac{1}{\sigma_e^2(1 + (z/\beta_e)^2)} \right] \right\} = 1 \cdot \frac{1}{w_0^2[1 + (z/z_0)^2]} \sigma_e^2 \frac{1}{1 + (z/\beta_e)^2}.
\]

After the radial integral part has been evaluated, the final integration is easily performed over \( z \), and the number of photons scattered per second in the head-on collision geometry is

\[
\frac{dN_x}{dt} = 2 \sqrt{\pi} \frac{N_e N_\gamma}{\Delta \tau} \exp \left\{ -\left[ \frac{t(1 + c/v)}{\Delta \tau} \right]^2 \right\} \frac{\sigma_T}{2 \sigma_e^2(1 + (ct/\beta_e)^2) + w_0^2[1 + (ct/z_0)^2]}.
\] (5.7)
A simple observation can be drawn from the result with respect to an electron pulse duration: the rate of Thomson x-rays produced is inversely proportional to $\Delta \tau$ so that a shorter electron beam produces more x-rays. In conjunction with the Thomson program, the velocity bunching technique was developed to produce femtosecond-range electron beams and is reported in [60].

By performing an integration of Eq. (5.7) overall time-domains, we obtain an equation for the total number of x-rays generated as

$$N_x = \sqrt{\frac{2}{\pi}} N_e N_\gamma \frac{\sigma_T}{2\sigma_e^2 + w_0^2}. \quad (5.8)$$

From this result, the importance of the electron beam spotsize is quite clear. The total x-rays produced could be increased easily by achieving a smaller electron spotsize at the interaction point. Hence, the justification for employing strong short focal length quads is demonstrated here by the expression $N_x$. However, the x-ray brightness does not necessary benefit from the spotsize alone, as it also depends on an angular divergence of the electron beam.

Rather than confining our discussion on various jittering effects to 180 degrees
scattering geometry alone, we can extend the result to all scattering angles. Then, with the generalized equation of the rate of x-rays produced, we study jitter effects in the flux scattered from a head-on and a side-on scattering geometries.

The generalization of Eq. (5.6), so that a result at any scattering angle can be obtained is accomplished by using the electron coordinates as defined in Eq. (5.1). The rate of x-ray flux produced is then in terms of these variables given by

\[
\frac{dN_x}{dt} = \eta_e \eta_c \sigma_T c \int_{-\infty}^{+\infty} dx \int_{-\infty}^{+\infty} dz \frac{\delta(ct-z)}{1 + z^2/z_0^2} \exp \left[ -\frac{2x^2}{w_0^2(1 + z^2/z_0^2)} \right] \\
\times \frac{1}{1 + (z_e/\beta_e)^2} \exp \left[ -\left( \frac{t - z_e/v}{\Delta \tau} \right)^2 \right] \exp \left[ -\frac{x_e^2}{\sigma_e^2(1 + (z_e/\beta_e)^2)} \right] \\
\times \int_{-\infty}^{+\infty} dy \exp \left\{ -y^2 \left[ \frac{2}{w_0^2(1 + z^2/z_0^2)} + \frac{1}{\sigma_e^2(1 + (z_e/\beta_e)^2)} \right] \right\}.
\]

Because of the Dirac delta function, the integral over the \( z \) component replaces all \( z \) with \( ct \), and the integral over the \( y \) component can be performed using the well-known integral identity: \( \int_{-\infty}^{+\infty} e^{-qx^2} dx = \sqrt{\pi/q} \). The remaining integral over the \( x \) component is too complicated to be evaluated, so we leave it in the
non-closed integral expression:

\[
\frac{dN_x}{dt} = \frac{2}{\pi^2} \frac{N_e N_e \sigma_T}{\Delta \tau} \frac{1}{w_0^2} \int_{-\infty}^{+\infty} \frac{dx}{\sigma_e} \frac{1}{\sqrt{1 + \left(\frac{z_e}{\beta_e}\right)^2}} \frac{1}{\sqrt{1 + \frac{2\sigma_e^2}{w_0^2} \left(1 + \frac{z_e}{\beta_e}\right)^2 \frac{1}{\frac{1}{1 + c^2 t^2 / z_0^2}}}}
\]

\[
\times \exp \left\{-\frac{2x^2}{w_0^2(1 + z^2 / z_0^2)} - \frac{x_e^2}{\sigma_e^2(1 + (z_e / \beta_e)^2)} - \left[ t \left(1 - \frac{\cos \theta}{\beta_e} \right) + \frac{z \sin \theta}{\Delta \tau} \right]^2 \right\}
\]

(5.9)

We verify that for \( \theta = \pi \) in Eq. (5.9), we recover the same result of the head-on collision case. Now, we are in position to examine the temporal jitter effects for the two limiting cases: 180 and 90 degrees scattering geometries.

### 5.1.2 Spatial & Temporal Jitter Effects

The successful operation of Thomson x-ray machine is based on keeping both spatial and temporal jitters between scattering beams below experimental error parameters. These jitter parameters sharply depend on the set-up of the scattering geometry. Until recently, Thomson x-rays have been mainly produced either in a 180 degrees mode or a 90 degrees mode depending on the desire to obtain the maximum number of scattered photons or the shortest x-ray duration. We will analyze the effects expected under various jitter conditions using the results that we have obtained in the last section.

In the last section, we derived the expression for the x-ray flux \( dN_x / dt \). In the expression of electron coordinates in Eq. (5.1), the time delay of beam interactions could be induced by introducing the new definition \( t \rightarrow t \pm \delta t \), where the positive sign is understood as a laser arriving at the interaction behind the electron beam and the negative sign vice versa. A similar change of the laser time definition in Eq. (5.9) gives full details of the x-ray flux effect at any given time delay.

Similarly, the spatial misalignment is introduced in the x-ray flux expression by redefining a transverse laser coordinate in Eq. (5.1) with \( x \rightarrow x \pm \delta x \). It is
<table>
<thead>
<tr>
<th>Parameter</th>
<th>Value</th>
</tr>
</thead>
<tbody>
<tr>
<td>Laser wavelength</td>
<td>800 nm</td>
</tr>
<tr>
<td>Laser pulse energy</td>
<td>200 mJ</td>
</tr>
<tr>
<td>Laser pulse duration</td>
<td>54 fs</td>
</tr>
<tr>
<td>Laser pulse size</td>
<td>20 µm</td>
</tr>
<tr>
<td>Electron beam charge</td>
<td>0.3 nC</td>
</tr>
<tr>
<td>Electron beam energy</td>
<td>30 MeV</td>
</tr>
<tr>
<td>Electron beam radius</td>
<td>20 µm</td>
</tr>
<tr>
<td>Electron beam duration</td>
<td>3 ps</td>
</tr>
<tr>
<td>Electron beam normalized emittance</td>
<td>5 π-mm-mrad</td>
</tr>
</tbody>
</table>

sufficient to induce a misalignment in one coordinate direction without having to do the same with the y-coordinate. Also, the transverse displacement of a laser beam is with respect to a fixed coordinate of an electron beam.

For the following studies of spacetime jitters in Thomson scattering, we used the beam parameters listed in the Table 5.1

In the time jitter studied, the time of arrival of a laser beam with respective to an electron beam is delayed by $\delta t$. In this way, the head of the electron beam tail sees the laser beam in the presence of a finite $\delta t$ or otherwise when $\delta t = 0$ centers of beam densities overlap in synchronous. For the head-on collision case, x-ray flux with respect to a laser time and the total time-integrated flux distributions with respect to time delay have been calculated. In Figure 5.3, x-ray flux produced from a 180 degrees scattering geometry for various $\delta t$ shows that as the time delay increases a peak amplitude of flux decreases. For instance,
for $\delta t = 1$ ps time delay, the peak amplitude is reduced almost by half that of the flux at zero time delay. Also the locations of each center peak occur at approximately 2, 4, 6, and 8 picoseconds ahead of a laser beam arriving at $t = 0$. Also, in Figure 5.3, the time integrated total x-rays plotted as a function of $\delta t$ is shown. The peak of total x-rays at zero time delay is approximately $5 \times 10^7$ photons. Gradually, the total x-ray decays over the time delay range of $-30$ ps $\leq \delta t \leq +30$ ps. At the half reduction of the peak x-ray photon numbers, the width of the curve (i.e. FWHM of $\delta t$) is approximately 18 ps, which indicates
Figure 5.4: Time-delay measurement for 180 degrees scattering geometry.

that an appreciable number of x-rays are scattered even in this large time delay. This is so because a large portion of scattering beams are still found within the Rayleigh range of a laser field, \( i.e. \), between \( t = -z_r/c \) and \(+z_r/c\).

For the 90 degrees scattering case, flux peaks are sharply peaked at a delay time with each width approximately equal to the duration of a laser beam. Also, the rate of reduction in flux is faster in this case than the 180 degrees case, as shown in Figure 5.3. The time integrated flux is also plotted as a function of the time-delay in Figure 5.3. Here, the FWHM of \( \delta t \) is approximately 1.2 ps, which is very narrow compared to that of the 180 degrees case. In addition, the peak integrated flux is found to be \( 3.6 \times 10^6 \) photons, which is about 15 times less than what was found in the 180 degrees scattering. For an appreciable Thomson scattering process to occur in this configuration, a laser beam should be present throughout the range of the whole electron beam duration.

The corresponding experimental time delay studies have been performed at PLEIADES by increasing the distance of a mirror located in the FALCON laser
compressor chamber. By adjusting the mirror mounted on a precision stepper motor stage riding on a rail, a real time delay of laser arrival time can be induced. While changing the arrival time of a laser beam, x-ray photon counts have been measured with an x-ray CCD camera. The photon count measurements were accumulated for various values of time-delay, and the data is plotted as shown in Figure 5.4. As can be seen in the data, the Gaussian curve fits well with the data in the negative delay time domain (where a laser arrival time is ahead of an electron). On the other hand the Lorentzian fits well in the positive delay time domain (opposite the negative case). The FWHM time delay width of $\sigma_t = 20$ ps is deduced, which in length is approximately 6 mm [40]. From the data, a measured FWHM width is 6.3 mm, which corresponds closely with the prediction. Also, the width of the interaction region calculated with the formula is $2z_r = 8$ mm, which is within the range of both the measured and the fitting parameters.

The spatial jitter studies show a stringent space limit for the case of a head-on collision. As stated earlier, a space displacement of a laser beam is applied in the $x$ direction, leaving the $y$ coordinate unchanged. The center beams on-axis yield a maximum x-ray flux, the same as that shown in the time jitter study, and the spatial FWHM width deduced is approximately equal to the rms spot size. But as a laser beam steers away from the interaction on-axis, the amplitude reduces and the width grows rapidly, as shown in Figure 5.5. At 35 $\mu$m transverse displacement, the x-ray flux is almost reduced to a very small number, compared to the 0 $\mu$m displacement case. The center of the x-ray beam is also shifted in the same direction as the laser beam displacement direction. From the integrated flux plotted versus the spatial displacement $\delta x$, the FWHM width is found to be approximately 30 $\mu$m, which explains why the peak flux decays rapidly. During the experimental run, tuning of the laser beam center showed evidence that the x-ray beam shifted in the direction of laser movement.
The spatial jitter for the 90 degrees scattering, on the other hand, is identical to the time-delay study of the 90 degrees case. Because in the side-on configuration a space displacement is identical in process as inducing the time-delay of a laser beam with respect to an electron beam. Therefore, the space displacement places a laser beam so that it can sample different sections of an electron beam region at each point the displacement value changes. Because the displacement occurs in a laser beam (not in the electron beam), the final position of the x-ray beam does not shift; rather one expects only the reduction in the x-ray flux. This
expectation is demonstrated in Figure 5.5. Also, the interaction region here is defined by the $\beta$-function of the electron beam, so that the FWHM of space jitter is about the length of the electron beam duration, and therefore wider in this case than that of the 180 degrees case.

5.2 PLEIADIES Experimental Set-up

The major layout of the PLEIADIES experiment was covered in Chapters 2 and 3 with an emphasis on the installation of the PMQ system, which left out discussions of the set-up of laser optics and beam optics components near the Thomson interaction region.

A high-vacuum chamber with dimensions of $32.5 \times 22.5 \times 15$ inches was used as a set-up chamber for four reflecting mirrors and a parabolic focusing mirror on an optical-mount platform. The incident IR laser, which was compressed to 54 fs long, propagates from the compression chamber in a pipeline 14 inches off-axis from the electron beamline. It then enters the vacuum folding mirror box, where the laser is reflected from three different mirrors specially designed for IR wavelength and focused by a 1.5 meter focal length parabolic mirror. The off-axis parabolic laser focusing mirror was designed to focus the laser to a 36 $\mu$m ($1/e^2$) spotsize at the IP. After the laser was focused, a final turning BK7 mirror (whose diameter is three inches) was employed to reflect the laser in the direction of the head-on collision with an oncoming relativistic electron beam. Also, this final turning-mirror, mounted on an optical mount, is equipped with stepper motors which adjust the laser spot transverse position at IP. The final-turning mirror, which is a half inch in thickness, is in the direct path of the x-ray beam. Therefore, x-rays passing through the material are expected to be absorbed, while effectively less x-rays reach the x-ray CCD camera. The material
Figure 5.6: Schematic of FALCON laser mirror high-vacuum chamber.

x-ray absorption effect is put in the x-ray count analysis to correctly obtain the total x-ray counts.

The electron beam at the interaction zone could be also maneuvered in transverse directions with the two steering magnets installed around the final focus system. We could then use these steering magnets to steer the beam in order to produce the highest intensity image detected by the x-ray CCD camera.

In addition, the same alignment procedure can be done by steering the FALCON beam against the fixed position of the electron beam spot for the highest x-ray image. The laser beam after the IP expands as it propagates further upstream and is then dumped along the walls of the beamline.
Figure 5.7: Steering magnets shown upstream and around the PMQ final focus system.
5.2.1 Alignment Diagnostics

A successful Thomson scattering experiment depends on the minimization of arrival time-delay between interacting beams. In order to adjust the time difference, various sophisticated fast-time detection instruments have been implemented around the interaction zone. Their read-outs are then used in the system tuning procedure. In this section, we provide details on instruments and detection methods.

The first step in the timing synchronization is the determination of the initial time separation between the FALCON laser and the PLS/electrons. The initial time separation in this case is approximately equal to the path length separation, which is about 70 meters. In order to detect the arrival time of the two beams, two different types of detectors are required. For the electron beam, a pair of two 100 ohm loops in parallel is used to detect a short magnetic field pulse (150 ps accuracy) from a passing electron beam. Simultaneously, a fast infrared UHS 016 photodiode detector is set up to detect the arrival time of an IR laser with a similar accuracy. The time difference readout is then used to provide the information necessary for the selection of a correct oscillator pulse to switch into the FALCON regenerative amplifier, relative to the one switched into the PLS regenerative amplifier. After the correction, the arrival time difference is reduced down to 12 ns or the time separation of two consecutive oscillator pulses.

In the second stage, the time difference is reduced further down from 12 ns. In order to perform this task, an Imacon 500 streak camera is set up to detect beam signals reflected off the aluminum cube. This streak camera is a high-resolution imaging instrument, equipped with an S20 photo-cathode with a quantum efficiency greater than 5% over the visible wavelengths and a 50 mm f/1.4 lens. At slow readout rate, the arrival time difference is measured, which is
shown as two beam images located at two different heights. The adjustment is then achieved by translating the retroreflecting roof mirror in the laser compressor on a 2 meter long rail, effectively changing the laser path length. A change in mirror position has no significant change in the characteristics of the compressed pulse, because of its location between the second and third grating strike in the compressor.

In the final stage of the timing synchronization, a streak camera is operated at the highest sweep speed of 18.7 ps/mm and performs the same procedure with the roof mirror. In this way, the time separation between the FALCON laser and the electron beam has been reduced to that of the streak camera resolution of 2 ps. An extra optimization step has been performed while observing the intensity of x-rays, but no improvement is observed.

5.2.2 X-ray Diagnostics

At an early stage of the PLEIADES experiment, an x-ray diagnostic instrument implemented in the experiment was a Princeton Instruments PI-SCX CCD. In the core of this camera is a 16 bit, 1340×1340 pixel charge-couple device semi-
conductor whose total detection area is 2.54×2.54 cm. A separate 3:1 taper fiber optic bundle was coupled to a 145 µm thick light producing CsI(T1) scintillator located on the front face of the camera, producing a field of view as large as 7×7 cm. Also, the scintillator was protected by a 0.5 mm beryllium filter to keep stray light out. The CCD was calibrated with a 60 keV $^{241}$Am source, where a sensitivity measurement of 7.4 counts/photon or 0.12 counts/keV at 60 keV was obtained.

A different type of x-ray CCD camera, whose working principle is based on the fast gated intensifier, has now been implemented at the facility. The use of this instrument has shown an improvement in detection efficiency from the previous camera primarily by cutting down the high electronic noise background. This gated CCD camera is a Princeton Instruments PI-MAX-512 CCD system tested to have a signal-to-noise ratios of 10:1, which is a huge improvement over the 1.5:1 ratios of the PI-SCX camera. A long exposure is also possible by activating the electronic gating function in the CCD camera, which then creates a window of opportunity for x-ray bursts to be read while shutting out other background noise. This type of instrument is most useful for the measurement of sensitive x-ray diffraction signals.

5.3 Thomson X-ray Productions

X-rays have been produced via the Thomson scattering process in a 180 degree collision geometry at PLEIADES. From the data that has been obtained emerge three very important questions to answer. First, what is the total x-ray flux produced? Second, what are the energies of the x-rays and are they in a good agreement with the Thomson x-ray energy equation? Third, what is the x-ray brightness achieved from this experiment? In order to answer these questions,
there are a number of calibration factors that we have to consider in the data analysis. For example, the transmittivity reduction due to the BK7 mirror in the x-ray beam path to the CCD detector and detection efficiencies. Details of data analyses and results are discussed in following sections.

5.3.1 Flux Measurement

The number of x-rays produced at the interaction point is obviously going to be different from how many actual x-rays are going to be detected by the x-ray CCD camera. How actual Thomson x-rays are detected is largely dependent on these known instrumental factors: x-ray transmission through the obstructing BK7 mirror, conversion of x-rays into visible photons by the CsI scintillator, and finally the sensitivity of the CCD. Therefore, in order to obtain the correct value of number of x-rays produced at the IP from a measurement, we need to find the normalization of the total number of x-rays predicted from 3D simulation code weighted by the theoretical fitting-curve. The theoretical fitting-curve used in the data analysis includes all real beam effects described in Chapter 1, and the x-ray transmission, CsI scintillation and detection sensitivity.

The transmission probability of a single x-ray making it through the turning mirror is dependent on what its given energy is. If the energy of a photon is not enough, this photon is likely to penetrate to a certain depth before being completely stopped by the material. A general theory of x-ray transmission is covered in the next section. The transmission curve is calculated based on the knowledge of material compositions of BK7. The BK7 mirror material is composed of 67% SiO$_2$, 12.6% B$_2$O$_3$, 8.1% Na$_2$O, and 12.3% K$_2$O by weight, and the total density of 2.51g/cm$^3$. The transmission curve as a function of x-ray energies of BK7 is shown in Figure 5.9. We note that in the spectrum below 30
Figure 5.9: (left) Transmission curve for BK7 and (right) CsI spectral response curve.

keV, x-rays are nearly all absorbed by the 1.6 cm thick BK7 mirror.

In addition, the spectral response of the CsI scintillator and the CCD sensitivity will be included in the total photon count equation. The spectral response of the CsI(T1) scintillator over a broad x-ray energy, 0-120 keV, is shown in Figure 5.9, along with $\eta_{CCD}(60\text{keV}) = 0.12$ count/keV from the x-ray CCD camera calibration with the $^{241}\text{Am}$ radioactive source and the GeLi-crystal detector.

The total counts that would be theoretically expected to be detected by the x-ray CCD camera are given by

$$N_{CCD} = \frac{\eta_{CCD}(\omega_c)}{P_{CsI}(\omega_c)} \int \int \int \hbar \omega_s T_{BK7}(\omega_s) P_{CsI}(\omega_s) \frac{dN_x}{d\Omega d\omega_s dt} d\Omega d\omega_s dt,$$

where for $\omega_c = 60\text{keV}$, $P_{CsI}(\omega_c) \approx 0.4$. Also, $T_{BK7}(\omega_s)$ and $P_{CsI}(\omega_s)$ are BK7 transmission and CsI spectral response respectively shown in Figure 5.9. The integrand term $\frac{dN_x}{d\Omega d\omega_s dt}$ in Eq. (5.10) is the intensity distribution resulting from a scattering of between $N$ total electrons with an emittance effect and a linearly polarized laser with a laser beam focusing and a laser bandwidth effects. This
expression is then used to fit with the actual x-ray measurement as shown in Figure 5.10. The adjusted total x-rays observed is obtained by weighting \( N_{\text{meas}} \) (the total area under the data curve) by the normalization factor. The normalization value is the ratio of \( N_{\text{sim}} \), total photons predicted by 3D simulation to \( N_{\text{CCD}} \), total counts representing the total area under the fitting-curve.

\[
N_x = \frac{N_{\text{sim}}}{N_{\text{CCD}}} N_{\text{meas}}. \tag{5.11}
\]

In Figure 5.10, an x-ray profile shown on the left hand-side is obtained at an electron beam energy of 57 MeV, which translates to peak 75 keV (on-axis) x-ray energy. The profile obtained also agrees well with the 3D simulation profile. The total counts measured are obtained from the data as a total area under the curve, and, from the data shown, about \( 4.5 \times 10^6 \) counts are measured. The normalization value obtained is 0.98 photons/count. By multiplying these two numbers, we then finally obtain a figure of \( N_x = 4.4 \times 10^6 \) photons produced from the linear Thomson scattering in a 180 degree geometry.

### 5.3.2 K-edge Radiographic Technique in X-ray Spectrum Measurement

We considered several methods in measuring the x-ray spectrum, which is crucial in quantifying the x-ray energy directly in proportional to the electron beam energy. We had three measurement methods that we considered: Ge(Li) detector, x-ray bent-crystal spectrometer [85], and K-edge radiography [86, 87]. The first method, the Ge(Li) detector, directly measures the amount of energy absorbed in the semiconductor during an x-ray imaging integration. The problem with this method is that a numerous noise is also registered along with an actual x-ray signal. Because of this unwanted background noise, it makes it very difficult to separate an analyzable x-ray image from a tarnished data image. On the other
Figure 5.10: (left) Actual x-ray profile detected with x-ray CCD camera and (right) a line (photon counts vs. angle) curve generated from the x-ray profile.

hand, an x-ray bent-crystal spectrometer diffracts only genuine x-ray signals onto the detector. A background noise coming from upstream is filtered out by a filtering slit plate made of a high-z metal such as titanium. The problem with this type of instrument is that the efficiency of the Bragg diffraction is very small. Therefore, narrow bandwidth x-rays are satisfied by the small window of Bragg angles available for photons to actually filter through a narrow slit. Even after a long integration time, a weak x-ray image is tarnished by radiation activated accelerator parts from other directions.

An ingenious technique has been devised by the PLEIADES team, which employs the use of K-edge radiography. The idea is that every metal made out of a pure element has a well-known characteristic K-edge line at a K-edge energy. Channeling a Thomson x-ray energy immediately below and above this K-edge produces different signature x-ray images. When an x-ray energy is tuned just below the K-edge it is highly transmitted through the metal, thus producing a full image. On the other hand, if it is just tuned above the characteristic energy,
x-rays are strongly absorbed producing a sort of dark hollow x-ray image (x-ray deficient) on-axis. Therefore, by knowing where the location of the K-edge energy of a material is, and channeling Thomson x-rays by tuning an electron beam energy, we can extract the information about the x-ray energy.

First, we describe the basics of how this K-edge effect works. The absorption of an x-ray is quantified by an absorption coefficient as a function of x-ray energies $\mu(\omega_s)$ and the thickness of the material. The attenuated intensity of x-rays in the material of thickness $dz$ can be given by $-dz = I(z, \omega_s) \mu(\omega_s)dz$. The solution $I(z, \omega_s)$ as a function of material depth is an exponential decay of $I(z, \omega_s) = I_0 e^{-\mu(\omega_s)z}$. But at a K-edge energy ($\omega_K$), an absorption cross section ($\mu(\omega_K)$) sharply increases almost vertically when the energy of an incident photon is taken, and at the same time the K-shell electron is energized to expel into a continuum state from its binding orbital state. At this moment this absorbed photon is never retrieved.

The x-ray source generated via linear Thomson scattering is quasi monochro-
matic, i.e., there is a correlation between the x-ray energy and the x-ray emission angle. The characteristics of this x-ray wavelength dependence on the emission of angle is described by \( \omega_s = 2\omega_L\gamma^2(1 + \cos \theta) \). For the K-edge experiment, we have chosen the 0.005" thick tantalum plate, which has a distinctive K-edge line at 67.46 keV. Similarly, to implement the K-edge transmission, Thomson x-rays of energies near this K-edge line are produced. With the 3D simulation, the energy dependence on emission angles is calculated, with the result shown in Figure 5.11. In an actual experiment, a tantalum foil is placed in the path of the x-ray, where x-rays of different wavelengths distributed outwardly, and similar in characteristics to those shown in Figure 5.11, is transmitted through to the x-ray CCD camera. By the combination of Ta K-edge and the Thomson x-ray spectral-angle dependence, a relatively easy experiment is accomplished, as shown in Figure 5.11. The 2D transmitted x-ray profile detected then can be used to infer the peak x-ray energy and verify its relation to the electron beam
energy.

In the experiment, the energy of the electron beam is tuned to 54.9 MeV to generate Thomson x-ray energies of 72.8 keV on-axis peak value. At this energy, x-ray wavelength is channeled slightly above the Ta K-edge where the attenuation is strongest. When x-rays are transmitted through the foil, a region of dark area is created in the center, while moving further outward two transmission regions reappear around the perimeter. This confirms that the peak x-ray energy at the center is strongly attenuated, whereas longer wavelengths on the top and bottom perimeters of the x-ray profile transmit unaffected. In Figure 5.12, it shows where the location of the K-edge is (67.46 keV) along with a region of strong attenuation [88]. The same experiment run is again performed, but this time at a higher x-ray energy. The energy is tuned to 78.2 keV at an electron beam energy of 56.9 MeV. At this energy, the x-ray transmission recovers substantially from the strong attenuation just above the K-edge line. So the dark region observed previously is filled up with transmitted x-rays in the center region, as shown in Figure 5.13. Qualitatively, these experimental runs shows that the K-edge transmission indeed worked as it should. Quantitatively, experimental observations are confirmed by 3D simulations performed with identical experimental parameters, as shown in the lower row of Figure 5.13. From this simple technique, there is a direct correlation between the peak x-ray energy and the electron beam energy. Also, the spectral-angle dependence shown in 3D simulation is in good agreement with the experimental observation.

For the PLEIADES Thomson x-ray source to be competitive against other x-ray sources, it is absolutely essential to deliver a stable, high x-ray flux over a wide range of x-ray energies. This x-ray tunability is achieved by tuning an electron beam energy to different settings. The electron beam linac currently existing
Figure 5.13: (top) Actual transmission x-ray profile detected with x-ray CCD camera and (top) 3D simulation predictions.
Figure 5.14: X-ray energy tunability.

at Livermore is capable of producing beam energies over the 20-120 MeV range. Using the simple Thomson energy equation, the corresponding range of x-ray energies are 10-160 keV. Furthermore, the whole linac system can be upgraded so that a gamma ray source can be produced. We have shown that in a single day of experiments, x-rays from as low as 40 to as high as 140 keV were produced with high efficiencies. Stable x-ray fluxes were also produced in this energy range. The experimental data is shown in Figure 5.14.
5.3.3 The Peak-Brightness of Thomson PLEIADES Source

In the categorization of x-ray sources, the source brilliance is often quoted as an accurate gauge of the quality of the x-ray beams produced. The majority of research experiments that are performed at any x-ray source facility require a high brilliance to successfully obtain meaningful experimental data. What creates the high brilliance? First, the rate of x-ray photons produced per unit time has to be high. This parameter is normally related to an x-ray flux, which marks the number of x-ray photons produced per second. Second, the x-ray source area has to be small. The source area for the Thomson scattering process is determined by the degree to which an electron spot and a laser spot of two different cross-sectional areas overlap at the interaction point. The total area of the scattering region is measured in units of mm$^2$. Third, the degree of collimation of the x-ray beams has to be high. As discussed before, due to real beam effects scattering angles tend to vary. The final tendency towards angular deviation is determined by the x-ray beam’s divergence or expansion over some propagation distance. The x-ray beam’s tendency to diverge is measured in units of mrad$^2$ (considering both transverse directions). Finally, the spectral bandwidth, which indicates an x-ray source’s degree of monochromaticity, has to be high. This x-ray spectral bandwidth is given in terms of a relative energy bandwidth of 0.1%. The brilliance can be defined by these parameters as follows:

$$brilliance = \frac{flux[\text{photons/sec}]}{(source \text{ area}[\text{mm}^2])(source \text{ divergence}[\text{mrad}^2])(0.1\% \text{ bandwidth})}.$$  

Currently, a peak brilliance achieved at PLEIADES is approximately about $10^{16}$ photons/s/mm$^2$/mrad$^2$/0.1%bw from a 0.25nC, 10mm-mrad, 20µm electron beam scattering with a 810 nm, 54 fs, 500 mJ laser beam in 40-140 keV x-ray energy range. The experiment can be improved in many areas to allow further
increases to upper brilliance ranges. Some of these improvements can be expected from better electron beam emittances, higher optimizations in small beam spot-sizes, shorter electron beam durations, higher particle density environments via a higher repetition scattering rate, and a higher total number of scattering particles (higher total charges and higher total laser energy). These are some of the major technological areas that are currently being researched to extend the capabilities beyond current technological limits.

One area we have already achieved improvement in is in the electron beam final-focusing system. As mentioned earlier, we have achieved a small electron beam spot as small as 10-20µm with 10-20 mm-mrad rms normalized emittance. The ability to hard focus a beam to a small spot size naturally increases the total x-rays generated (Eq. (5.8)). Also, the Thomson source area contribution in the denominator of Eq. (5.12) computed from an integral of a convolution of two Gaussian beam areas is

\[ A_s \approx \frac{\sigma_l^2 \sigma_e^2}{\sigma_l^2 + \sigma_e^2}, \]  

which approaches 10^{-3} mm² values as an electron beam spot-size approaches that of a small laser spot-size. A combination of these improvements from achieving a smaller electron beam spot-size is expected to increase the source brilliance. However, the hard-focusing also introduces high deviations in scattering angles, in turn producing a broader source divergence angle. In several recent reports [89, 83], the degree to which the final electron beam spot-size affects the brilliance is reported to be insignificant, which puts more emphasis on the importance of the beam emittance. Central to this core argument is the assumption that the laser focusing effect has a negligible term, and therefore little impact on the brilliance. On the contrary, it can be shown that a laser focusing effect does play an important role in the overall source divergence term and consequently in the
source brilliance property. The end of this section will show the ways in which the complex intertwining relationships among beam parameters affect the properties of the source brilliance.

The present analysis will be concentrated on the characteristics of the source divergence generated from two initial beams. The source brilliance, which is inversely proportional to the source divergence, can be approximately examined for various beam spot-sizes and emittances. We assume that the beam spatial distribution is Gaussian for both an electron beam and a laser beam. Similarly, particle directions in a beam bunch are described by a Gaussian function. For electrons, this direction distribution is given by

$$f_e(\varphi) = \exp \left( -\frac{(\varphi - \pi)^2}{\sigma_e'^2} \right), \quad (5.14)$$

where I have included \(\pi\) to indicate that a 180 degree scattering is involved. Remember that the electron beam geometric emittance relation with other terms is \(\varepsilon_x = \sigma_x\sigma_e'.\) For a laser beam divergence distribution, this function is given by

$$f_l(\varphi) = \exp \left( -\frac{2\varphi^2}{\sigma_l'^2} \right), \quad (5.15)$$

where \(\sigma_l' = M^2\lambda/\pi\sigma_l.\) The free variable \(\varphi\) is defined to be the direction of the source x-ray. The source divergence is derived as usual from the integration of a convolution of two functions as written by

$$f_s(\sigma_e', \sigma_l') = \int_0^{2\pi} f_e(\varphi) f_l(\varphi) d\varphi, \quad (5.16)$$

where the final expression computed is given by

$$f_s(\sigma_e', \sigma_l') = \frac{\sqrt{\pi}}{2} \frac{\sigma_e' \sigma_l'}{\sqrt{2\sigma_e'^2 + \sigma_l'^2}} \left[ \text{erf} \left( \frac{4\sigma_e'^2}{2\sigma_e'^2 + \sigma_l'^2} \right) + \text{erf} \left( \frac{\pi \sigma_l'^2}{2\sigma_e'^2 + \sigma_l'^2} \right) \right]$$

$$\times \exp \left( -\frac{2\pi^2}{2\sigma_e'^2 + \sigma_l'^2} \right). \quad (5.17)$$
The total x-ray divergence angle as shown can be expressed in terms of beam spot-size terms given by $\sigma_e = \sigma'_e/\varepsilon_x$ and $\sigma_l = M^2\lambda/\pi\sigma'_l$. Using this expression, the behavior of total x-ray divergence angle for different laser spot sizes was studied. In Figure 5.15, x-ray divergence curves for a 30 MeV electron beam of varying beam spot-sizes with a normalized rms emittance of 10 mm-mrad colliding head-on with 2, 3, 5, 10, 15 and 30 $\mu$m laser beams are plotted. It is shown that down to 20 $\mu$m, x-ray divergence curves behave identically. With further compression of an electron beam below that limit, we see x-ray divergence curves start to spread apart for different laser spot-sizes. We can see that the electron beam size effect progressively becomes larger for a smaller laser spot size below 20 $\mu$m electron beam limit. We note that the 30 $\mu$m laser beam has the least x-ray angular divergence, whereas a strong x-ray divergence is noticed for when a 2 $\mu$m laser beam is colliding with the electron beam. Overall, we observe that for a laser beam down to 10 $\mu$m colliding with an electron beam below the limit, an x-ray divergence retains a finite range of $20 \leq \sigma'_s \leq 70$ mrad.

Amid beam divergence and spot-size effects, we have left out emittance’s role in a total x-ray divergence. As reported in [89, 83] at the limit where the laser angular spread is negligible such that the total x-ray divergence only depends on the electron beam angular spread, the x-ray brilliance scales inversely quadratically with the electron beam geometric emittance. Thus the final x-ray brilliance is independent of the electron beam focal spot-size, because this parameter is absorbed into the emittance term along with the electron beam angular spread term. Nonetheless, emittance’s effect on the total divergence has been investigated with both electron beam and laser beam angular spreads, and are shown in Eq. (5.17). From the calculations plotted in Figure 5.16, for various emittance values, we find that curves are well separated regardless of given electron beam spot-size ranges. Also, total x-ray divergence worsens as the emittance value
gets larger. This calculation shows that a low emittance has a huge impact on increasing an x-ray brilliance via minimization of a total x-ray divergence. In addition to the various curves plotted for different emittances, we have also plotted curves for two different laser beam spot-sizes. The peak of the x-ray divergence curve almost doubled when a laser beam size reduced from 20 to 10 µm. So, a hard focusing of a laser beam does have noticeable effect in creating a large x-ray divergence angle, consequently reducing the peak brilliance.

In this last section, the x-ray beam divergence effect on the peak brilliance has been analyzed. It shows that focal spot-sizes of both colliding beams are significant in relation to the final x-ray angular spread, and, therefore, the peak
Figure 5.16: The source divergence as a function of electron beam spot-sizes and the emittance effects.

brilliance. Also, it was found out that the electron beam emittance has a significant role in improving the peak brilliance via a high-quality, low emittance beam. However, complete evaluations, both theoretical and practical, and which study every possible aspect of real beams are needed to realistically verify the experimental results.
CHAPTER 6

Conclusion

Over the last couple years, we developed a novel, hard x-ray source based on the linear Thomson scattering process. This x-ray source combined the existing 100 MeV photoinjector electron beam linear accelerator and the cutting-edge high intensity CPA laser system to produce a high flux, quasi-monochromatic x-ray beam. High-quality electron beams produced from the low emittance photoinjector system allowed high x-ray productions, up to the system’s limit. For further increase in x-ray flux, we later developed and installed the highly reliable, sophisticated final focusing magnet lens based on the permanent-magnet quadrupoles. Each PMQ’s magnetic gradient field was measured to be an unprecedented value of 570T/m, which means that, in a triplet configuration, a focal beam spot-size as small as 10 µm could be produced over the 20-120 MeV beam energy range. Thus, in comparison to previously reported 50-60 µm beam focal spot focused with standard 15 T/m EM quadrupoles, this system enabled a focal spots to be almost 5 to 6 times smaller. Consequently, we observed higher x-ray flux than before.

With this method, the x-ray flux can also be increased via tight focusing of the incident laser at the interaction point, therefore providing more photons to interact with electrons. The laser intensity defined by $I_L [\text{W/cm}^2] = 2P_L / \pi w_L^2$ shows that the incident photon density can be increased through providing either a larger laser power, $P_L$, or a smaller laser spot size, $w_L$. From the intensity
expression given above, we see that a smaller laser produces a higher concentration of power in a small region so that the laser field strength defined by $a_L$ is increased. A normalized laser vector potential defined as a function of $I_L$ is given by $a_L = 0.85 \times 10^{-9} \lambda_L I_L^{1/2}$ which indicates that a non-linear condition $a_L \geq 1$ is manifested if a sufficiently high laser intensity is focused in a small region. In the non-linear Thomson regime, harmonic generations and the depression of the x-ray energy practically make further improvement of x-ray brightness unattainable.

At $a_L \geq 1$ a strong laser field is established, and the magnetic field component of the laser becomes an effective contribution to the electron’s longitudinal motion in addition to the transverse linear motions due to the electric field component. In the stronger laser field, the electron trajectory exhibits a “figure-8” motion which produces non-linear signatures: namely, a reduction in scattered x-ray energy as a function of laser potential $a_L$ and the generation of x-ray harmonics over a wide energy spectrum range. We conclude the chapter with discussions on the experimental scheme of the proposed non-linear Thomson experiment at the PLEIADES and the proposed non-linear Thomson spectrum measurement.

6.1 The Design of the Non-linear X-ray Thomson Experiment

The non-linear Thomson scattering experiment has been proposed at LLNL PLEIADES facility [1, 14]. The non-linear Thomson x-ray experimental parameters are listed as shown in Table 6.1.

The experimental chamber has been designed to contain the laser and electron optics necessary to produce both small laser and electron beam spots at the interaction points required for the production of this type of non-linear Thomson
Table 6.1: Non-linear Thomson experiment parameters

<table>
<thead>
<tr>
<th>Parameter</th>
<th>Value</th>
</tr>
</thead>
<tbody>
<tr>
<td>Laser wavelength</td>
<td>820 nm</td>
</tr>
<tr>
<td>Laser peak power</td>
<td>10 TW</td>
</tr>
<tr>
<td>Laser rms radius</td>
<td>9 µm</td>
</tr>
<tr>
<td>Laser periods</td>
<td>20</td>
</tr>
<tr>
<td>Electron beam energy</td>
<td>40 MeV</td>
</tr>
<tr>
<td>Electron rms radius</td>
<td>10 µm</td>
</tr>
<tr>
<td>Normalized emittance</td>
<td>5 mm mrad</td>
</tr>
<tr>
<td>Energy spread</td>
<td>2 %</td>
</tr>
<tr>
<td>Total electron charge</td>
<td>300 pC</td>
</tr>
</tbody>
</table>

x-rays (refer to Figure 6.1). The standard 60″ (f/30) focal length parabolic mirror used in the linear Thomson x-ray production is to be replaced with a 12″ (f/6) focal length parabolic mirror. The short focal length parabolic mirror is capable of producing a tighter laser beam, one whose rms radius is approximately 9 µm, which means the laser potential parameter $a_L$ would exceed one. Because the new parabolic mirror’s focal point would be located inside the experimental chamber, the electron final-focus system has to be moved close inside, and to a precise position. The counter-propagating 40 MeV electron beam to be collided with the high-intensity laser will be focused to a typical rms spot-size of about 10-20 µm using the novel final beam focusing system developed at UCLA. The electron beam energy is determined by the design of the permanent-magnet dipoles that would be used to dump the beam into the heavily shielded beam dump. In consequence, the low-energy beam operation of the current final focus system cannot properly focus the beam, requiring the focal length of the final focus
system to be reconfigured for a longer focal length ($\beta$-function). The permanent magnet quadrupole (PMQ) final focusing system, which is replaced and currently in operation at PLEIADES, employs the two strong quads with a 560 T/m field gradient and one weak quad with a reduced 260 T/m field gradient. The three individual PMQs are arranged in a triplet configuration in the final focusing system, which focuses the electron beam in both $x$ and $y$ transverse directions.

Two permanent magnet dipoles (PMD), composed of high magnetization neodymium iron boride material will be used for the deflection of electron beams into the beam dump, which will be located outside and away from the experimental chamber. In the initial design study, we found that a single PMD design scheme would require a large volume of material to deflect the beam in a single stage. Given the limited space inside, this would physically interfere with op-
tical components closely set up nearby and, moreover, be dangerously close to the high power IR laser. Instead we have chosen to go with a scheme where two smaller PMDs will be set up to deflect the beam in two separate stages. This design scheme fits inside the small experimental chamber and circumvents the interference with the functions of the other optical components. The combined PMDs produce a final electron beam deflection angle of 33 degrees through the 6" diameter port located on the right-hand side of the chamber wall and along into the beam dump pipe.

The tapered PMD is designed, as shown in Figure 6.3, to deflect the electron beam after the Thomson interaction to a 15 degree deflection angle in the direction of the second PMD. Since the tapered PMD is located a few centimeters after the IP, the magnet blocks need to be shaped exactly to allow the high-power, converging laser beam to propagate through safely avoiding a damage to the magnet material caused from an intense laser heat. The yoke is a single piece of iron material formed in a cylindrical round shape on both the outside and inside. On the inside surface, two flat surfaces on exact opposite sides are cut.
Figure 6.3: (left) Tapered permanent-magnet dipole & (right) Radia simulation showing magnet magnetization.

to mount the two oppositely polarized magnet pieces. In addition, two separate aluminum block pieces, cut to fit in with sloped edges of top and bottom magnets, are pushed in from the two opposite sides of the gap by aluminum screws which hold the magnets securely in their respective places. The whole assembly is then mounted on the aluminum pole, securely joined from the bottom, which then will be mounted on the optical table. The variable gap created by the installation of the two slanted magnetic blocks produces a varying longitudinal dipole profile in which the entering electron beam sees the strong dipole field and gradually becomes weak as it exits.

Additional kick in the beam deflection is provided by the second C-PMD where the deflection angle produced is 18 degrees, as indicated in the design shown in Figure 6.4. In contrast to the tapered shape PMD, the magnet design will be the standard C-shape. The yoke of the dipole is cut from iron in order to maximize the magnetic field saturation level. Similarly, we avoided using a
strong adhesive glue for attaching the two strong permanent magnets by using separable aluminum cases. Each magnet is placed securely in its right place by clamping a pair of aluminum cases on both sides of each one, and locking them in securely with non-magnetic screws. The design of this clamp cases ensures the magnets will not be knocked out of place. In addition, this design scheme solves the problem of compromising the high-vacuum state inside the chamber when an adhesive glue substance is used. The second PMD then finally guides the electron beam into the beam-dump pipe. The beam dump site is approximately 40 cm distance away from the experimental chamber, and housed in lead brick stacks to effectively shield the bremsstrahlung radiation produced from the intense relativistic electron beam, and to prevent it from interfering with the x-ray diagnostic instrument located just outside, in the direction of Thomson x-ray beam.
Figure 6.5: (left) The longitudinal magnetic field of C-shape permanent-magnet dipole & (right) Tapered permanent-magnet dipole.

After the complete assembly of each PMD, we measured the respective longitudinal magnetic fields. These measurements were performed with a sensitive magnetic Hall probe mounted on a high-precision stepper motor stage, controlled remotely by a Labview computer station. These measurements are shown in Figure 6.5, and compared with the Radia simulations. Each measurement agrees reasonably well with the expected simulations, with a minor disagreement shown in the C-shape PMD. This disagreement, as we understand, comes from the misalignment of the probe translation direction with respect to the PMD’s main axis so that a slightly broader $B_z$ profile is obtained. Another disagreement is shown in the tapered PMD where the amplitude appears slightly less, due to the computation limit. This minor problem can be adjusted by increasing the level of the magnetic simulation elements, which would require more computer memories, as discussed in Chapter 3.

A non-linear Thomson x-ray detection scheme employing the K-edge transmission measurements has been proposed. With the low $a_L < 1$ obtained by
Figure 6.6: Barium mass attenuation coefficient curve in 30-45 keV range. A clear presence of K-edge line at 37.6 keV.

under-compressing the laser pulse, the x-rays tuned just above the K-edge are expected to produce very low transmission through a filter material. As the laser pulse is fully compressed back, the $a_L$ would exceed one, and the x-ray energy is depressed below the K-edge line giving rise to an increase in x-ray transmissions. For the experimental beam energy of 40 MeV, the x-ray energy produced is 37.6 keV. The corresponding candidate for the K-edge filter material is Barium, whose K-edge absorption occurs at around 37.5-37.6 keV. The mass attenuation coefficient plotted as a function of x-ray energy is shown in Figure 6.6 [88]. Barium is a soft metallic material, which when pure is silvery white in color. Also, the substance easily oxidizes in air and dissolves in either water or alcohol. Thus, the material should be used in a high-vacuum environment [90]. We will employ a fast-gated x-ray CCD camera, as we described in Chapter 5, to obtain the
transmitted x-ray profile images.

6.2 The Current & the Future Experiment Prospects

The linear Thomson x-ray source facility, PLEIADES, has successfully been demonstrated as a high brightness, ultra-short and energy tunable x-ray source. It employs various novel technologies such as the photoinjector system for production of low-emittance, high-charge electron beams, the CPA Ti:Sapphire fs high-intensity laser system and the UCLA ultra-strong gradient PMQ final focus system for production of ultra-small 10-20 $\mu$m electron beams. The linear Thomson source produced at PLEIADES has enabled studies of various diffraction effects, and allowed the probing of various different heavy-metals.

The behavior of the non-linear Thomson effect in a head-on collision of a high-field laser ($a_L > 1$) and a high-density small electron beam still needs to be investigated. The theoretical study of this experiment has revealed frequency downshifted harmonic lines with the ponderomotive broadening effect [91, 92]. But the necessary hardware such as the new short-focal length laser parabolic mirror, PMQ final focus system, PMDs, and other optical items have been designed and fabricated. The new beamline design has been reconfigured for the non-linear Thomson experiment.

A similar experiment has also been planned at the UCLA Neptune Laboratory, with the difference that the collision geometry proposed is 90 degrees scattering angle, which employs 15 MeV electron beam and 500 GW CO$_2$ Mars laser producing soft x-rays [93].
References


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