



Title	Optimization of micro and nanoscale targets for ultra-high field generation and ion acceleration
Author(s)	Zosa, Allen Hermoso Myles
Citation	大阪大学, 2022, 博士論文
Version Type	VoR
URL	https://doi.org/10.18910/91767
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DOCTORAL DISSERTATION

Optimization of micro and
nanoscale targets for ultra-high
field generation and ion
acceleration

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August 16, 2022

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Acknowledgements

I extend my sincerest and warmest thanks to my supervisor, Prof. Dr. Masakatsu Murakami, for giving me this golden opportunity and teaching me so much throughout my years in Japan. His knowledge and advice helped me grow rapidly.

I am also very grateful to my panel, Prof. Dr. Ryosuke Kodama and Prof. Dr. Yasuhiro Kuramitsu for their valuable input and dedication in reading my manuscript.

I would also like to thank Didar Shokov. Our academic discussions helped me refine my ideas.

I am also thankful to Dr. Yanjun Gu. Our conversations broadened my perspective of research and helped me expand my field of interest.

I would also like to acknowledge all the past and present members of the TLP group. Everyone helped me when I had trouble during my stay in Japan.

I give my special thanks to Dr. Melvin Empizo, who has helped me through some of my difficult times in Japan.

I'm also grateful to the Filipino community in Osaka for always giving me a warm welcome when I needed it.

Finally, I would like to specially acknowledge my fiance, Dr. Maricar Rosete, for supporting me through all my troubles and hardships for the past 13.5 years.

Abstract

The emergence of ultra-intense lasers enables new technologies and applications for relativistic laser-plasma interactions. Ion acceleration, magnetic field generation, and high-energy-density physics benefit from these developments. For ion acceleration, generating quasi-monoenergetic ions is difficult for laser-based techniques such as target normal sheath acceleration. However, quasi-monoenergetic ions can be obtained through the Coulomb explosion of hollow nanospheres. When the nanospheres are correctly optimized, they can be more than twice as efficient as the previous designs. These make the Coulomb explosion of hollow nanospheres attractive for medical applications (cancer therapy). However, while hollow nanospheres exhibit Coulomb explosion, hollow microspheres will display both Coulomb explosion and microbubble implosion. In this case, a larger target can hold more electrons in its cavity, enough to initiate a rapid implosion process. When undergoing microbubble implosion, the inner wall of the microsphere will compress to a density that is 10^5 times higher than solid density. This high compression will

generate highly accelerated flashed ions and a very intense electric field during maximum compression. Thus, this is particularly interesting for applications involving vacuum polarization and pair creation. Finally, while nanoscale hollow targets demonstrate Coulomb explosion and microscale hollow targets demonstrate both microbubble implosion and Coulomb explosion, designing a target with nanoscale and microscale components can produce a different phenomenon. In this case, a novel paisley design composed of a nanoscale tip and a microscale body can generate a highly intense sub-mega tesla magnetic field. These magnetic fields are challenging to produce and are usually found in extreme astrophysical environments. Hence, this novel design helps study fundamental physical phenomena such as magnetic reconnection or other laboratory astrophysics applications.

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The study of laser plasmas is an exciting and diverse field with multiple applications. For example, laser plasmas are helpful for nuclear fusion, cancer therapy, particle acceleration, magnetic field generation, and fundamental quantum electrodynamics (QED) experiments [1–5]. Furthermore, with the modern developments in ultra-intense laser technology, the scope of laser plasmas has expanded, allowing the discovery of new phenomena [3, 6, 7]. Among the advances in ultraintense laser technology, chirped pulse amplification (CPA) has been one of the most impactful [8]. With CPA, it is possible to amplify the intensity of femtosecond laser pulses drastically. Also, the emergence of ultraintense-femtosecond lasers sparked interest in phenomena such as Coulomb explosion (CE), non-linear QED effects, and x-ray generation, as well as allowing the emergence of novel phenomena like microbubble implosion (MBI) and microtube implosion (MTI) [1–5].

1.1 Electron heating using ultraintense-femtosecond laser pulses

When an ultraintense-femtosecond laser pulse interacts with a solid density plasma, it heats the plasma, with the electrons receiving most of the transferred energy. The following ponderomotive scaling law can estimate the electron temperatures:

$$T_e = m_e c^2 \left(\sqrt{1 + \frac{I_L (\text{W cm}^{-2}) \lambda_L (\mu\text{m})}{1.37 \times 10^{18}}} - 1 \right), \quad (1.1)$$

where T_e is the electron temperature, m_e is the electron mass, c is the speed of light, I_L is the laser intensity, and λ_L is the laser wavelength [9]. When $I_L \gg 10^{18} \text{ W cm}^{-2}$ and $\lambda_L \sim 1 \mu\text{m}$, Eq. 1.1 will simplify to

$$T_e (\text{MeV}) = 44 \sqrt{I_L (10^{22} \text{ W cm}^{-2}) \lambda_L^2 (\mu\text{m})} \quad (1.2)$$

According to Eq. 1.2, the electron temperature increases with laser intensity. How the electron temperature affects the system depends on the target size in relation to the laser skin depth. A target that is much larger than the laser skin depth will keep most of its electrons, while a target comparable to the laser skin depth can lose almost all of its electrons. Typically, the electron density profile is modeled using the

Boltzmann relation given by,

$$n_e = n_{e0} \exp\left(\frac{e\phi}{kT_e}\right) \quad (1.3)$$

where n_e is the electron density, n_{e0} is the reference electron density, and k is the Boltzmann constant. However, the Boltzmann relation assumes an isothermal profile which may not be valid in non-quasineutral expanding plasma. According to Eq. 1.3, as $n_e \rightarrow 0$, $\phi \rightarrow -\infty$. Thus, the Boltzmann relation will not hold when considering a system with empty regions [10].

1.2 Electron distributions under non-isothermal conditions

Calculating the electron distribution of non-quasineutral expanding plasma requires a general equation not limited to isothermal conditions. When electrons are in hydrostatic equilibrium with a static ion sphere,

$$\nabla p_e = en_e \nabla \phi, \quad (1.4)$$

where p_e is the electron pressure, and ϕ is the potential. According to the polytropic relations, pressure and density obeys the following

relationship,

$$p_e = C n_e^\gamma, \quad (1.5)$$

where C is a constant. Applying the boundary conditions, $\phi_0 = 0$, $n_0 = n_{e0}$, $T_e = T_{e0}$ and using Eq. 1.5 in Eq. 1.4, the following equation is obtained assuming ideal gas behavior,

$$n_e = n_{e0} \left[1 + \frac{\gamma - 1}{\gamma} \frac{e\phi}{T_{e0}} \right]^{\frac{1}{\gamma-1}}, \quad (1.6)$$

where n_{e0} and T_{e0} are the electron density and temperature at $\phi = 0$, respectively. In contrast to Eq. 1.3, Eq. 1.6 gives a finite ϕ if $n_e = 0$.

Using Eq. 1.6 in the Poisson equation, the electrostatic potential obeys the following equation,

$$\nabla^2 \phi = 4\pi e \left(n_{e0} \left[1 + \frac{\gamma - 1}{\gamma} \frac{e\phi}{T_{e0}} \right]^{\frac{1}{\gamma-1}} - Z n_i \right), \quad (1.7)$$

where Z is the ionization state, and n_i is the ion density. Considering only one-dimensional geometries and using normalized units, Eq. 1.7 simplifies to following equation,

$$\frac{d^2 \tilde{\phi}}{d\xi^2} + \frac{\nu}{\xi} \frac{d\tilde{\phi}}{d\xi} = \frac{1}{\Lambda^2} \left(\left[1 + \frac{\gamma - 1}{\gamma} \tilde{\phi} \right]^{\frac{1}{\gamma-1}} - \tilde{n} \right), \quad (1.8)$$

where $\tilde{\phi} = e\phi/T_{e0}$, $\xi = r/R$, R is the normalizing length, $\Lambda = R/\lambda_D$, $\lambda_D = \sqrt{T_{e0}/4\pi e^2 n_{e0}}$ is the electron Debye length, $\tilde{n} = n_i/n_{e0}$, and ν

is a geometric factor equal to 0 for planar, 1 for cylindrical, and 2 for spherical symmetry.

1.3 Electron distributions in spherical symmetry

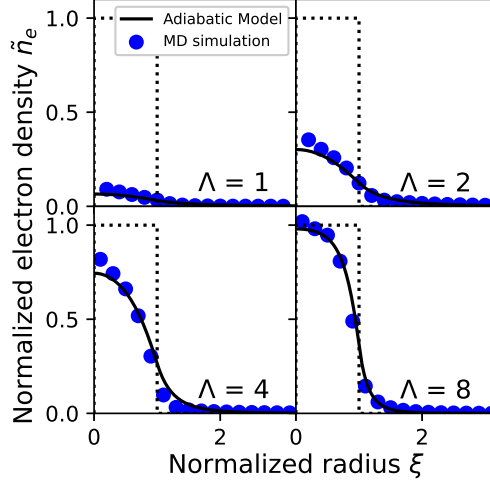
Spherical plasmas are a topic of interest because of their various applications in ion acceleration, nuclear fusion, quantum electrodynamics, and pair-creation [2, 4–6]. Thus, finding the electron distribution in spherically symmetric targets is a topic of interest because it is an essential aid in calculating the dynamics of spherical plasmas. However, current models require complex and time-consuming calculations or unrealistic assumptions about the system [10, 11]. Equation 1.8 provides a simple and accurate model to calculate the electron distribution around a spherical ion distribution for a given temperature. When considering an ion sphere of radius, $R = R_0$, Eq. 1.8 becomes:

$$\frac{\partial^2 \tilde{\phi}}{\partial \xi^2} + \frac{2}{\xi} \frac{\partial \tilde{\phi}}{\partial \xi} = \begin{cases} \Lambda^2 \left(1 + \frac{\gamma-1}{\gamma} \tilde{\phi}\right)^{1/(\gamma-1)} - \Lambda^2, & \xi \leq 1 \\ \Lambda^2 \left(1 + \frac{\gamma-1}{\gamma} \tilde{\phi}\right)^{1/(\gamma-1)}, & \xi > 1 \end{cases} \quad (1.9)$$

If there are no external sources of energy, the system is adiabatic, and γ will be the adiabatic index— 5/3 for classical electrons and 4/3 for relativistic electrons.

Figure 1.1 shows the electron distribution around a solid ion sphere at varying Λ with the normalizing length, R , set to be the ion sphere radius. One can verify the accuracy of Eq. 1.8 using three-dimensional particle-particle (3D-PP) simulations. The 3D-PP simulations used a uniformly charged cloud to model the ion sphere. Additionally, the simulation used 1000 computational particles for the electrons with an initial velocity based on the given Λ . Comparing the calculated electron density obtained using Eq. 1.8 and the results of the 3D-PP simulations, seen in Fig. 1.1, the two datasets agree well for different values of Λ . As Λ decreases, the electron Debye length to ion sphere radius ratio increases. Thus, when Λ decreases, the ion sphere has fewer bound electrons, decreasing the electron density profile. Conversely, as Λ increases, the electrons become more bound to the ion sphere, resulting in an electron density profile similar to the ion density.

The charge separation is also an essential factor when discussing laser plasmas because the degree of charge separation influences the plasma dynamics. For example, when the target has almost no electrons remaining, the dominant process is generally Coulomb explosion. Conversely, when most electrons are bound to the ions but create an electron sheath close to the surface, other phenomena such as target-normal-sheath acceleration (TNSA) and microtube implosion (MTI) can be observed. The charge separation is the ratio of the total charge of the electrons outside the ion sphere and the total ion sphere charge. Figure 1.2

Figure 1.1: Electron distribution at different Λ .

shows how the charge separation changes with Λ . With increasing Λ , the charge separation decreases because as the electron Debye length decreases, more electrons are bound to the ions. Comparing the Λ dependence of the charge separation obtained using Eq. 1.8 and the semi-analytic model obtained by Peano et al. [11], seen in Fig. 1.2, the two datasets agree well for different values of Λ . The good agreement further confirms the accuracy of Eq. 1.8.

1.4 Electron distribution of a hollow sphere

Hollow spherical plasmas have applications that solid spherical plasmas cannot replicate. For example, hollow microspheres can generate ultra-high electric fields and GeV ions through a process called microbubble

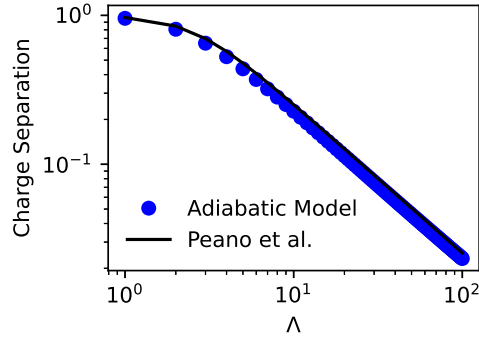


Figure 1.2: Charge separation of a solid ion sphere at various Λ .

implosion (MBI) [6, 7]. Additionally, it is easier to ionize a hollow target than a solid target of the same mass. Figure 1.3 shows the two-dimensional ionization profile of different hollow carbon targets with the same mass when irradiated with a $10^{21} \text{ W cm}^{-2}$ laser. As the ratio between the inner and outer radius (R_{in}/R_{out}) increases, the ionization profile becomes more uniform. When the target is completely solid ($R_{in}/R_{out} = 0$), the core of target is not completely ionized. However, the ionization profile becomes more uniform as the structure becomes thinner. A thinner target means it is closer to the laser skin depth, resulting in higher laser penetration. Thus, more of the material is ionized as the structure becomes more hollow.

Using Eq. 1.8, the corresponding equation for the electron density profile for a hollow sphere with inner radius, R_{in} , and outer radius, R_{out} ,

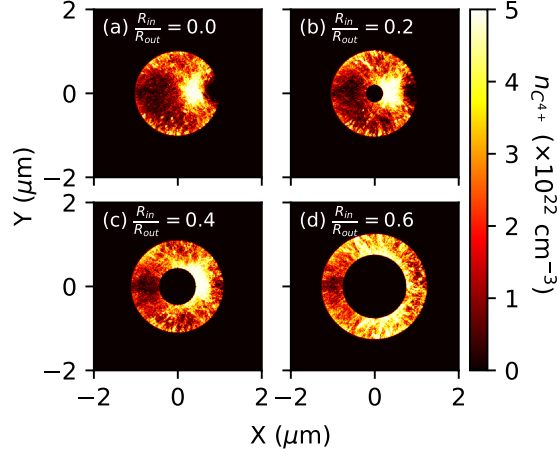


Figure 1.3: Ionization profile of different hollow targets.

is the following function:

$$\frac{\partial^2 \tilde{\phi}}{\partial \xi^2} + \frac{2}{\xi} \frac{\partial \tilde{\phi}}{\partial \xi} = \begin{cases} \Lambda^2 \left(1 + \frac{\gamma-1}{\gamma} \tilde{\phi}\right)^{1/(\gamma-1)} - \Lambda^2, & 1 \leq \xi \leq \xi_{ar} \\ \Lambda^2 \left(1 + \frac{\gamma-1}{\gamma} \tilde{\phi}\right)^{1/(\gamma-1)}, & otherwise \end{cases} \quad (1.10)$$

where $\xi = r/R_{in}$ and $\xi_{ar} = R_{out}/R_{in}$. Figure 1.4 shows the electron density profile solved using Eq. 1.6 and Eq. 1.10. The behavior is similar to Fig. 1.1, where the electron density decreases with increasing Λ . However, the electron density within the cavity does not follow a similar trend for the hollow target. Initially, the electron density in the cavity increases with decreasing Λ . Then, past a Λ value, the cavity electron density decreases with Λ , which implies an optimal value of Λ that maximizes the number of electrons in the cavity exists. Finding the optimal Λ value is crucial because many exciting phenomena such as

MBI and MTI rely on having electrons inside the cavity. MBI relies on as much electrons as possible to initiate a rapid implosion process. On the other hand, MTI relies on a lower electron temperatures where only an electron sheath is formed, initiating a slower isothermal expansion into the cavity. Additionally, there are two instances where almost no electrons are in the cavity. One instance is if the electron temperature is low enough that most electrons stay close to the ions. The other is if the electrons temperature is very high and almost all electrons can escape the ion sphere.

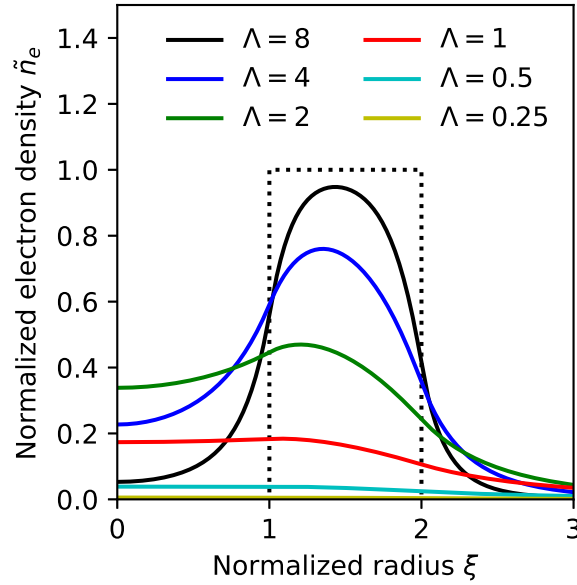


Figure 1.4: Electron density profile of a hollow sphere

The phenomena observed in hollow spheres depend on the number of electrons within the cavity and the value of Λ . The charge separation

is significant if there are almost no electrons in the cavity and $\Lambda \ll 1$. Thus, the target will undergo Coulomb explosion. If there are plenty of electrons in the cavity and $\Lambda \sim 1$, the cavity electrons produce a strong Coulomb force, which violently attracts the wall ions and initiates MBI. Lastly, suppose there are only a few cavity electrons in a cylindrical target and $\Lambda \sim 10$, which is large enough to avoid significant charge separation. In that case, the electrons produce a Coulomb force sufficient to attract the wall ions. If a seed magnetic field is present, this process initiates MTI, significantly amplifying the seed magnetic field. Figure 1.5 presents a rough diagram of the different regimes discussed. The following chapters will discuss the different regimes and their potential applications in detail.

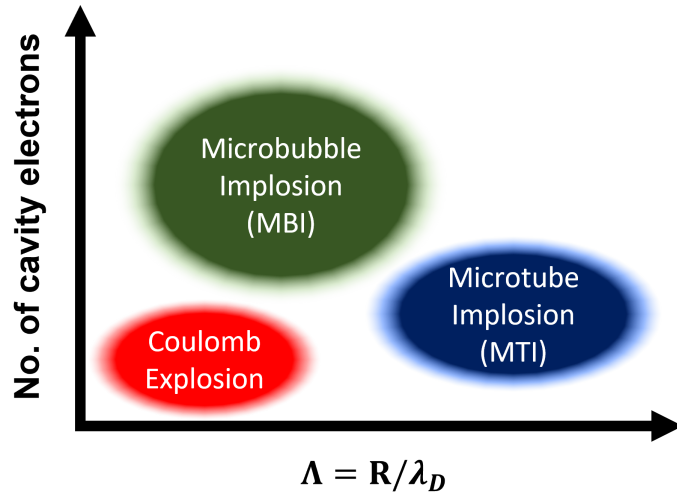


Figure 1.5: Different phenomena involving hollow targets.

Although Fig. 1.5 presents the distinction of each phenomenon, there are cases where overlaps occur. For example, a microsphere with an electron-filled cavity can undergo both MBI and Coulomb explosion. However, nanospheres will not undergo any significant MBI process. This is because the number of electrons contained in a nanosphere is substantially smaller. MBI is a process that relies on the total number of electrons inside the cavity, and this electron requirement does not necessarily scale proportionally to its size. In fact, a smaller target requires more electrons relative to its size to initiate MBI. Thus, a defining factor in MBI is the total number of electrons in its cavity. This will be discussed in more detail in the next chapter.

The topology and design of the targets are also crucial factors. For example, upon irradiation with an ultra-intense laser, nano-sized hollow targets are comparable to the relativistic skin depth of an ultra-intense laser. Thus, the laser can remove most of the electrons, and the Coulomb explosion is the dominant process. On the other hand, micron-sized hollow targets are much larger, resulting in more hot electrons filling the cavity, making them conducive for MBI. Additionally, the target symmetry, whether hollow sphere or cylinder, influences the mechanics of the system. For example, the electrostatic force under spherical symmetry has a $1/r^2$ scaling and a $1/r$ scaling under cylindrical symmetry. Therefore, phenomena requiring intense electrostatic interactions, such as MBI and Coulomb explosion, are more pronounced in hollow spheres.

However, hollow cylinders are more suitable for applications that prefer a specific direction such as magnetic field generation, because they are axially symmetric. Thus, the target structure, target size, electron Debye length, and the total number of electrons within the cavity are all essential factors determining the system dynamics.

2.1 Coulomb Explosion

The first regime to discuss is the Coulomb explosion (CE) regime. CE is an ion acceleration scheme that uses electrostatic repulsion to accelerate ions. In CE, an ultraintense laser heats the electrons, giving them enough energy to escape the ions. For spherical targets, the following condition needs to be satisfied for CE to occur [2]:

$$\frac{I_L}{1 \times 10^{18} \text{ W cm}^{-2}} \left(\frac{\lambda_L}{1 \mu\text{m}} \right)^2 \gg 4.0 \frac{\bar{Z} n_{i0}}{50 \text{ nm}^{-3}} \left(\frac{R_0}{50 \text{ nm}} \right)^2, \quad (2.1)$$

where I_L , λ_L , $\bar{Z}$, n_{i0} , and R_0 are the laser intensity, laser wavelength, average charge state, initial ion density, and initial radius, respectively. During CE, the charge separation process happens within $\sim 2R_0/c$, where c is the speed of light. Thus, the ejection of electrons occurs

within a few hundred attoseconds for ~ 100 nm spheres. Due to their large mass, the ions will have almost no interaction with the laser.

2.2 Energy spectrum for uniform sphere

Because the laser rapidly strips most electrons from the target, the expansion of an ion sphere is an approximation to the CE process [2, 12, 13]. In this case, the final kinetic energy of the accelerated ions is approximately its initial potential energy, i.e.:

$$\mathcal{E} = \frac{e^2 Z^2}{r} N(r) = \frac{4}{3} \pi e^2 Z^2 n_{i0} r^2, \quad (2.2)$$

where $N(r)$ is the total number of ions within a radius r . Thus, the energy spectrum as a function of energy can be expressed as,

$$\frac{dN}{d\mathcal{E}} = \frac{dN}{dr} \frac{dr}{d\mathcal{E}} = \frac{3}{4e^3 Z^3} \sqrt{\frac{\mathcal{E}}{\pi n_{i0}}}. \quad (2.3)$$

Thus, for spherical CE of a single ion species, $\frac{dN}{d\mathcal{E}} \propto \sqrt{\mathcal{E}}$.

2.3 Monoenergetic ions from two-species CE

While the energy spectrum of single ion species CE is better than standard target normal sheath acceleration (TNSA), it is still not monoenergetic. However, Murakami and Mima have demonstrated that two-species CE can produce monoenergetic ions in solid spheres [2]. Equation 2.2 shows why single ion species CE cannot generate monoenergetic ions [2]. Because $\mathcal{E} \propto r^2$, the outer ions have higher energies than inner ions. Thus, if an additional energy source increases the energies of the inner ions, the accelerated ions can be monoenergetic. This is achieved in two species CE.

The target is composed of light and heavy ions in two species CE. If the heavy ions are massive enough, they are approximately stationary relative to the timescale of the light ions. In this case, the light ion energy is the sum of contributions from the other light and heavy ions. For the light ion contribution, the total energy from the other light ions, $\mathcal{E}_l$, is the same as Eq. 2.2,

$$\mathcal{E}_l = \frac{e^2 Z_l^2}{r} N_l(r) = \frac{4}{3} \pi e^2 Z_l^2 n_l r^2. \quad (2.4)$$

Because the heavy ions are approximately stationary, the heavy ion

contribution to the light ion energy can be modeled as a uniform charged cloud. The electric field, E_h , due the heavy ion cloud is

$$E_h(r) = \frac{Q_h(r)}{r^2} = \begin{cases} \frac{4}{3}\pi e Z_h n_h r, & r < R_0 \\ \frac{4}{3}\pi e Z_h n_h \frac{R_0^3}{r^2}, & r > R_0 \end{cases} \quad (2.5)$$

Thus, the total energy from the heavy ions, $\mathcal{E}_h$, is

$$\mathcal{E}_h = \frac{4}{3}\pi n_h e^2 Z_l Z_h \int_r^{R_0} r dr + \int_{R_0}^{\infty} \frac{R_0^3}{r^2} dr = \frac{4}{3}\pi n_h e^2 Z_l Z_h \left(\frac{3}{2}R_0^2 - \frac{r^2}{2} \right). \quad (2.6)$$

Therefore, the total energy of the light ions, $\mathcal{E}$, is

$$\mathcal{E} = \mathcal{E}_l + \mathcal{E}_h = \frac{e Z_l}{R_0} \left(\frac{Q_l r^2}{R_0^2} - \frac{Q_h r^2}{2R_0^2} + \frac{3Q_h}{2} \right). \quad (2.7)$$

Equation 2.7 shows that if $Q_h = 2Q_l$, $\mathcal{E}$ is no longer dependent of r and will have the same value regardless of the initial position. Thus, the light ions will be accelerated monoenergetically if $Q_h = 2Q_l$.

2.4 Microbubble Implosions

The next regime is Microbubble implosion (MBI). MBI is a novel phenomenon observed in hollow structures with a substantial amount of cavity electrons [6]. A typical MBI setup is seen in Fig. 2.1(a). Consider a solid target with a micron-sized cavity and hot electrons [Fig. 2.1(b)].

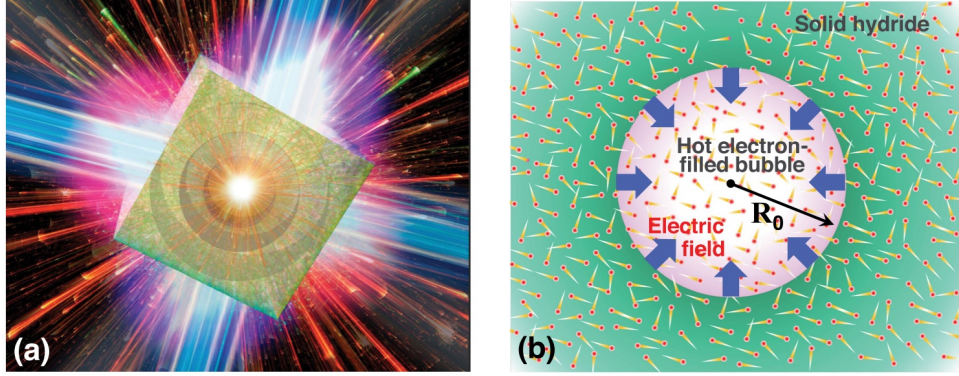


Figure 2.1: (a) Artistic representation of the MBI process, i.e. laser illumination, hot electron generation, proton implosion, and proton flashing. (b) A diagram of the mechanism involved in MBI. A cavity with hot electrons create a negative electric field, attracting the wall ions and initiating an implosion process [7].

The hot electrons can fill the hole with enough temperature, creating a negatively charged region that attracts the wall ions, initiating an implosion process. The implosion continues until the ions compress to a nanometer-sized radius and halt because of their intense Coulomb repulsion. This implosion results in a highly dense ion core, generating an electric field within a few orders of magnitude from the Schwinger limit [7].

2.5 Dynamics of MBI

To understand the dynamics of MBI, consider an electron-filled spherical cavity with electrons with a radius of R_0 and an average density of $\bar{n}_e$. Within the cavity, these electrons produce an electric field, E_{cavity} , given

by the following equation,

$$E_{cavity}(r) = -\frac{4}{3}\pi e\bar{n}_e r, \quad (2.8)$$

where e is the elementary charge. A test proton on the cavity wall will accelerate to a maximum implosion energy, $\mathcal{E}_{imp}$, of

$$\mathcal{E}_{imp} = \int_0^{R_0} eE(r)dr = \frac{2}{3}\pi e^2\bar{n}_e R_0^2 \quad (2.9)$$

$$= 150 \text{ MeV} \left(\frac{\bar{n}_e}{5 \times 10^{22} \text{ cm}^{-3}} \right) \left(\frac{R_0}{1 \text{ } \mu\text{m}} \right)^2. \quad (2.10)$$

Each proton on the innermost layer of the wall has a potential energy of $U(r_{ion}) = \frac{e^2 N_a}{2r_{ion}}$, where r_{ion} is the radius of the ion layer. N_a can be estimated by $N_a \simeq 4\pi R_0^2 d^{-2}$, where the interatomic spacing, d , is roughly $d \simeq n_{i0}^{-1/3}$ and n_{i0} is the initial ion density. By using energy conservation, i.e. $U(r_{min}) = \mathcal{E}_{imp}$, the minimum radius, r_{min} of the imploding wall ion can be expressed as

$$r_{min} = \frac{3n_{i0}^{2/3}}{\bar{n}_e}. \quad (2.11)$$

This result can be extended to different ion species and would similarly yield the following equation,

$$r_{min} = \frac{3Z^2 n_{i0}^{2/3}}{\bar{n}_e}, \quad (2.12)$$

where Z is the ionization state. Although Eq. 2.12 shows a Z^2 dependence, because $\bar{n}_e \propto Z$, $r_{min} \propto Z$. Thus, a carbon target ($Z = 6$) will have a minimum radius that is six times larger than a hydrogen target with the same n_{i0} . For a solid density proton target with $\bar{n}_e \approx n_{i0} = 5 \times 10^{22} \text{ cm}^{-3}$, $r_{min} = 0.81 \text{ nm}$ [6]. Remarkably, r_{min} does not depend on the initial cavity radius R_0 and only depends on n_{i0} , $\bar{n}_e$, and Z . Thus, a 100 nm cavity and a 1 μm cavity will implode to the same minimum radius if they have similar target compositions and densities.

Consequently, it is difficult to observe MBI especially in high- Z -hollow nanospheres. According to Fig. 1.4, a highly charged separated hollow sphere ($\Lambda < 1$) has cavity electron density that is 10% the initial electron density. Thus, when $\bar{n}_e \sim 0.1Zn_{i0}$, a fully ionized carbon sphere with $n_{i0} = 5 \times 10^{22} \text{ cm}^{-3}$ compresses to $r_{min} = 48.6 \text{ nm}$. The implosion is not prominent especially in targets with $\sim 10 - 100 \text{ nm}$ cavities. Thus, targets with cavities of the order 1 μm or greater is essential to observe MBI. Smaller targets will not have enough electrons in the cavity to violently attract the ions and initiate MBI.

Another way to understand this is by considering the net force on the innermost ion layer. Initially, the net force on an ion in the innermost layer is expressed as

$$F_{net}|_{t=0} = 2\pi e^2 \left(n_{i0}^{2/3} - \frac{\bar{n}_{e0}}{3} R_0 \right). \quad (2.13)$$

Interestingly, this equation shows that the repulsive term of the initial net force only depends on the ion density. It also shows that if the initial radius does not overcome a critical value, no implosion can occur. Unsurprisingly, the critical value is just r_{min} . It also shows that the radially-inward net force on the ion increases as R_0 increases. Thus, the acceleration of the ions gets more violent as the cavity radius increases because the repulsive term does not change with R_0 .

2.6 Ultrahigh electric field and ion flashing in MBI

When the ion layers compress to their minimum radius, they form a highly dense core with a density profile given by [6],

$$n(r) = \frac{n_{i0}}{6} \left(\frac{R_0}{r} \right)^2. \quad (2.14)$$

For a 1 μm solid hydrogen target, the minimum radius is 0.81 nm. Thus, according to Eq. 2.14, the maximum density achieved is $\sim 10^5$ times the initial ion density, n_{i0} . This ion density results in an electric field, E_f , given by the following equation

$$E_f(r) = \frac{eN_{e0}}{2R_0r_{min}} \left(\frac{r_{min}}{r} - \frac{r_{min}^2}{r^2} \right), \quad (2.15)$$

where $N_{e0} = \frac{4}{3}\pi\bar{n}_{e0}R_0^3$ is the initial number of electrons inside the cavity. Following Eq. 2.15, the maximum electric field occurs at $r = 2r_{min}$. Thus, the equation for the maximum electric field generated is

$$E_{f,max} = \frac{eN_{e0}}{8R_0r_{min}}. \quad (2.16)$$

For a solid density target with a 1 μm cavity, the maximum electric field is $\sim 10^{16} \text{ V m}^{-1}$, which is within 2 orders of magnitude from the Schwinger field, $1.32 \times 10^{18} \text{ V m}^{-1}$.

As the innermost ions halt and expand, the dense ion core formed during MBI further accelerates the innermost ions. The expanding inner core experiences the intense electric field given by Eq. 2.15. Thus, the innermost ions gain addition kinetic energy significantly more than their initial implosion energy. Murakami et al. demonstrated that flashed proton will be accelerated to GeV levels [6].

2.7 Microtube implosion

In the microtube implosion regime (MTI), implosion occurs but not to the same extent as MBI. In MTI, the implosion model is an isothermal hydrodynamic expansion into the cavity driven by an electron sheath surrounding the cavity walls rather than a result of the intense electrostatic attraction between the electron-filled cavity and the ions. In MTI,

$\Lambda \sim 10$, which is not enough to fill the cavity with electrons, but it is enough to induce an electron sheath close to the cavity walls. Figure 2.2 shows an overview of the MTI process. As the lasers heat the target electrons, they induce an electron sheath close to the cavity walls, which drives the isothermal implosion of the wall ions [3, 14]. A seed magnetic field bends the trajectory of the imploding particles, creating a Larmor hole. The collective electron and ion motion around the Larmor hole generates an amplified magnetic field in the same direction as the seed. Depending on the size and electron density of the target, the amplified magnetic field can reach the mega tesla level. The Larmor hole radius is known to scale as:

$$\frac{R_H(t)}{1 \text{ nm}} \simeq 43 \frac{Z}{A} \left(\frac{B_0}{1 \text{ kT}} \right) \left(\frac{R_0}{3 \mu\text{m}} \right)^2 \left(\frac{10^9 \text{ cm s}^{-1}}{c_s (1 + \xi_c(t))} \right), \quad (2.17)$$

where R_H is the Larmor hole radius, Z is the ionization state, A is the mass number, B_0 is the seed magnetic field strength, R_0 is the cavity radius, c_s is the ion speed of sound, and $\xi_c(t) = R_0/c_s t$. For the amplified magnetic field, B_c , its scaling is the following equation,

$$\frac{B_c}{1 \text{ MT}} = \left(\frac{Z}{6} \right)^{\frac{3}{2}} \left(\frac{A}{12} \right)^{-\frac{1}{2}} \left(\frac{n_{i0}}{10^{23} \text{ cm}^{-3}} \right) \left(\frac{R_0}{3 \mu\text{m}} \right) \sqrt{\frac{\mathcal{E}_{he}}{6 \text{ MeV}}}, \quad (2.18)$$

where n_{i0} is the initial ion density, and $\mathcal{E}_{he}$ is the hot electron temperature. Interestingly, this model predicts that the amplified magnetic field

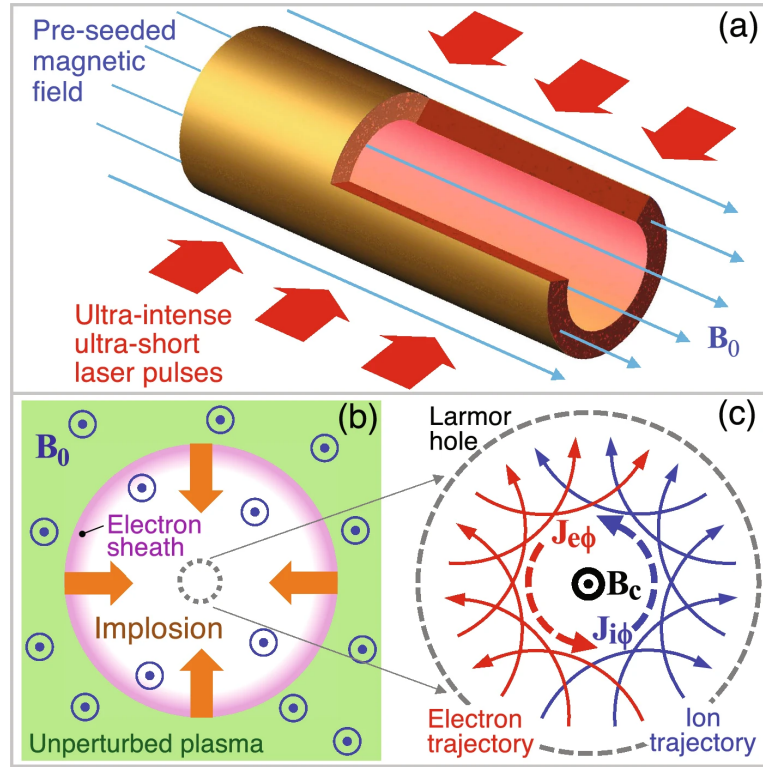


Figure 2.2: (a) Microtube irradiated by ultraintense laser pulses. A uniform seed magnetic field B_0 is present during irradiation (b) A diagram of the plasma dynamics. A hot electron sheath drives the isothermal ion implosion. (c) A Larmor hole and an amplified field, B_c , is formed at the center because of the collective electron and ion motions [3].

is independent of the seed magnetic field strength. For a $3\text{ }\mu\text{m}$ carbon target, to produce 1 MT, the hot electron temperature should be 6 MeV.

2.8 Polarity switching in MTI

The current MTI mechanism predicts that the induced magnetic field from the collective electron and ion motion has the same direction as the seed field. However, MTI can also generate an amplified magnetic field opposite the seed field. This phenomenon was reported and explained by Gu and Murakami. When the ions implode and form a highly charged ion core, their Coulomb field is strong enough to attract electrons, confining them in orbit around the ion core. As shown in Fig. 2.3, the electron orbit around the ion core is in the counter-clockwise direction, producing a clockwise current. The clockwise current induces a magnetic field along the $-z$ —*axis*— opposite the seed magnetic field. The anti-clockwise electron orbit is explained in Fig. 2.4. In the presence of a seed magnetic field along the $+x$ -axis, the Lorentz force rotates the electrons in the counter-clockwise direction. However, when an intense, attractive Coulomb field is present, the electrons sustain their anti-clockwise rotation but shift their trajectory to center around the Coulomb field. Thus, the orbital motion of the electron is consistent with their initial cyclotron motion.

The induced magnetic field opposite the seed field is primarily due

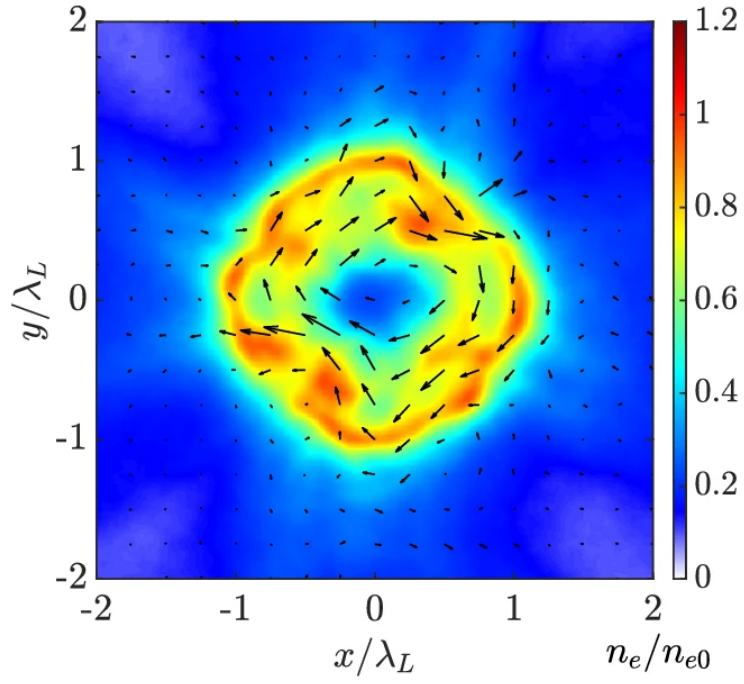


Figure 2.3: The electron density profile of MTI during its magnetic field generation phase with a seed magnetic field along the $+z$ -axis. The arrows represent the current vectors [15].

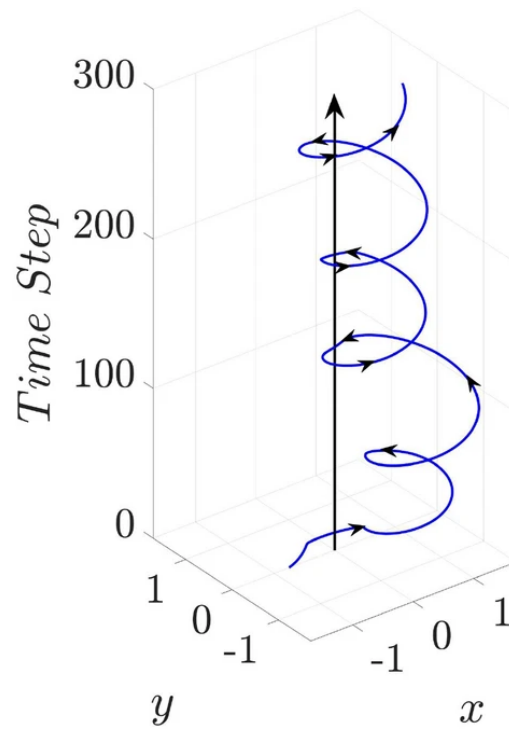


Figure 2.4: The trajectory of an electron in a seed magnetic field along the $+z$ -axis and a strong attractive coulomb field at the origin [15].

to the electron motion. Because the ion Coulomb field sustains the electron orbits, the induced magnetic field dissipates once the ion core expands.

2.9 MTI vs MBI

Although MTI and MBI share many common factors such as electron heating and ion implosion, the key difference between the two is that MTI can efficiently amplify a seed magnetic field, but MBI cannot. As explained in Fig. 2.2, the magnetic field is amplified because the electron and ion trajectories are bent, forming a Larmor hole inducing an amplified magnetic field in the same direction as the seed. Thus, it is clear that no amplified magnetic field is generated if the trajectories are not bent enough to form Larmor hole. In other words, to observe magnetic field amplification, the cyclotron radius, $r_{cyc} = \frac{\gamma m v}{e B}$, should be comparable or less than the cavity radius, R_0 . If $r_{cyc} \gg R_0$ then no amplified magnetic field is expected.

In MBI, the electron temperatures are much higher than MTI because MBI requires electrons to fill the cavity, not to create an electron sheath. Because the cavity is filled with more electrons, the ion implosion in MBI is much more violent with a much higher implosion velocity than MTI. Thus, for both the electrons and ions the cyclotron radius for MBI is significantly larger than in MTI. Additionally, it is unlikely for MTI

to occur in nanosized targets. According to Eq. 2.18, $B_c \propto R_0 \sqrt{\mathcal{E}_{he}}$, scaling down the radius by a factor of 10 means increasing the hot electron energy by a factor of 100 to maintain the same magnetic field strength.

Monoenergetic Ions via the Coulomb Explosion of Hollow Nanospheres

The discussion and figures below were lifted from M. A. H. Zosa and M. Murakami, Generation of quasimonoenergetic ions using optimized hollow nanospheres, *Physics of Plasmas*, 2020.

Ion acceleration has many applications in medicine, neutron generation, and fusion research [9, 16–22]. Target normal sheath acceleration (TNSA), radiation pressure acceleration (RPA), and Coulomb explosion (CE) are examples of laser-driven ion acceleration mechanisms [9, 18–20, 23–26]. However, medical applications require that the ions have a quasi-monoenergetic energy distribution [22].

3.1 Hollow Nanospheres

Section 1.4 discussed the advantages of using hollow targets. An important advantage is that hollow targets produce a more uniform ionization

profile. This is important because as discussed in Section 2.3, having an additional ion species can alter the acceleration profile of the light ions. Many methods to fabricate hollow nanospheres have already been reported [27–31]. Their most common application is in fuel storage research [28, 29]. However, we can also use these hollow nanospheres to enhance the ion acceleration profile of CE.

As discussed in Chapter 2.1, single species CE does not produce monoenergetic ions because the outer ions gain more energy. However, adding heavy ions results in monoenergetic light ions because the heavy ions contribute more energy to the inner ions, lowering the energy difference between the inner and outer light ions until no difference is observed. Thus, three different hollow nanosphere designs can be considered to give an enhanced ion acceleration profile (Fig. 3.1). For convenience, three target designs are named as type-I (Fig. 3.1a), type-II (Fig. 3.1b) and type-III (Fig. 3.1c). Each target type was designed to ensure that the heavy ions contribute more energy to the inner light ions. For this reason, a design with a heavy ion outer layer and a light ion inner layer was not considered. In this design, at best the heavy ions will contribute the same amount of energy to the light ions equally. In reality, due to the eventual expansion of the heavy ions, the heavy ions will contribute more energy to the outer light ions, increasing the energy difference between the inner and outer light ions.

If a hollow target and a solid target have equal mass, hollow targets

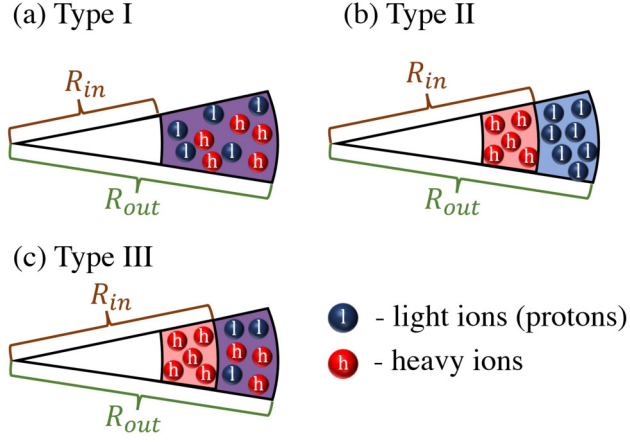


Figure 3.1: The different target types: (a) type-I, (b) type-II, and (c) type-III. A type-I target is a hollow nanosphere consisting of a proton-heavy-ion mixture. A type-II target is a layered target with an inner layer of heavy ions and an outer layer of protons. Lastly, type-III is a layered target with an inner layer of heavy ions and an outer layer of protons and heavy ions.

tend to be easier to ionize. It is one advantage that hollow targets have over solid targets. This is seen in Fig. 3.2, which is the ionization profile simulated using the two-dimensional (2D) particle-in-cell (PIC) code, EPOCH [32]. The simulation parameters are shown in Table 3.1. The solid target has a radius of $0.5 \mu\text{m}$, while the hollow targets were designed to ensure the total number of carbon particles is constant.

Figure 3.2 shows the advantages of using a hollow target. The solid target offers limited ionization coverage. Only the region close to the surface is completely ionized. Thus, the ionization profile is not uniform throughout the target. However, increasing the hollow radius also increases the uniformity of the ionization profile. This trend can

Table 3.1: EPOCH simulation parameters.

Simulation Parameters	
Laser intensity (I_L)	$2.2 \times 10^{21} \text{ W cm}^{-2}$
Laser wavelength (λ_L)	$1 \mu\text{m}$
Number of Grids	512×512
Grid size	$\lambda_L/128$
Particles per cell	30
Waist size	plane wave
FWHM	30 fs

be explained by considering the skin depth of the laser, given by the following equation

$$l = \frac{c}{\omega_{pe}} = \sqrt{\frac{\gamma m_e c^2}{4\pi e^2 n_e}}, \quad (3.1)$$

where c is the speed of light in vacuum, ω_{pe} is the electron plasma frequency, γ is the relativistic Lorentz factor, m_e is the electron mass, e is the elementary charge, and n_e is the electron density. γ can be approximated as $\gamma = \sqrt{1 + \frac{a_0^2}{2}}$, where a_0 is the normalized laser intensity. a_0 is given by the following equation

$$a_0 = \frac{eE}{m_e \omega_l c} \simeq 0.85 \sqrt{I_{18} \lambda_\mu}, \quad (3.2)$$

where E is the laser electric field amplitude, ω_l is the laser angular frequency, I_{18} is the laser intensity in $10^{18} \text{ W cm}^{-2}$, and λ_μ is the laser wavelength in microns. For a fully ionized carbon target with $n = 5 \times 10^{22} \text{ cm}^{-3}$ and a $1 \mu\text{m}$ laser with an intensity of $2.2 \times 10^{21} \text{ W cm}^{-2}$, the corresponding skin depth is $\sim 60 \text{ nm}$. The solid target in Fig. 3.2

has a radius of 500 nm, and the hollow shells have shell thicknesses of 408 nm, 327 nm, and 250 nm. As the shell thickness nears the skin depth, the ionization profile is more uniform because a larger fraction of the laser penetrates the target. Additionally, if the target size is comparable to the skin depth, the majority of the electrons in the target will be removed. If a laser can penetrate a target, it can transfer its energy to electrons more efficiently, resulting in higher electron temperatures.

This makes hollow targets more attractive because achieving a high ionization to laser energy ratio is easier. Additionally, for spherically symmetric CE, the uniformity of the ionization profile is critical to properly control the quality of the accelerated ions. As discussed in Chapter 2.1, adding a different ion species can alter the energy profile of the accelerated ions. Thus, limiting the number of unwanted ion species helps control the quality of accelerated ions.

3.2 Simple Analytic Models

In order to derive a simple analytic model for CE, some simplifying assumptions are made. Firstly, because most of the electrons escape during the CE process, to calculate the final kinetic energy of the ions, it is convenient to assume that the final kinetic energy is simply equal to the initial potential energy. For this to be true, it needs to be stipulated that the inner ions cannot overtake the outer ions.

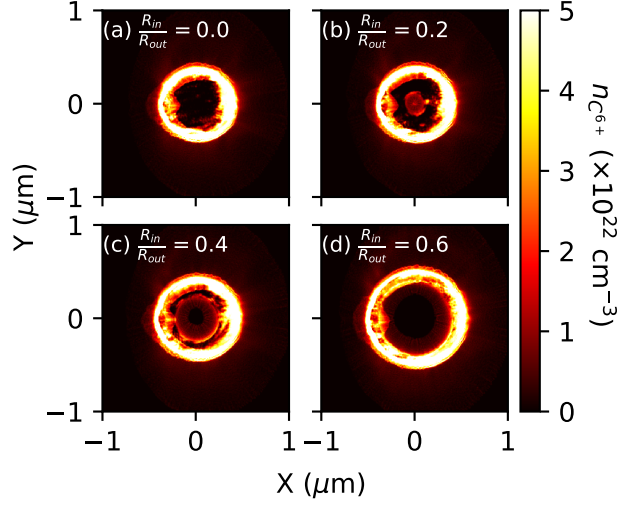


Figure 3.2: the C^{6+} density profile after laser irradiation, $n_{C^{6+}}$, for (a) solid target and hollow targets with (b) $R_{in}/R_{out} = 0.2$, (c) $R_{in}/R_{out} = 0.4$, and (d) $R_{in}/R_{out} = 0.6$.

For convenience, a parameter associated with the light and heavy ion species will have the subscripts l and h , respectively. If a parameter does not have any subscripts, they are implied to be the sum of the light and heavy ions contributions. The inner radius, R_{in} , and the outer radius, R_{out} , are shown in Fig. 3.1. It is also important to introduce the light ion charge fraction,

$$\beta = \frac{Q_l}{Q_l + Q_h} = \frac{Q_l}{Q}, \quad (3.3)$$

where $Q = Q_l + Q_h$ is the total charge of the system. Additionally, the

“cavity parameter”, γ , is defined as

$$\gamma = 1 - \xi_{\min}^3, \quad (3.4)$$

where $\xi_{\min} = \frac{R_{in}}{R_{out}}$. γ gives a rough idea of the structural properties of the target. A completely solid target will have $\gamma = 1$, while a very thin shell will have $\gamma \approx 0$.

The final kinetic energy of the light ions, $\mathcal{E}$, is the sum of the initial potential energy due to other light ions, $\mathcal{E}_l$, and the initial potential energy due to the heavy ions, $\mathcal{E}_h$, i.e. $\mathcal{E} = \mathcal{E}_l + \mathcal{E}_h$.

In order to generate quasi-monoenergetic ions using spherical targets, the inner ions should not overtake the outer ions. However, $\mathcal{E}_l$ for spherical targets with uniform densities increases away from the center, i.e., the ions closest to the surface have the largest $\mathcal{E}_l$. To counteract this behavior, $\mathcal{E}_h$ needs to decrease as the distance from the center increases, i.e., the surface ions have the smallest $\mathcal{E}_h$.

3.2.1 Type-I target

For type-I targets, the light ions initially located at r_0 will have $\mathcal{E}_l(r_0) = eZ_l \int_{r_0}^{\infty} Q_l(r_0)/r'^2 dr'$. $Q_l(r_0)$ is the total light ion charge enclosed by a Gauss sphere of radius r_0 , which is $Q_l(r_0) = Q_l(r_0^3 - R_{in}^3)(R_{out}^3 - R_{in}^3)^{-1}$ for a hollow sphere of uniform light ion density. Hence, $\mathcal{E}_l$ can be written

as

$$\begin{aligned}\mathcal{E}_l &= \int_{r_0}^{\infty} eZ_l \frac{Q_l}{r'^2} \frac{r_0^3 - R_{in}^3}{R_{out}^3 - R_{in}^3} dr' \\ &= \frac{eZ_l Q_l}{\gamma R_{out}} \left(\xi_0^2 + (\gamma - 1)\xi_0^{-1} \right),\end{aligned}\quad (3.5)$$

where r_0 is the initial radial position, $\xi_0 = r_0/R_{out}$, and Z_l is the light ion charge state. For $\mathcal{E}_h$, if the heavy ions are assumed to be practically immobile, $\mathcal{E}_h$ can easily be written as

$$\begin{aligned}\mathcal{E}_h &= \frac{eZ_l Q_h}{R_{out}} + eZ_l Q_h \int_{r_0}^{R_{out}} \frac{r'^3 - R_{in}^3}{R_{out}^3 - R_{in}^3} \frac{dr'}{r'} \\ &= \frac{eZ_l Q_h}{\gamma R_{out}} \frac{3 - \xi_0^2 + 2(\gamma - 1)\xi_0^{-1}}{2},\end{aligned}\quad (3.6)$$

Hence, the normalized light ion kinetic energy, $\tilde{\mathcal{E}}$, can be calculated in terms of composition or β , structure or γ , and ξ_0 , which is given by

$$\tilde{\mathcal{E}} = \frac{\mathcal{E}}{\mathcal{E}^*} = \frac{3(1 - \beta) + (3\beta - 1)\xi_0^2 + 2(\gamma - 1)\xi_0^{-1}}{2\gamma}, \quad (3.7)$$

where $\mathcal{E}^* = \frac{eZ_l Q}{R_{out}}$ is the surface light ion energy. Interestingly, when $\gamma = 1$ and $\beta = 1/3$, $\mathcal{E}$ is independent of ξ_0 . Thus, the energy spectrum is monoenergetic [2]. However, this is under the assumption that the heavy ions are immobile. In other words, this is under the condition that $\alpha = \frac{q_l/m_l}{q_h/m_h} \gg 1$. If the heavy ions and light ions expand at a similar rate (e.g., a target composed of fully ionized deuterium and

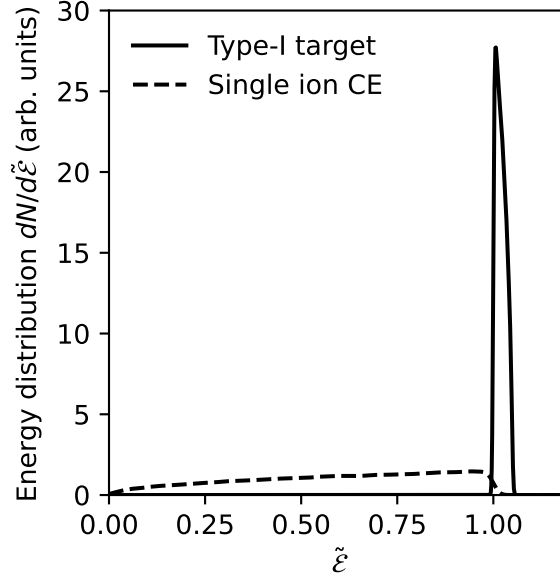


Figure 3.3: The energy spectrum calculating using Eq. 3.7 via Monte-Carlo simulation. Solid line for a type-I target with $\gamma = 1$ and $\beta = 0.3$. Dashed line for single ion species CE calculated using Eq. 3.5

oxygen), the energy spectrum will not be monoenergetic. The energy spectrum, in this case, is similar to single species Coulomb explosion, i.e., $dN/d\mathcal{E} \propto \sqrt{\mathcal{E}}$.

Figure 3.3 shows a sample energy spectrum of a type-I target with $\gamma = 1$ and $\beta = 0.3$. It shows a type-I target with ion energies that are significantly more monoenergetic than a single-ion species CE of a target with the same structure.

3.2.2 Type-II target

For type-II targets, Eq. 3.5 can still be used to calculate $\mathcal{E}_l$. In this case, the heavy-ion contribution, $\mathcal{E}_h$, can be approximated to be a point charge, i.e.,

$$\mathcal{E}_h = \frac{eZ_l Q_h}{R_{out}\xi_0}. \quad (3.8)$$

Therefore, $\tilde{\mathcal{E}}$ is written as the following equation:

$$\tilde{\mathcal{E}} = \frac{\beta\xi_0^2 + (\gamma - \beta)\xi_0^{-1}}{\gamma}. \quad (3.9)$$

However, ξ_0 cannot be eliminated for any specific combination of β and γ in Eq. 3.9.

In this case, one method to estimate the combination of β and γ that results in a quasi-monoenergetic energy distribution is by doing a series expansion. This series expansion yields,

$$\tilde{\mathcal{E}} = \left(1 - \frac{3\beta}{\gamma}\right) \Delta + 1 + \frac{\beta}{\gamma} + \mathcal{O}(\Delta^2), \quad (3.10)$$

where $\Delta = 1 - \xi_0$. However, Eq. 3.10 only holds when the ions are close to the surface ($\Delta \ll 1$). Thus, this is particularly useful when considering thin spherical shells. In this case, Eq. 3.10 is constant when $\beta\gamma^{-1} = 1/3$.

One feature of the Type-II target is that, while type-I targets require

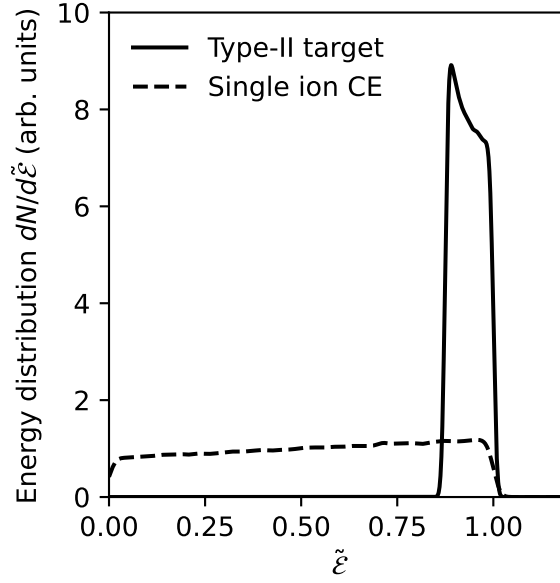


Figure 3.4: The energy spectrum calculating using Eq. 3.9 via Monte-Carlo simulation. Solid line for a type-II target with $\gamma = 0.488$ and $\beta = 0.3$. Dashed line for single ion species CE calculated using Eq. 3.5

static heavy ions ($\alpha \gg 1$), static heavy ions are not required for type-II targets as long as the heavy ions do not overtake the light ions.

Figure 3.4 shows a sample energy spectrum of a type-II target with $\gamma = 0.488$ and $\beta = 0.3$. Although, type-II targets do not have an exact monoenergetic condition, Fig. 3.4 shows that the type-II target ion energies are significantly more monoenergetic than a single-ion species CE of a target with the same structure.

3.2.3 Type-III target

Finally, the last target type we will discuss is the Type-III target. For this design, the total charge of the heavy ions is referred to as Q_h , the total charge of the heavy ions on the outer layer is $Q_{h_{out}}$, and the total charge of the heavy ions on the inner layer is $Q_{h_{in}}$. Here, $\mathcal{E}$ is written as $\mathcal{E} = \mathcal{E}_l + \mathcal{E}_{h_{out}} + \mathcal{E}_{h_{in}}$. Eq. 3.5 is still valid for the light-ion contribution, while the heavy-ion contributions are the following:

$$\mathcal{E}_{h_{out}} = \frac{eZ_l Q_{h_{out}}}{\gamma R_{out}} \frac{3 - \xi_0^2 + 2(\gamma - 1)\xi_0^{-1}}{2}, \quad (3.11)$$

$$\mathcal{E}_{h_{in}} = \frac{eZ_l Q_{h_{in}}}{R_{out} \xi_0}. \quad (3.12)$$

Thus, calculating $\tilde{\mathcal{E}}$ as a function of ξ_0 will give

$$\tilde{\mathcal{E}} = \frac{3(1 - \beta') + (3\beta' - 1)\xi_0^2 + 2(\gamma - 1 + \gamma\tilde{\beta})\xi_0^{-1}}{2\gamma(1 + \tilde{\beta})}, \quad (3.13)$$

where $\beta' = Q_l/(Q_l + Q_{h_{out}})$ and $\tilde{\beta} = Q_{h_{in}}/(Q_l + Q_{h_{out}})$. Additionally, β can also be written as $\beta = \beta'/(1 + \tilde{\beta})$.

Unlike Eq. 3.9, Eq. 3.13 has conditions to make it independent of ξ_0 . This happens when $\beta' = 1/3$ and $\gamma = (1 + \tilde{\beta})^{-1}$. However, the Eq. 3.7 condition requires γ to be dependent on $\tilde{\beta}$, while the value of β' should be $1/3$.

Comparing Fig. 3.5, Fig. 3.4, and Fig. 3.3 reveal that the type-I

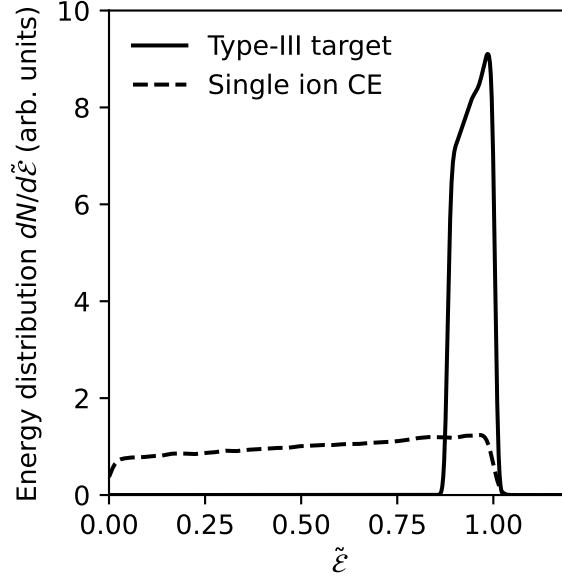


Figure 3.5: The energy spectrum calculating using Eq. 3.13 via Monte-Carlo simulation. Solid line for a type-III target with $\gamma = 0.595$, $\tilde{\beta} = 0.64$, and $\beta' = 0.4$. Dashed line for single ion species CE calculated using Eq. 3.5

target has the narrowest energy distribution because the type-I target has an exact monoenergetic condition with $\gamma = 1$ and $\beta = 1/3$ while the requirement for type-II targets is only approximate. In addition, the parameters used for the type-I target are closer to the ideal parameters, resulting in a very narrow energy spectrum.

3.3 Optimal Target Parameters

This section will discuss how to better estimate the optimal parameters. Previously, the optimal parameters for type-I targets were $\gamma = 1$ and $\beta = 1/3$. These conditions restrict both the cavity parameter and composition to a specific value. For type-II targets, the optimal parameters exist only for limitingly thin shells. Lastly, for type-III targets, the condition is $\beta' = 1/3$, restricting the outer layer's composition to a specific value.

$\tilde{\mathcal{E}}$ is normalized to ensure the outermost light ions have $\tilde{\mathcal{E}} = 1$. When an exact monoenergetic condition exists, $\tilde{\mathcal{E}} = 1$ for all light ions. Thus, a better optimization can be derived by minimizing distance between each ion energy and 1. This achieved by minimizing the total least-squared error term, $E = \int_{\xi_{\min}}^1 (\tilde{\mathcal{E}} - 1)^2 d\xi_0$. However, this equation does not account for the larger number of ions as ξ_0 increases. Because $\frac{dN}{d\xi_0} \propto \xi_0^2$, ξ_0^2 can be multiplied to the integrand to weight the squared-error term accordingly. Thus, the weighted-least-squared error is

$$\mathbb{E} = \int_{\xi_{\min}}^1 \xi_0^2 (\tilde{\mathcal{E}} - 1)^2 d\xi_0, \quad (3.14)$$

where Eqs. 3.7, 3.9, or 3.13 can be used for $\tilde{\mathcal{E}}$. Thus, minimizing Eq. 3.14 with respect to β will yield the values of beta that ensure that the ions have energies as close to the surface energy as possible.

Because the following equations will consist of long polynomial terms, it is convenient to write the equations in matrix form with the vector X defined as

$$X = \begin{bmatrix} 1 \\ \xi_{\min} \\ \xi_{\min}^2 \\ \xi_{\min}^3 \\ \xi_{\min}^4 \\ \xi_{\min}^5 \end{bmatrix}. \quad (3.15)$$

For the type-I targets, minimizing Eq. 3.14 with respect to β , i.e., setting $\frac{d\mathbb{E}}{d\beta} = 0$, yields

$$\beta = -\frac{X^\top B_1}{X^\top B_2} \quad (3.16)$$

where

$$B_1 = \begin{bmatrix} -16 \\ -48 \\ -96 \\ -41 \\ 117 \\ 84 \end{bmatrix}, B_2 = \begin{bmatrix} 48 \\ 144 \\ 288 \\ 270 \\ 90 \\ 0 \end{bmatrix}. \quad (3.17)$$

Interestingly, in the limiting case where $\xi_{\min} = 0$ ($\gamma = 1$), Eq. 3.16 will give $\beta = 1/3$, consistent with the previous section's result.

Repeating the minimization process for the type-II targets will yield

$$\beta = -\frac{28\xi_{\min}^5 + 49\xi_{\min}^4 + 63\xi_{\min}^3 - 28\xi_{\min}^2 - 49\xi_{\min} - 63}{20\xi_{\min}^4 + 60\xi_{\min}^3 + 120\xi_{\min}^2 + 130\xi_{\min} + 90}. \quad (3.18)$$

Additionally, Eq. 3.18 is consistent with Eq. 3.10. When considering the thin shell limit ($\xi_{\min} \rightarrow 1$), $\beta\gamma^{-1} \rightarrow 1/3$, which is shown in Fig. 3.6.

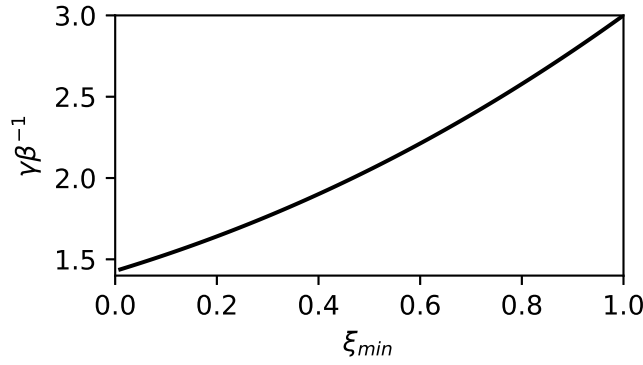


Figure 3.6: the graph of $\gamma\beta^{-1}$, calculated from Eq. 3.18, vs $\xi_{\min}$.

Interestingly, when $\xi_{\min} \rightarrow 1$, $\gamma \rightarrow 0$. This implies that $\beta \rightarrow 0$. Hence, Eq. 3.18 shows that thinner targets require a much higher heavy-to-light-ion ratio. Consequently, it can be inferred that this would degrade the overall energy efficiency because as more heavy ions are added to the system, a larger portion of the total energy is stored as heavy ion potential energy. Although a single atomic layer of hydrogen around a heavy ion sphere produces monoenergetic protons, the energy efficiency is low because most of the total energy belongs to the heavy ions.

Lastly, for the type-III targets, $\mathbb{E}$ is minimized with respect to $\tilde{\beta}$.

The result of the minimization yields

$$\tilde{\beta} = -\frac{X^\top B_1}{X^\top B_2} \quad (3.19)$$

where

$$B_1 = \begin{bmatrix} 48\beta'^2 + 17\beta' - 11 \\ 144\beta'^2 + 51\beta' - 33 \\ 288\beta'^2 - 108\beta' + 4 \\ 270\beta'^2 - 131\beta' + 107 \\ 90\beta'^2 + 87\beta' - 39 \\ 84\beta' - 28 \end{bmatrix}, B_2 = \begin{bmatrix} -49\beta' - 77 \\ 49 - 147\beta' \\ 28 - 84\beta' \\ 49\beta' + 77 \\ 147\beta' - 49 \\ 84\beta' - 28 \end{bmatrix}. \quad (3.20)$$

By setting $\beta' = 1/3$, Eq. 3.19 will reduce to $\gamma = (1 + \tilde{\beta})^{-1}$, the optimal conditions given in the previous section.

The least-squares approach gives a more robust estimate for the optimal composition or β for a given γ . Hence, these results are useful for target design purposes because the “ideal” parameters are not constrained to be specific values but are instead functions. This gives more freedom experimentally because an optimal target can be designed for a set composition or structure.

3.4 One-dimensional particle-particle simulation

To simulate the CE dynamics, it is computationally efficient to use a one-dimensional spherically symmetric particle-particle (PP) simulation. The following equation gives the equation of motion for each computational particle:

$$\ddot{r}_i = \frac{Q_i}{M_i} \left(\sum_{j \neq i, r_j < r_i}^N \frac{Q_j}{r_i^2} \right) + \frac{Q_i^2}{2M_i r_i^2}, \quad (3.21)$$

where Q_i , M_i , and r_i are the total charge, mass, and position of each psuedo-particle, respectively, while N is the total number of spherical particles in the system. The first term of Eq. 3.21 is the acceleration from inter-particle interactions, while the second term is the particle's acceleration from its own electric field. As shown in Fig. 3.7, the ions expand in roughly a spherically symmetric manner, which means that a spherically symmetric code can roughly approximate the process. This method was used by multiple research articles to study the Coulomb explosion dynamics of nanospheres [2, 12, 13]. However, this method is unable to reproduce theoretical results when simulating CE shocks caused by an initial density gradient [13]. D'Angola et al. observed that the spherically symmetric PP algorithm significantly softens the

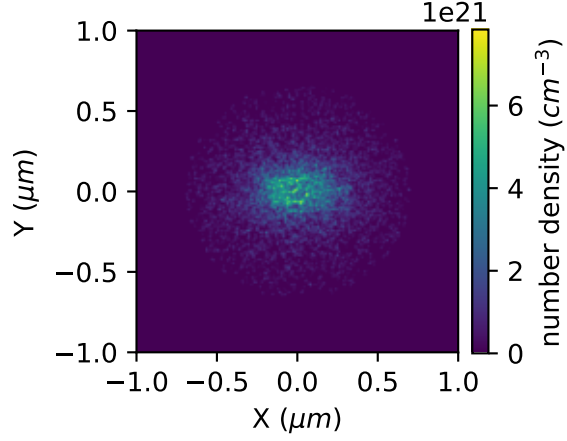


Figure 3.7: The carbon density from a three-dimensional PIC simulation of a fully ionized 50 nm solid carbon sphere irradiated by $1 \times 10^{21} \text{ W cm}^{-2}$ 1 μm laser with an FWHM of 50 fs. The snapshot is along the $z = 0$ plane at $t = 30$ fs.

density discontinuity associated with a CE shock front. However, the phase-space of the simulation particles roughly matches the theoretical predictions. Thus, for energy calculations, this method should suffice.

A total of $N = 1000$ computational particles were used. The light ions are protons, and the heavy ions are fully ionized carbon ions. For every target, radii were calculated to ensure $\mathcal{E}^* = 2 \text{ MeV}$. The initial density was $n_0 = 5 \times 10^{22} \text{ cm}^{-3}$. As shown in Fig. 3.8, only a small number of electrons remain close to the ion sphere after laser irradiation. In this case, the electrons are neglected because, in CE, the electrons

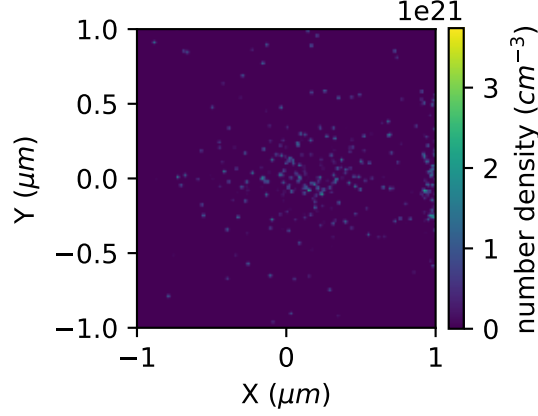


Figure 3.8: The electron density from a three-dimensional PIC simulation of a fully ionized 50 nm solid carbon sphere irradiated by $1 \times 10^{21} \text{ W cm}^{-2}$ $1 \mu\text{m}$ laser with an FWHM of 50 fs. The snapshot is along the $z = 0$ plane at $t = 30 \text{ fs}$.

have negligible contributions to the ion motion.

The simulation results show that the type-I targets have a low energy tail in their energy spectrum, as shown in Fig. 3.9. However, this low energy tail is absent in the type-II and type-III targets. Of the three designs, the type-II target has the lowest peak. This indicates that this target type is the most inefficient of the three designs. Additionally, only the type-II target has an energy distribution well approximated by the theoretical models. Each model assumes static heavy ions. This assumption does not work well when the light ions and heavy ions are

mixed in the same layer. Because the energy contribution of the heavy ions is calculated by integrating the Coulomb field of the ion cloud, an expanding ion cloud expands, decreases total energy transfer to the light ions decreases. On the other hand, type-II targets have the heavy ions on the inner layer. In this case, the assumption is valid as long as the heavy ions cannot overtake the light ions because according to Gauss' law $\oint \vec{E} \cdot d\vec{A} = 4\pi\rho$, a spherically symmetric charge distribution behaves as a point charge outside its surface. Thus, the expansion of carbon ions has no effect on type-II targets because the proton have a q/m ratio twice of carbon. The higher q/m ratio ensures that the carbon ions cannot overtake the protons. However, all three designs display a quasi-monoenergetic energy spectrum. Thus, quantifying the efficiency of each target type is very important to know which design yields a more efficient transfer of energy to monoenergetic light ions. Additionally, the carbon ions have roughly twice the maximum energy of the protons because the maximum energy is proportional to the charge state, Z , of the ions. The inset in Fig. 3.9(c) also shows a jump in the carbon energy spectrum at ~ 7 MeV because the Type-III targets have two different carbon density profiles. The energies above ~ 7 MeV corresponds to the outer carbon layer, and the rest correspond to the inner carbon layer.

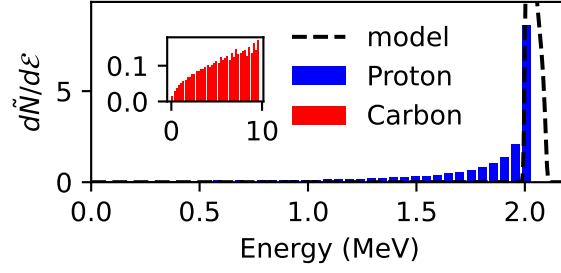
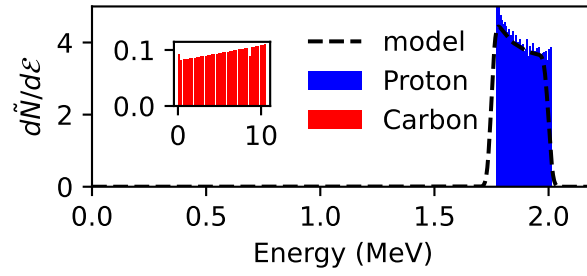
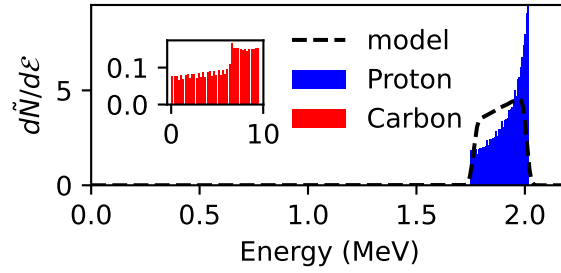
(a) Type-I energy spectrum for $\beta = 0.3$ and $\gamma = 1$ (b) Type-II energy spectrum $\beta = 0.3$ and $\gamma = 0.488$ (c) Type-III energy spectrum for $\beta' = 0.40$, $\tilde{\beta} = 0.64$, and $\gamma = 0.595$

Figure 3.9: the energy spectrum of the protons and carbon ions (inset) for the different target types. The spectrum is normalized such that $\int (d\tilde{N}/d\mathcal{E}) d\mathcal{E} = 1$. The dashed line is the theoretical energy distribution, the blue bars and red bars are the energy distributions of the protons and carbon ions, respectively.

3.5 Energy Coupling Efficiency

Due to the asymmetry of the energy distributions, to measure the monoenergeticity, a new parameter called the energy coupling efficiency, η , is defined as

$$\eta(\delta) = \frac{\int_{\mathcal{E}_{max}(1-\delta/2)}^{\mathcal{E}_{max}(1+\delta/2)} \mathcal{E} \frac{dN_l}{d\mathcal{E}} d\mathcal{E}}{\int_{-\infty}^{\infty} \mathcal{E} \frac{d(N_l + N_h)}{d\mathcal{E}} d\mathcal{E}}, \quad (3.22)$$

where $dN/d\mathcal{E}$ is the energy spectrum, δ is the fractional energy width, and $\mathcal{E}_{max}$ is the reference energy maximizing $\eta(\delta)$ which is generally close to the peak of the energy spectrum. The coupling efficiency parameter gives us the ratio of the total light ion energy corresponding to a $\pm\delta/2$ bandwidth with respect to the total energy of all the ions.

The maximum efficiency that can be achieved for a solid sphere with light and heavy ions can be approximated as

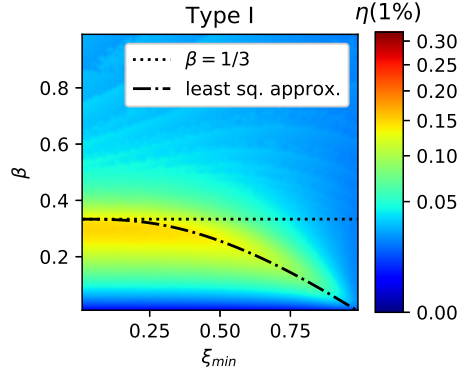
$$\eta_{max} = \frac{\frac{Q^2}{R} - \frac{Q_h^2}{R}}{\frac{Q^2}{R}} = 2\beta - \beta^2 \quad (3.23)$$

For a solid target with $\beta = 1/3$, the maximum efficiency is $\eta_{max} = 56\%$. This can be used to benchmark the performance of the target with respect to ideal conditions.

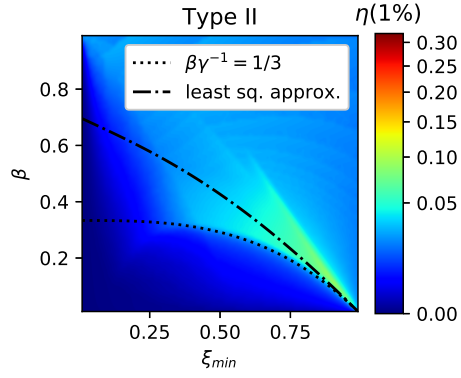
Figure 3.10 shows the simulation result. The type-I targets have high efficiency ($\eta(1\%) > 0.10$) for a large range of β and ξ_{min} values. This

implies that the target design is robust and can utilize many different structures and light-to-heavy-ion ratios. On the other hand, the type-II target has high coupling efficiency for thin targets ($\xi_{\min} > 0.75$). It has a narrow range of β and $\xi_{\min}$ values where it has reasonably high coupling efficiency. Thus, this target type may only be useful if the light ion layer is very thin. Finally, the type-III targets show the highest coupling efficiency values with a broad range of high efficiency based on structure and composition. Thus, type-III target are the most favorable because of its high efficiency. It is more efficient than the optimal type-I target, with the maximum efficiency being 15%, and 32% for type-I and type-III, respectively. Interestingly, the efficiency of the type-III target is slightly more than half the maximum efficiency of 56%. However, type-III target also has the highest deviation from the calculated optimal efficiency from the model because of the static heavy ion assumption. Because the type-III target has two heavy ion layers, it is more susceptible to errors caused by heavy ion expansion. Additionally, the no overtaking assumption in the models will have an effect especially in cases, where inner ions have more initial potential energy.

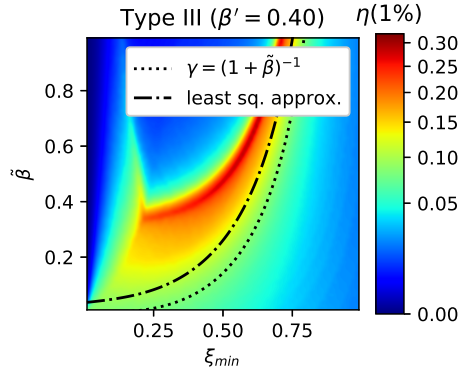
Finally, Fig. 3.10 shows that the regions of high efficiency for all three target types agree well with the optimal parameters predicted by the simple analytic and the least-squares minimization models. This confirms that the models are useful for estimating the optimal target structure and composition when designing the hollow nanospheres.



(a) Type I energy efficiency



(b) Type II energy efficiency



(c) Type III energy efficiency

Figure 3.10: the energy coupling efficiency of the different target types

3.6 Stability

To test the stability of the targets when perturbations are present in the spherical structure, three-dimensional particle-particle (3D-PP) simulations were utilized. In 3D-PP simulations, Coulomb collisions are not softened, which is a prominent problem in mesh-based algorithms such as particle-in-cell simulations. For this 3D-PP simulation, the binary Coulomb force on each particle was calculated using the following equations motion

$$\frac{d\vec{x}_i}{dt} = \vec{v}_i \quad (3.24)$$

$$\frac{d\vec{v}_i}{dt} = \sum_{j=1}^{N-1} \frac{q_i q_j}{|\vec{x}_i - \vec{x}_j|^3} (\vec{x}_i - \vec{x}_j), \quad (3.25)$$

where q is the particle charge, x is the position vector, and v is the velocity vector.

The perturbations were incorporated by using the following equation:

$$r = r_0 + r_0 \Delta \sin(n\pi x/r_0) \sin(n\pi y/r_0) \sin(n\pi z/r_0), \quad (3.26)$$

where r_0 is the unperturbed particle radius, Δ is the perturbation amplitude, and n is the mode frequency of the perturbations. Figure 3.11 shows a graphical representation of different perturbations applied to a sphere. A 50-nm solid target with $\beta = 1/3$ was simulated using 5000

particles.

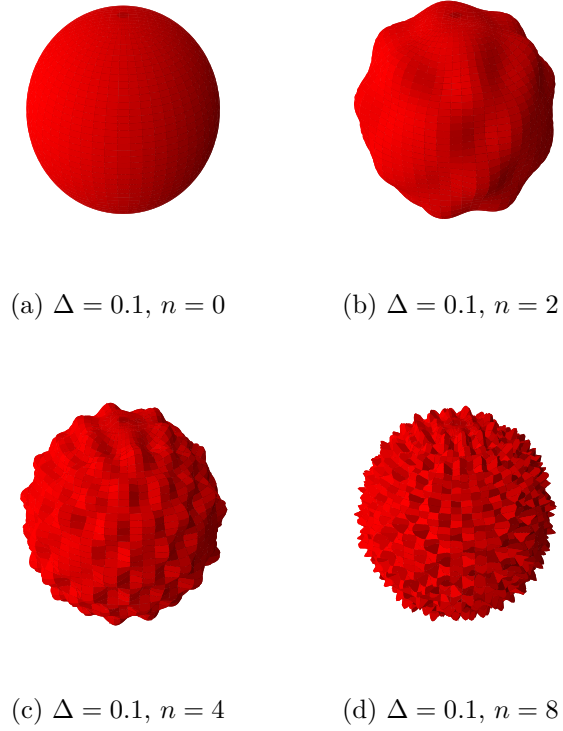


Figure 3.11: an illustration of the different types of perturbations used

Figure 3.12 shows no noticeable change in the energy spectrum even for $\Delta = 0.1$. Additionally, even when both Δ and n are varied, results show no appreciable difference in the efficiency with or without perturbations. This implies that the system is stable despite any asymmetry. This is supported by the experimental results where the CE of water droplets produced a quasi-monoenergetic energy spectrum [33].

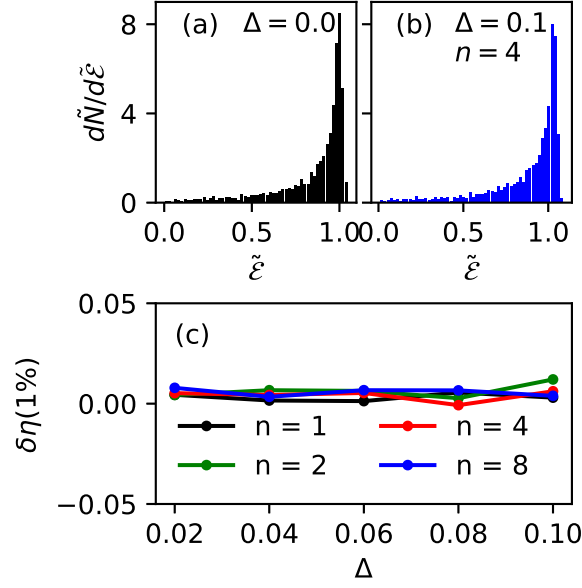


Figure 3.12: (a) and (b) show the energy spectrum for the unperturbed and perturbed spheres, respectively. (c) shows the plots of difference in coupling efficiency with respect to the unperturbed case, i.e. $\delta\eta(1\%) = \eta_{(\Delta=0)}(1\%) - \eta_{\Delta}(1\%)$, for different perturbation amplitude and mode numbers.

3.7 Section Remarks

This section shows that nanospheres are viable as laser targets to produce quasi-monoenergetic ions. The simple analytic and the least-squares models agree well with the simulation results. Thus, the optimal cavity parameter and composition can be calculated using these models. However, type-II targets can only have quasi-monoenergetic ions when the light ion layer is very thin. The simple analytic models are a general guide for the nanosphere design, but the least-squares models are more

robust and accurate in predicting the optimal parameters.

Lastly, nanospheres are also shown to be robust in the presence of perturbations. The new type-III target design has a maximum coupling efficiency of 32%, which is more than twice the type-I efficiency of 15%. Although the CE of hollow nanospheres is isotropic, the ions can be guided using a magnetic field or by adding the targets inside a carbon nanotube [22].

Microbubble Implosions in finite Microspheres

In the previous chapter, we discussed the applications of hollow spherical targets in the context of CE. However, the scale of the targets were too small to demonstrate the violent ion implosion in microbubble implosion (MBI). Using 50-100 nm targets will not show any significant implosion because of the following reasons. Firstly, as discussed in Section 2.5, the minimum radius of MBI is independent of target size and only depends on the initial ion density and the initial average electron density in the cavity. Thus, a 100 nm bubble and a 1 μm bubble will compress to the same minimum radius given the same densities. The latter might not be noticeable especially if the predicted minimum radius is a few tens of nanometers. Secondly, as discussed in Section. 3.1 small targets ($\sim 100\text{ nm}$ have comparable sizes to a ultraintense laser's skin depth. Because their size is roughly on the same scale as the laser's skin depth, it is easier for the laser to heat and remove the electrons from the

targets. This chapter considers $\sim 1 \mu\text{m}$ targets where there are enough cavity electrons to ensure a violent implosion of ions. Although MBI was previously studied by Murakami et al., that study assumes a quasi-neutral plasma [6]. Thus, no significant charge separation is observed. This chapter discusses the case where the plasma is charge separated and the outer ions are allowed to Coulomb explode while inner ions undergo MBI.

The discussion and figures below were lifted from M. A. H. Zosa and M. Murakami, Microbubble implosions in finite hollow spheres, *Physics of Plasmas*, 2022.

4.1 Microbubble Implosion

The emergence of relativistic laser facilities [34–37] ($> 10^{18} \text{ W/m}^2$) is a crucial factor in facilitating experiments involving various physical phenomena such as ion acceleration [9, 12, 18, 25, 38–42], relativistic electron acceleration [43–47], nuclear fusion [1, 48–51], or even quantum electrodynamics (QED) [52–58]. A more recent proposition, which takes advantage of relativistic laser intensities, is microbubble implosion (MBI) [6, 7]. MBI is a new and exciting proposition because it has various scientific and technological applications such as ion acceleration [7], gamma-ray lensing [4], and even pair creation [5].

In MBI, a target with a spherical cavity is irradiated by an intense

laser, which heats the electrons. Once the electrons are heated, they start to fill the cavity and initiate the implosion of ions close to the cavity walls. The rapid implosion in MBI is similar to another interesting phenomenon called sonoluminescence [59, 60]. However, in MBI, maximum compression is characterized by a highly dense but low-temperature proton core [6]. Conversely, the collapsed bubble in sonoluminescence has a high temperature but relatively low density.

During maximum compression, the number density, n , as a function of r close to the core of the implosion is given by

$$n(r) = \frac{n_{i0}}{6} \left(\frac{R_0}{r} \right)^2 \quad (4.1)$$

where n_{i0} is the initial ion density and R_0 is the initial cavity radius [6, 7]. In the case where $n_{i0} = 5 \times 10^{22} \text{ cm}^{-3}$, the minimum radius is roughly 0.8 nm [6]. This suggests that the ions are compressed to 10^5 times that of their original density. This high compression should yield electric fields one order of magnitude lower than the Schwinger limit [7]. Hence, vacuum polarization and even pair creation could occur in this scenario [4, 5]. However, it should be noted that Eq. 4.1 was derived for an infinite quasi-neutral plasma.

MBI has been investigated when the ion distribution is semi-infinite or periodic or when there are multiple laser beams to create a sufficient counterflow of electrons. Thus, most of the electrons would be confined

within the target, creating a quasi-neutral plasma.

To study the case where the electrons are free to redistribute themselves after they are given an initial temperature, a few factors have to be considered. In this case, the electrons can escape to infinity [11]. Thus, this gives the target a net charge. This makes calculating the electron distribution non-trivial. Although the electron distribution was successfully calculated in previous studies [10, 11], these studies either assume a background ion density or are computationally intensive.

4.2 Electron distribution model

Determining the equilibrium electron distribution requires solving the Poisson's equation, which is given as,

$$\nabla^2 \phi = 4\pi e (n_e - Zn_i), \quad (4.2)$$

where e is the elementary charge, Z is the ionization state of the ions, ϕ is the electrostatic potential, and n is the number density. The subscripts e and i denote electrons and ions, respectively. Typically, solving Eq. 4.3 involves using the isothermal condition. Thus, n_e should follow the Boltzmann relation, which is expressed as

$$n_e = n_{e0} \exp(e\phi/T) \quad (4.3)$$

where T is the temperature and n_{e0} is the density as $T \rightarrow \infty$. In spherically symmetric cases, Eq. 4.2 and Eq. 4.3 can be solved numerically for a static ion distribution. However, for finite-sized targets, the Boltzmann relation encounters a problem. As noted by Kanopathipillai et al., for finite-sized spherical targets, $n_e \rightarrow 0$ as $r \rightarrow \infty$, which implies that $\phi(\infty)$ diverges [10]. In this case, the electron expansion for finite-sized spherical plasmas cannot be assumed to be isothermal. A quick remedy is to add a small background ion density. This ensures that $\phi(\infty)$ does not diverge, and Eq. 4.3 holds [10].

Calculating n_e in a vacuum without background ions requires a more generalized form of Eq. 4.3. Assuming hydrostatic equilibrium, an expression for n_e can be obtained by utilizing the polytropic equation of state, $p_e = cn_e^\gamma$, where p is the pressure, γ is the polytropic index, and c is a constant [61, 62]. Thus, n_e can be written as

$$n_e = n_{e0} \left(1 + \frac{\gamma - 1}{\gamma} \frac{e\phi}{T_{e0}} \right)^{1/(\gamma-1)} \quad (4.4)$$

where the subscript 0 refers to an unperturbed homogeneous plasma (i.e., $n_{e0} = Zn_{i0}$). Unsurprisingly, when $\gamma \rightarrow 1$, Eq. 4.4 is equal to Eq. 4.3. Additionally, $\phi(\infty)$ no longer diverges under Eq. 4.4.

4.3 Electron distribution for a solid sphere

Utilizing Eq. 4.2 and Eq. 4.4, Poisson's equation for a solid sphere with radius R_0 is

$$\frac{\partial^2 \tilde{\phi}}{\partial \xi^2} + \frac{2}{\xi} \frac{\partial \tilde{\phi}}{\partial \xi} = \begin{cases} \Lambda^2 \left(1 + \frac{\gamma-1}{\gamma} \tilde{\phi}\right)^{1/(\gamma-1)} - \Lambda^2, & \xi \leq 1 \\ \Lambda^2 \left(1 + \frac{\gamma-1}{\gamma} \tilde{\phi}\right)^{1/(\gamma-1)}, & \xi > 1 \end{cases} \quad (4.5)$$

where $\xi = r/R_0$, $\tilde{\phi} = e\phi/T$, $\Lambda = R_0/\lambda_{De}$, and λ_{De} is the electron Debye length given by $\sqrt{T_{e0}/(4\pi e^2 n_{e0})}$. Solving Eq. 4.5 can be accomplished by utilizing a finite-difference relaxation algorithm.

4.3.1 Numerically solving the Poisson equation

Equation 4.5 can be discretized in space using the following transformation.

$$\left(\frac{1}{\Delta \xi^2} + \frac{1}{\xi_i \Delta \xi}\right) \tilde{\phi}_{i+1} + \left(\frac{-2}{\Delta \xi^2}\right) \tilde{\phi}_i + \left(\frac{1}{\Delta \xi^2} - \frac{1}{\xi_i \Delta \xi}\right) \tilde{\phi}_{i-1} = \Lambda^2 \left(1 + \frac{\gamma-1}{\gamma} \tilde{\phi}_i\right)^{\frac{1}{\gamma-1}} - \zeta_i \Lambda^2, \quad (4.6)$$

where $\Delta \xi$ is the normalized grid spacing and ζ_i is a constant whose value is either 0 or 1 according to Eq. 4.5. The matrix form of this system of equations is the following equation

$$A\tilde{\Phi} = \Lambda^2 \left(1 + \frac{\gamma-1}{\gamma} \tilde{\Phi}\right)^{\frac{1}{\gamma-1}} - Z, \quad (4.7)$$

where

$$A = \begin{pmatrix} \frac{2}{\Delta\xi^2} & \frac{1}{\Delta\xi^2} - \frac{1}{\xi_1\Delta\xi} & 0 & 0 & 0 & \dots & 0 \\ \frac{1}{\Delta\xi^2} + \frac{1}{\xi_2\Delta\xi} & \frac{2}{\Delta\xi^2} & \frac{1}{\Delta\xi^2} - \frac{1}{\xi_2\Delta\xi} & 0 & 0 & \dots & 0 \\ 0 & \frac{1}{\Delta\xi^2} + \frac{1}{\xi_3\Delta\xi} & \frac{2}{\Delta\xi^2} & \frac{1}{\Delta\xi^2} - \frac{1}{\xi_3\Delta\xi} & 0 & \dots & 0 \\ \ddots & \ddots & \ddots & \ddots & \ddots & \ddots & \ddots \end{pmatrix} \quad (4.8)$$

$$\tilde{\Phi} = \begin{pmatrix} \tilde{\phi}_1 \\ \tilde{\phi}_2 \\ \tilde{\phi}_3 \\ \vdots \end{pmatrix} \quad (4.9)$$

and

$$Z = \begin{pmatrix} \zeta_1\Lambda^2 \\ \zeta_2\Lambda^2 \\ \zeta_3\Lambda^2 \\ \vdots \end{pmatrix}. \quad (4.10)$$

A relaxation algorithm can be used to get a solution for Eq. 4.7. Relaxation algorithms start with an initial guess that is corrected after every iteration until it converges to a solution. Consider an initial guess, Φ_g , that differs from the real solution, Φ_{real} by ϵ , i.e.

$$\Phi_{real} = \Phi_g + \epsilon. \quad (4.11)$$

The real solution should satisfy Eq. 4.7. Thus, Eq. 4.7 and Eq. 4.11 will

yield

$$A(\tilde{\Phi}_g + \epsilon) = \Lambda^2 \left(1 + \frac{\gamma-1}{\gamma} (\tilde{\Phi}_g + \epsilon) \right)^{\frac{1}{\gamma-1}} - Z. \quad (4.12)$$

The right-hand side can be linearized by doing a series expansion where

$$\left(1 + \frac{\gamma-1}{\gamma} (\tilde{\Phi}_g + \epsilon) \right)^{\frac{1}{\gamma-1}} \approx \left(1 + \frac{\gamma-1}{\gamma} \tilde{\Phi}_g \right)^{\frac{1}{\gamma-1}} + \frac{\epsilon}{\gamma} \left(1 + \frac{\gamma-1}{\gamma} \tilde{\Phi}_g \right)^{\frac{2-\gamma}{\gamma-1}} + \mathcal{O}(\epsilon^2). \quad (4.13)$$

Thus, an approximate solution for ϵ is obtained by using Eq. 4.13 in Eq. 4.12 and isolating ϵ , which gives the following equation

$$\left[A - \text{diag} \left(\frac{\Lambda^2 \left(1 + \frac{\gamma-1}{\gamma} \tilde{\Phi}_g \right)^{\frac{2-\gamma}{\gamma-1}}}{\gamma} \right) \right] \epsilon = \Lambda^2 \left(1 + \frac{\gamma-1}{\gamma} \tilde{\Phi}_g \right)^{\frac{1}{\gamma-1}} - Z - A\tilde{\Phi}_g. \quad (4.14)$$

Equation 4.14 can be solved for ϵ using matrix inversion methods, Gaussian elimination, LU decomposition, or the tridiagonal algorithm.

Once the value for ϵ is known, an updated guess is calculated using the following recursion formula

$$\tilde{\Phi}_g^{i+1} = \tilde{\Phi}_g^i + \epsilon^i. \quad (4.15)$$

This process is repeated until the desired level of convergence is reached.

4.3.2 Electron distribution of a solid sphere using the numerical solution

For the following calculations, the polytropic index is $\gamma = 5/3$. The electron densities obtained by solving Eq. 4.5 are compared with three-dimensional (3D) particle-particle (PP) simulations of electrons. The 3D-PP simulations used 10,000 particles, and the ions are assumed to be a static uniform cloud of charge. The electrons are initially uniformly distributed in space. Their velocities are randomly given according to the Maxwell-Boltzmann distribution. The electron temperature is set to match the desired Λ . The time-averaged electron density is calculated once the system has reached equilibrium. This 3D-PP simulation used the following equations motion

$$\frac{d\vec{x}_i}{dt} = \vec{v}_i \quad (4.16)$$

$$\frac{d\vec{v}_i}{dt} = \sum_{j=1}^N \frac{e^2}{|\vec{x}_i - \vec{x}_j|^3} (\vec{x}_i - \vec{x}_j) - e\vec{E}(r), \quad (4.17)$$

where q is the particle charge, x is the position vector, v is the velocity vector, N is the total number of particles, and $\vec{E}(r)$ is the electric field due to the ion cloud. The equation for $\vec{E}(r)$ is given by

$$\vec{E}(r) = \begin{cases} \frac{eNr}{R_0^3} \hat{r} & , r < R_0 \\ \frac{eN}{r^2} \hat{r} & , r \geq R_0 \end{cases} \quad (4.18)$$

where R_0 is the radius of the ion cloud. Because 3D-PP simulations calculate the binary Coulomb forces for each pair of particles, they do not soften the Coulomb collision term unlike mesh-based codes like particle-in-cell (PIC). Thus, using a 3D-PP code will also show if the model is accurate even when electron-electron collisions are considered.

Figure 4.1 shows the comparison of $\tilde{n}_e = n_e/n_{e0}$ from the 3D-PP simulation results and the numerically solved Poisson's equation. The results both results agree very well. Thus, the numerical solution can accurately predict the behavior of electron density bound to the ion sphere. Interestingly, the number of electrons bound to the ion sphere decreases with increasing temperature (Λ decreases). For high temperatures ($\Lambda = 1$), most of the electrons have escaped, and only a small fraction is bound to the ion sphere.

To further verify the accuracy of the model, the charge separation, β , of the electrons and ion sphere was calculated and compared with a semi-analytic fitted model described by Peano et al. [11], which is given by $\beta = F_{2.6}(2.573/\Lambda)$, where $F_\mu(x) = x/(1 + x^\mu)^{1/\mu}$. In this case, the charge separation β is the ratio of the total charge of the electrons within the ion sphere to the ion sphere's total charge. The results show that our numerical results are almost identical to Peano's fitted model (Fig. 4.2), confirming that this model is suitable for predicting the equilibrium electron density for finite-sized ion spheres.

Although the current method closely reproduces the results of Peano

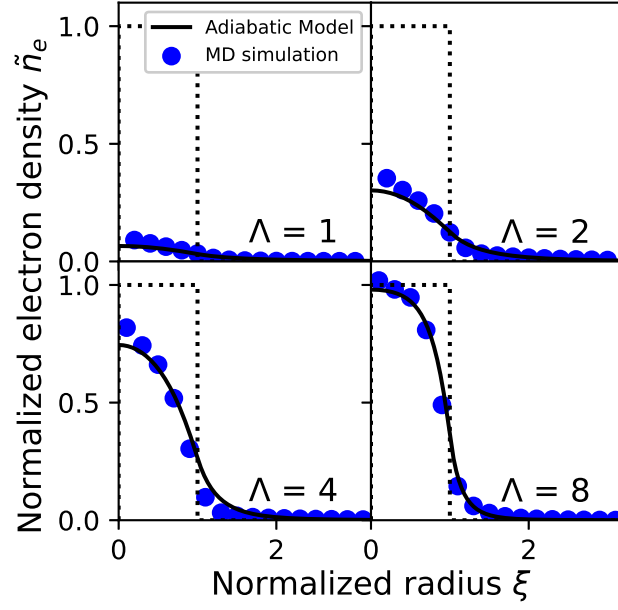


Figure 4.1: Calculated normalized electron density profile for the model and MD simulations for different values of Λ .

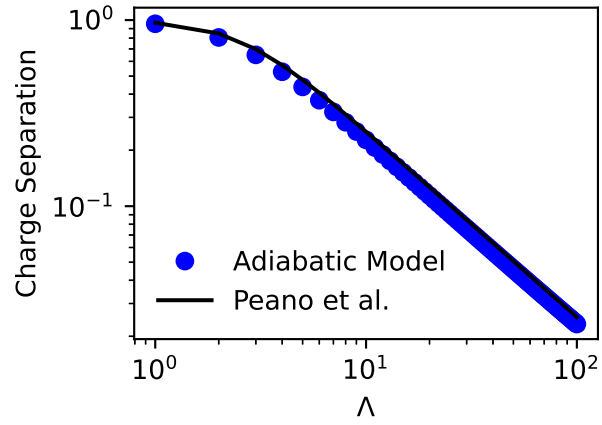


Figure 4.2: Calculated charge separation from the numerically solved Poisson's equation and the model derived by Peano et al. [11]

et al., there is a significant difference in calculating the electron distributions. In the method of Peano et al., the electron density is calculated using the following equations [11]

$$M \frac{\partial^2 r}{\partial t^2} = -Ze \frac{\partial \phi}{\partial r}, \quad (4.19)$$

$$\frac{1}{r^2} \frac{\partial}{\partial r} \left(r^2 \frac{\partial \phi}{\partial r} \right) = 4\pi e (n_e - Zn_i), \quad (4.20)$$

$$n_i(r) = n_{i0}(r_0) \frac{r_0^2}{r^2} \left(\frac{\partial r}{\partial r_0} \right)^{-1}, \quad (4.21)$$

$$n_e = \frac{1}{4\pi r^2} \int \rho_0(\epsilon_0) \mathcal{P}(r, \epsilon; \{\phi\}) d\epsilon_0, \quad (4.22)$$

$$\frac{d\epsilon}{dt} = -e \int \frac{\partial \phi}{\partial t} \mathcal{P}(r, \epsilon; \{\phi\}) dr, \quad (4.23)$$

where M is the mass, ϕ is the electrostatic potential, ϵ is the electron energy, ρ is the energy distribution of the electrons, and $\mathcal{P}$ is the probability of finding an electron at radius r and energy ϵ . In this model, multiple quantities need to be calculated, and the electron distribution needs to evolve in time before it reaches equilibrium. Thus, this method is computationally expensive when compared to our current model. Iteratively solving for Φ by using Eq. 4.14 and Eq. 4.15 leads to quick convergence depending on the initial guess. When the initial guess is close to the solution, the algorithm converges within one or two iterations. When the initial guess is set to 0, the algorithm will converge in not more than five iterations. Thus, the advantage of our method is its quick convergence and simple implementation.

4.4 Electron occupancy in the cavities of hollow spheres

The Poisson equation for hollow spheres with inner radii of R_{in} and an outer radii of R_{out} , can be expressed as,

$$\frac{\partial^2 \tilde{\phi}}{\partial \xi^2} + \frac{2}{\xi} \frac{\partial \tilde{\phi}}{\partial \xi} = \begin{cases} \Lambda^2 \left(1 + \frac{\gamma-1}{\gamma} \tilde{\phi}\right)^{1/(\gamma-1)} - \Lambda^2, & 1 \leq \xi \leq \xi_{ar} \\ \Lambda^2 \left(1 + \frac{\gamma-1}{\gamma} \tilde{\phi}\right)^{1/(\gamma-1)}, & otherwise \end{cases} \quad (4.24)$$

where $\xi = r/R_{in}$ and $\xi_{ar} = R_{out}/R_{in}$. In this section, all normalizations involving R_0 in the previous section are replaced with R_{in} .

Figure 4.3 shows the electron density profiles. As λ_{De} of the electrons increases from $R_{in}/8$ to $R_{in}/2$, the number of electrons inside the cavity increases. This behavior is expected because cold electrons are more likely to stay close to the ion sphere. However, increasing λ_{De} past a certain point decreases the electron distribution because most of its electrons have enough energy to escape the ion sphere completely. Contrary to the semi-infinite case where the optimum number of electrons bound inside the cavity occurs when $\lambda_{De} \rightarrow \infty$ [6], finite spheres will have optimum Λ values that maximize the number of bound electrons found inside the cavity. Additionally, the optimum value of Λ should depend on $\xi_{ar} = R_{out}/R_{in}$ because, as $\xi_{ar} \rightarrow \infty$, $\Lambda \rightarrow 0$ should maximize the

number of bound electrons inside the cavity. Therefore, Λ_{opt} can be expressed as

$$\Lambda_{opt}(\xi_{ar}) = \operatorname{argmax}_{\Lambda} \int_0^1 n_e(\xi, \Lambda, \xi_{ar}) \xi^2 d\xi \quad (4.25)$$

where Λ_{opt} is the optimum value of Λ for a given ξ_{ar} .

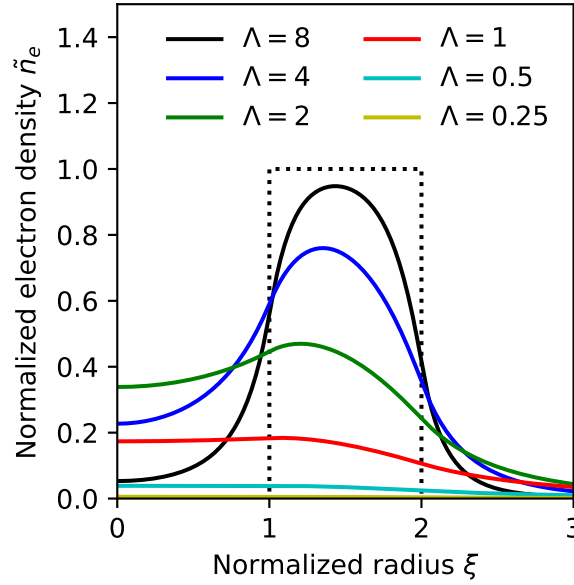


Figure 4.3: Electron density profile of a hollow target for various values of Λ . Dotted line denotes the ion density profile.

Figure 4.4 shows that as the hollow sphere thickness decreases ($\xi_{ar} \rightarrow 1$), $\Lambda_{opt} \rightarrow \infty$. Additionally, it shows that $\Lambda_{opt} \rightarrow 0$ for $\xi_{ar} \rightarrow \infty$. The power-law fit $\Lambda_{opt} = 2.875(\xi - 1)^{-0.5257}$ agrees well with the numerically calculated values for Eq. 4.25. Additionally, as the target thickness

increases, the optimum number of bound electrons inside the cavity, $N_e(\xi \leq \xi_{ar})$, increases. This is an important factor to consider when designing targets for MBI. The compression ratio of MBI depends on the total number of bound electrons inside the cavity. Therefore, although thinner targets need lower electron temperatures to obtain the optimal number of bound electrons inside the cavity, the number of bound electrons would also be much smaller.

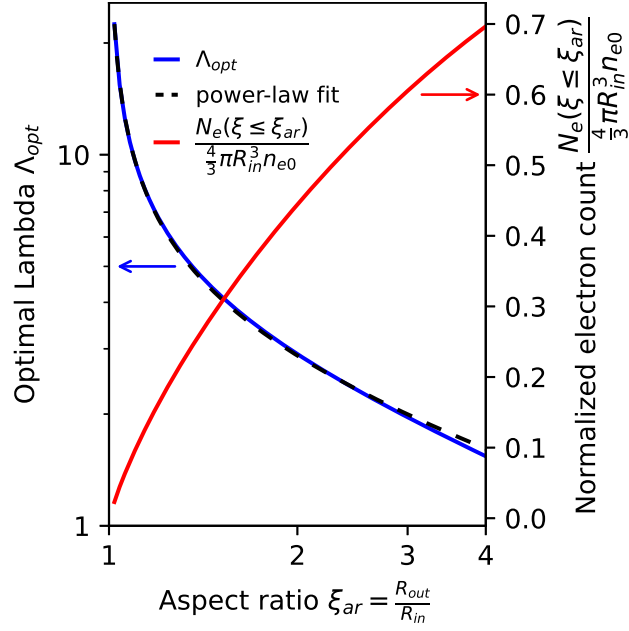


Figure 4.4: Numerically calculated Λ_{opt} as a function of ξ_{ar} along with the power-law fit and the normalized number of electrons within the cavity.

4.5 MBI Simulations

The current model does not diverge as $r \rightarrow \infty$. Thus, it can be used to simulate the dynamics of a finite MBI target. To simulate the MBI dynamics of a finite spherical shell, a hybrid one-dimensional (1D) PP code was implemented. Similar to previous works [2, 12, 63], every computational particle is a spherical shell.

In the simulation, only the protons were treated as computational particles. The electron density was directly calculated using the electron model discussed previously. Each particle trajectory was solved using the equation of motion given by

$$\frac{dp_i}{dt} = \sum_{r_j < r_i} \frac{Q_i Q_j}{r_j^2} + \frac{Q_i^2}{2r_i^2} + Q_i E(r_i) \quad (4.26)$$

where the simulation particles' charge, momentum, and radius are expressed as Q , p , and r , respectively. $E(r)$ is the electric field of the electrons, and it is calculated by solving Eq. 4.2 and Eq. 4.4 at every time-step. Using a mesh-based algorithm to calculate the electron density and treating the protons as computational particles maintain the overall accuracy of the simulation while reducing the total number of simulation particles needed because particles to simulate the electrons are not necessary. Although particle-in-cell (PIC) simulations are common when simulating laser plasma, it is difficult to implement a (PIC) in

spherical MBI without consuming a large amount of computational resources. To properly simulate a system, multiple things need to be considered in PIC codes. Two of the most basic considerations are the time-step, Δt , and the grid spacing, Δx . In order to ensure proper convergence to a solution, the time step in the PIC code should follow the Courant–Friedrichs–Lewy (CFL) condition given by

$$\frac{\Delta x}{\Delta t} \leq c, \quad (4.27)$$

where c is the speed of light in vacuum. Additionally, as a general rule of thumb the grid spacing should be comparable to the Debye length, i.e. $\Delta x \sim \lambda_D$, to avoid numerical instabilities [64]. In MBI, the ions halt and compress to minimum radius of ~ 1 nm and a density of up to 10^5 times solid density ($n_{solid} = 5 \times 10^{22} \text{ cm}^{-3}$) [6]. This makes the final stages of the implosion particularly difficult to simulate using PIC methods because the Debye length is too small. The problem of having a small spatial resolution also affects the time-step. Because the time-step needs to satisfy the CFL condition, small grid spacings need proportionally small time-steps. Thus, such a simulation can cost an unreasonable amount of resources especially if it is three-dimensional. For a one-dimensional spherical PIC, it cannot accurately simulate the implosion dynamics because in PIC, the forces are calculated using the fields on the grid. This softens the inter-particle interactions especially

at close distances. Because particle-particle methods directly calculate the binary force from every particle, they do not have the same problem as PIC. However, in cases where the laser-plasma dynamics is essential and the system does not compress to a high degree, PIC calculations should be considered.

The targets in the simulations were assumed to be composed of pure hydrogen. The electron temperatures were assigned to be the optimum temperature as given by Eq. 4.25. In all simulations, a cavity radius of $1\text{ }\mu\text{m}$ was set, and the initial density was $n_i = 5 \times 10^{22}\text{ cm}^{-3}$. As a benchmark, in ideal MBI, when the initial density is $5 \times 10^{22}\text{ cm}^{-3}$ and the bubble radius is $1\text{ }\mu\text{m}$, the bubble will be compressed to a minimum radius of 0.8 nm . The compression rate, in this case, is $\sim 10^5$.

4.6 Simulation Results

The trajectories of the MBI simulations are shown in Fig. 4.5. The results show that MBI occurs even for a finite hollow sphere. Furthermore, all three cases considered in Fig. 4.5 show indications of proton flashing. However, unlike the typical MBI results for semi-infinite geometries, in these cases, the charge separation resulting from high electron temperatures causes the CE of the outer layer. The expanding outer layer is also a source of accelerated ions along with the flashed ions from the imploding inner layer. Interestingly, the implosion time decreases as the

thickness of the hollow sphere increases (Fig. 4.5).

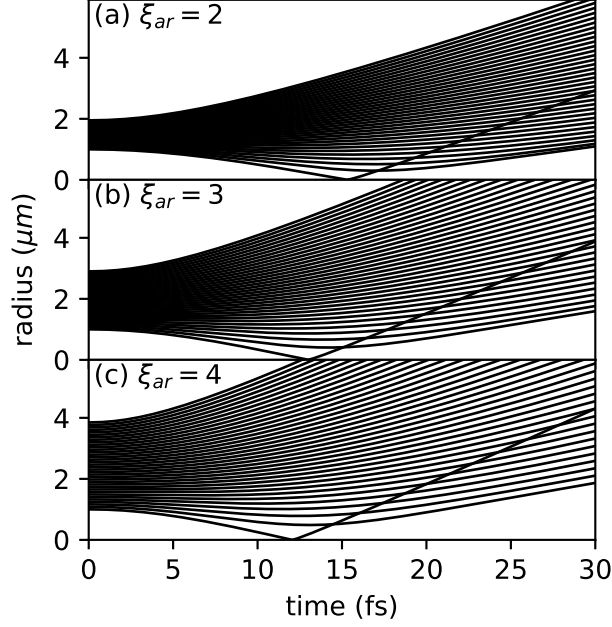


Figure 4.5: Trajectory of a hollow sphere with 1- μm cavity undergoing MBI for different aspect ratios: (a) $\xi_{ar} = 2$, (b) $\xi_{ar} = 3$, and (c) $\xi_{ar} = 4$.

The maximum number of bound electrons inside the cavity increases (Fig. 4.4) as the aspect ratio increases. Because the MBI is implosion velocity is heavily affected by the number of bound electrons inside the cavity, more electrons present within the cavity should increase the implosion velocity of the ions, implying a shorter implosion time [6].

In all three cases, a density compression ratio of roughly 10^4 times that of the solid density (Fig. 4.6) is achieved. Although this is one order of magnitude lower than the 10^5 compression ratio predicted using an ideal semi-infinite system (Eq. 4.1), Fig. 4.6 shows that high compression

ratios are possible even with a finite target. More importantly, the agreement between the density profile and the theoretical model given by Eq. 4.1 improves as the aspect ratio increases.

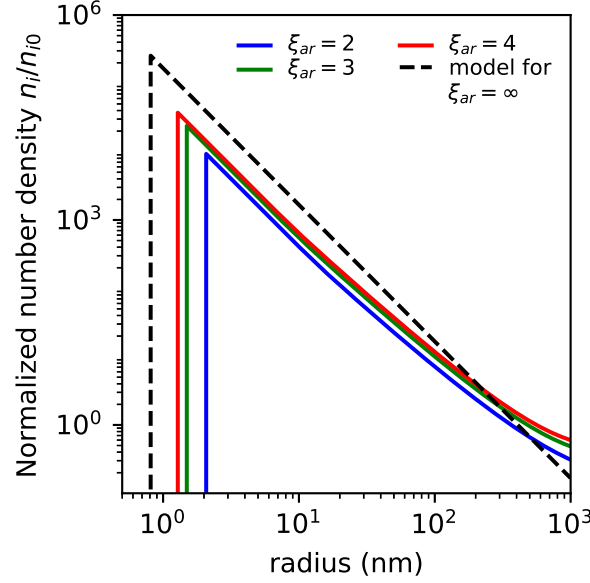


Figure 4.6: Density profile during maximum compression of a hollow sphere with 1- μm cavity undergoing MBI for different aspect ratios. The dashed line is the model given by Eq. 4.1.

Evidence of ion flashing described by earlier studies is shown in the energy distribution plots (Fig. 4.7). The small spikes in the graphs agree well with the energies of the inner layers. Thus, these spikes are due to the flashing of the inner ion layers. An interesting observation is that most flashed ion layers have lower energies than the outermost coulomb exploding layer. Hence, distinguishing the flashed layers from the normal Coulomb exploding layer is difficult if the flashed layers

have comparable energies to the Coulomb exploding layers. However, for $\xi_{ar} = 2$, the innermost layer has an energy higher than that of the outermost Coulomb exploding layer. Thus, detecting ion flashing in thick targets ($\xi_{ar} > 2$) will be very difficult. The maximum energy of the Coulomb exploding front, $\mathcal{E}_{ce}$, of a hollow sphere with a inner radius of R_{in} and outer radius of R_{out} is given by

$$\mathcal{E}_{ce} = \frac{4}{3}\pi e^2 n_i \beta \frac{R_{out}^3 - R_{in}^3}{R_{out}}, \quad (4.28)$$

where β is the charge separation and n_i is the ion density. Thus, the maximum energy of the Coulomb exploding front, $\mathcal{E}_{ce}$, is roughly proportional to R_{out}^2 . For example, consider a hollow sphere with a radius of $2\text{ }\mu\text{m}$ and a $1\text{ }\mu\text{m}$ cavity radius. Increasing the outer radius to $4\text{ }\mu\text{m}$ will increase $\mathcal{E}_{ce}$ by roughly a factor of 4. However, according to Fig. 4.4, this change will only give a 50% increase in cavity electrons. Thus, in this case, increasing the outer radius strongly favors the enhancement of the Coulomb exploding front because MBI depends on the number of cavity electrons. Conversely, making the targets thinner will be more detrimental to the energy of Coulomb exploding front because of its $\sim R_{out}^2$ scaling.

Targets with $\xi_{ar} > 2.2$ will not produce flashed ions according to Fig. 4.8 because Λ_{opt} is larger for thicker targets. The charge separation between electrons and ions increases as Λ_{opt} increases. Thus, the total

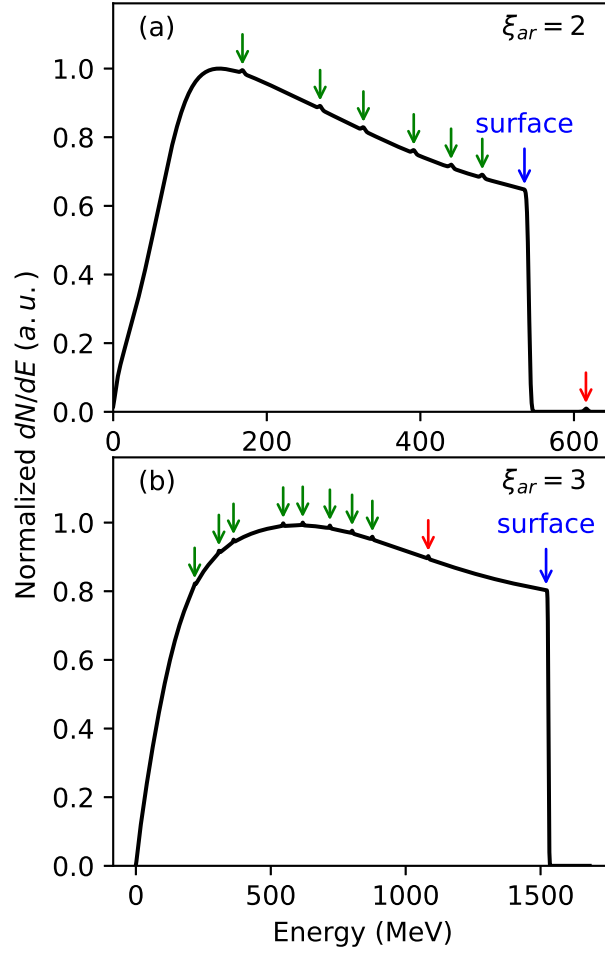


Figure 4.7: Energy distribution of a hollow sphere with 1- μm cavity undergoing MBI for: (a) $\xi_{ar} = 2$ and (b) $\xi_{ar} = 3$. Green arrows indicate flashed ions, the red arrow indicates the innermost layer, and the blue arrow indicates the surface layer.

energy of the Coulomb exploding outer layers also increases. On the other hand, the acceleration of the inner layers is influenced by the total number of bound electrons inside the cavity [6, 7]. Thus, $\xi_{ar} = \frac{R_{out}}{R_{in}} = 2.2$ marks the point where the flashed protons from MBI will no longer have enough energy to overtake the Coulomb exploding front. On the other hand, thinner targets have fewer electrons, according to Fig. 4.4. Thus, if the target is too thin, the number of electrons within the cavity will not be enough to properly initiate the MBI process. In this system, because the system dynamically changes according to not only MBI but also Coulomb explosion, it is difficult to find an analytical estimate for the thickness that provides the best contrast between the flashed protons and the Coulomb explosion front. Figure 4.8 shows that when $\xi_{ar} < 1.3$, the energy ratio between the innermost flashed ion and the outmost Coulomb exploding layer starts to decrease. Hence, the maximum contrast between the surface ions and the flashed layers occurs when $\xi_{ar} = 1.3$. At this point, the target is thin enough to minimize $\mathcal{E}_{ce}$, but it still has enough cavity electrons to initiate MBI and produce flashed ions.

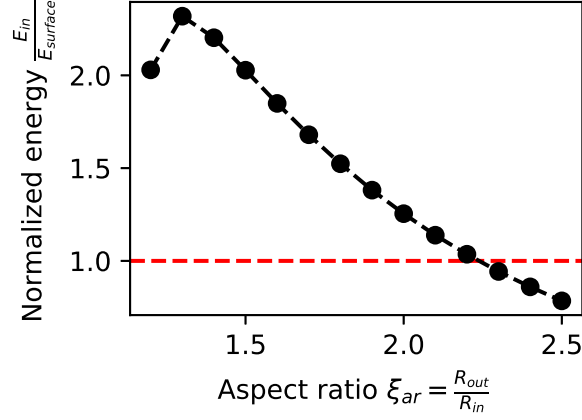


Figure 4.8: The ratio of the innermost ion energy, E_{in} , to the surface ion energy, $E_{surface}$ versus ξ_{ar} .

4.7 Section remarks

The one-dimensional hybrid code that was developed to simulate MBI for finite spheres was successfully employed to simulate the MBI dynamics for finite hollow spheres. It confirmed that MBI occurs for finite spheres. For all simulated targets, the density compression ratio is roughly $\sim 10^4$. Interestingly, thicker targets were observed to produce higher density compression. However, the coulomb exploding outer layers make it more difficult to observe ion flashing in thicker targets.

A key observation from the simulation results is that the target's aspect ratio plays an important role in the different applications of MBI. As observed, targets with high aspect ratios have better compression ratios. Unfortunately, ion flashing is difficult to detect in high aspect

ratio targets because of the Coulomb exploding front. Increasing the target thickness significantly favors increasing the energy of the Coulomb exploding front. Targets with large aspect ratios also have a lower Λ , requiring a higher intensity laser. This implies that while thick targets are good at generating an ultra-dense core for gamma-ray lensing and pair creation, thinner targets are more suited for ion acceleration applications for MBI because, as long as $\xi_{ar} < 2.2$ and Λ_{opt} is used, there will be contrast between the flashed ions and the Coulomb explosion front. Thus, for finite MBI targets, the designed thickness of the MBI targets should be tailored to the specific application to be observed.

In Chapter 3 and 4, Coulomb explosion and microbubble implosion (MBI) were discussed. In Coulomb explosion, a large charge separation creates a highly positive ion sphere, which results in the rapid outward acceleration of the ions due to the Coulomb field. In MBI, a high electron temperature leads to an electron-filled cavity in a hollow spherical target, creating a highly negatively charged region. This negative region violently attracts the ions initiating an implosion process that results in a highly dense core and energetic flashed ions.

Finally, microtube implosion (MTI) is another phenomenon that occurs in hollow targets. MTI is observed in hollow micron-sized cylindrical targets. In MTI, the electron temperature is lower, thus only an electron sheath close to the cavity walls is generated. This allows the hydrodynamic expansion of electrons and ions towards the center. In the presence of a seed magnetic field, the particle trajectories are

bent, generating an amplified magnetic field in the megatesla domain. MTI typically requires a seed magnetic field. However, if the target can generate an initial magnetic field serving as the seed, the setup would be more straightforward, and megatesla-level magnetic fields occurs without needing an external seed magnetic field.

The target should generate a polar current in the $\hat{\theta}$ -direction to create an intense magnetic field in the z-direction. Because electrons are lighter and have a shorter timescale, guiding the electrons to follow a curved path to generate the polar current is essential. MBI and MTI use a negatively charged region to direct the motion of the ions. However, it should also be possible to use a positively charged region to direct the motion of electrons. As discussed in Section 3.1, when a target is comparable to the laser skin depth, most of its electrons will be removed, resulting in a large charge separation. If a target has a segment thin enough to be comparable to the laser skin depth, the thin region will be highly charge separated upon laser irradiation. The highly charge separated region can attract the bulk electrons towards it, and if the surface is curved, a strong magnetic field can be generated.

In this chapter, we will discuss how charge separation and skin depth can direct electron flow, and produce a strong magnetic field. This strong magnetic can be used as an initial seed magnetic field that can be amplified to an even stronger field.

The discussion and figures in this chapter were lifted from M. A. H.

Zosa, Y.J. Gu and M. Murakami, Sub-megatesla Magnetic field generation using paisley targets by femtosecond laser-plasma interactions, *Applied Physics Letters*, 2022.

5.1 Magnetic Field Generation

With the modern developments in relativistic femtosecond lasers [34,65] and microstructure fabrication [66,67] new regimes of high energy density physics are being discovered [68–70]. Recent studies have used relativistic femtosecond lasers to investigate ion acceleration [71], magnetic field generation [72], ultra-high density compression [6, 7], and pair-creation [5]. Generating magnetic fields on the sub-megatesla scale has garnered much attention because it enables the study of fundamental phenomena such as magnetic reconnection [15,73,74]. These intense magnetic fields are observed in the accretion disks of black holes, which makes them valuable for laboratory astrophysics experiments [75].

Irradiating an “escargot” target with a laser can produce a strong magnetic field [76,77]. It was used as a magnetic field source in a magnetic reconnection laboratory experiment [75]. Microtube implosion (MTI) is another method for generating intense magnetic fields [3,14]. In MTI, a seed magnetic field is amplified to the megatesla scale through the implosion of the inner layer of a microtube. Its strength is enhanced by a few orders of magnitude [78].

5.2 The Paisley Design

Both setups use different schemes to generate a magnetic field. However, a common factor is in both cases is the formation of a surface current [14, 76, 79]. This chapter will introduce a new design to generate a magnetic field without a seed magnetic field. One way to produce a current is to create a region with large charge separation so that the electrons will flow towards that area. To ensure that an region will have large charge separation, its thickness should be comparable to the laser skin depth. Thus, this region should be considerably thin ~ 100 nm. However, if the entire target is thin, the majority of the electrons will be removed and none will be left to sustain the surface current. Thus, the rest of the target should be much larger than the laser skin depth. Additionally, the surface where the electrons will flow towards should be curved to efficiently produce a magnetic field. Taking all these requirements into consideration, a design using circles would be good. Because circles provide a curved surface, the electrons moving along its surface will naturally produce a magnetic field. If we consider a large circle inscribed with two smaller circles with half the radius of the larger circle as shown in Fig. 5.1(a), the shaded region shows all the qualities required. It has a thin region while the rest of the target is thick, and it has a curved surface that produces a magnetic field once the thin region is charge separated. In three-dimensions, the paisley

target extends along the z-axis like a long cylinder [Fig. 5.1(b)].

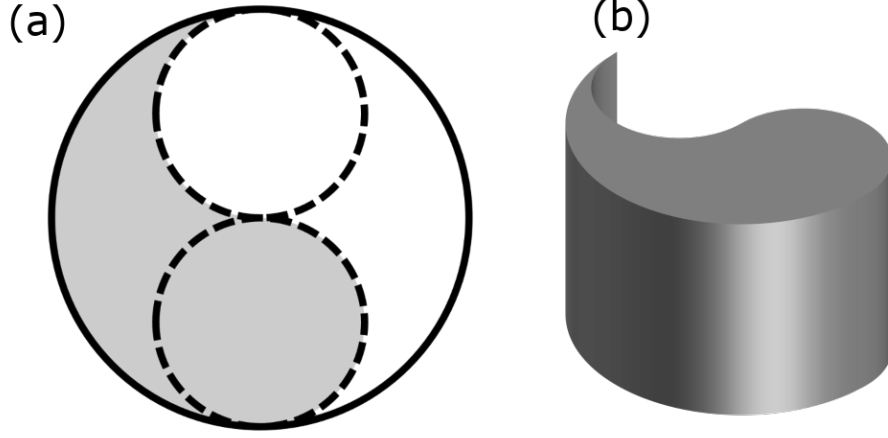


Figure 5.1: (a) Geometric design of the paisley target. The shaded region is the paisley target. Dotted lines are inscribed circles with half the radius. (b) Three-dimensional isometric representation of the paisley target.

The new design called the “paisley design” is described mathematically by the following function

$$f(k) = \begin{cases} -\frac{R_0}{2} (\exp [4\pi i k - \pi/2] + 1) & k \in (0, \frac{1}{4}) \\ -\frac{R_0}{2} (\exp [4\pi i k - \pi/2] - 1) & k \in (\frac{1}{4}, \frac{1}{2}) \\ R_0 \exp (2\pi i k - \pi/2) & k \in (\frac{1}{2}, 1) \end{cases}, \quad (5.1)$$

where R_0 is the radius, and $k \in (0, 1)$. The x- and y-coordinates are the real and imaginary parts of Eq. 5.1, respectively. Figure 5.2(a) is a graphical representation of Eq. 5.1. To generate a magnetic field, surface currents produce a magnetic field on the concave side of the target.

The magnetic field generated on an open area makes it easily accessible. Thus, it is easier to utilize in various experimental designs. Additionally, two or more targets to be connected in a modular fashion, and various arrangements can be used to study different types of magnetic field interactions.

To simulate the magnetic field generation of the paisley design, the 2.5D particle-in-cell (PIC) program “EPOCH” was used [80]. In this case, because the laser-plasma dynamics strongly influence system, PIC simulations are needed. The laser parameters were $\lambda_L = 800$ nm, $I_L = 1 \times 10^{21}$ W cm⁻², and $\tau_L = 100$ fs for the wavelength, peak intensity, and full-width at half-maximum (FWHM), respectively. These parameters were chosen so that the simulations are comparable to similar works affiliated to our research group [14, 15]. The simulations used a $30 \mu\text{m} \times 30 \mu\text{m}$ box with a cell size of $\frac{\lambda_L}{100}$, 100 pseudo-ions and 200 pseudo-electron per cell, and a laser propagating in the +x direction. The target was fully ionized carbon with a density of 1×10^{23} cm⁻³.

5.3 Magnetic Field Generation

Local charge separation is essential to generate surface current in the paisley design. Because the thickness gets much thinner closer to the tip, the charge separation will be greater closer to the tip compared to the main body [Fig. 5.2(b)]. The laser can strip most of the electrons from

the thinner sections of the paisley target. However, it cannot induce much charge separation to the thicker areas. This causes electrons to flow towards the highly charge separated tip [Fig. 5.2(c)]. Paisley target's curvature creates a positive (negative) magnetic field on its concave (convex) side [Fig. 5.2(d)]. Unsurprisingly, the polarity of the magnetic fields flips if the target is rotated 180° about $y = 0$.

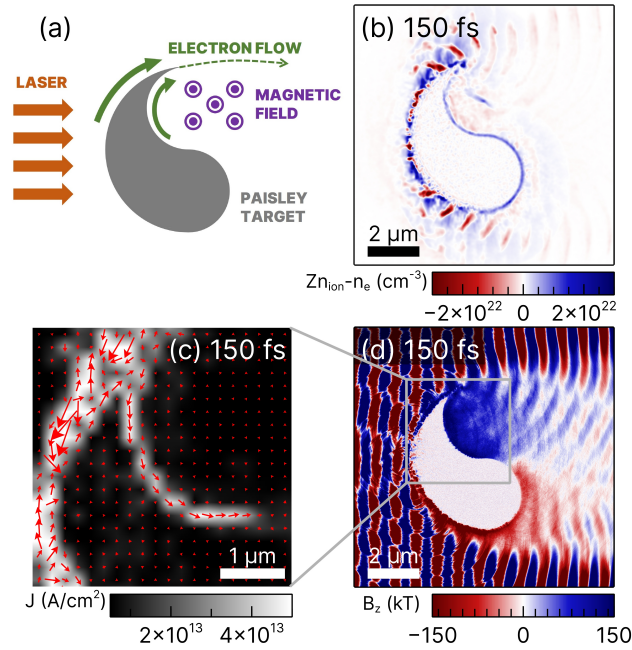


Figure 5.2: (a) Illustration of the magnetic field generation mechanism. At $t = 150$ fs, (b) charge separation profile, (c) vector diagram for the net current, and (d) magnetic field profile along the z-axis.

A simple analytic model was developed to predict magnitude of the magnetic field in kT. The magnetic field strength should be $B_z \sim j_e R_0$, and the current density can be estimated by, j_e , is $j_e \sim n_{he} c$. The hot-electron density, n_{he} , is related to I_L by $n_{he} = \eta_a I_L / \mathcal{E} c$ [81], where

η_a is the absorption efficiency, I_L is the laser intensity, $\mathcal{E}$ is the average kinetic energy, and c is the speed of light in vacuum. When considering relativistic electrons, the average kinetic energy, $\mathcal{E}$, is approximately $3T_e$, where T_e is the electron temperature. The electron temperature can be estimated using the ponderomotive scaling [82]. Thus, the model will be

$$B_z(\text{in kT}) = 30.3 \frac{\eta_a \sqrt{I_{L20}} R_{0\mu m}}{\lambda_{\mu m}}, \quad (5.2)$$

where I_{L20} is the laser intensity normalized to $10^{20} \text{ W cm}^{-2}$, η_a is the absorption efficiency, $R_{0\mu m}$ is characteristic radius given by Eq. 5.1 in μm , and $\lambda_{\mu m}$ is the laser wavelength in μm . From Fig. 5.3(a), the peak magnetic field increases with the laser intensity. The FWHM for the magnetic field is roughly twice the laser's FWHM ($\sim 2\tau_L$). Because the laser generates the charge separation across the tip, once it leaves, the magnetic field will start to decay. Without a highly positively charged region to direct the surface, the magnetic field cannot be sustained under this mechanism. Additionally, the PIC simulation results agree well with Eq. 5.2 for $\eta_a = 0.4$ [Fig. 5.3(b)]. From the model, an absorption efficiency of 0.8 or higher is required to have megatesla scale magnetic fields for $I_L = 10^{22} \text{ W cm}^{-2}$.

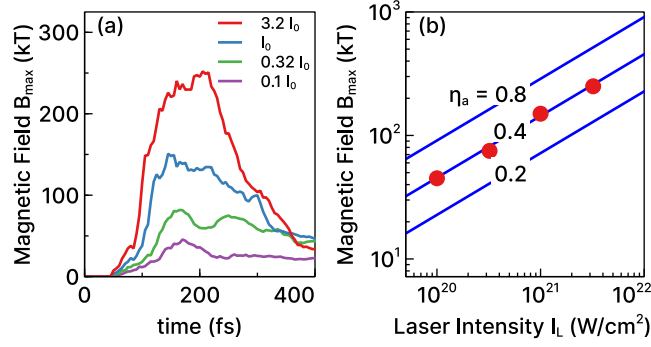


Figure 5.3: (a) Maximum magnetic field generated by the paisley target vs time for various laser intensities ($I_0 = 10^{21} \text{W cm}^{-2}$). (b) Simple analytic model (blue lines) for varying η_a and the maximum magnetic field from the simulations results (red dots).

5.4 Double Paisley

Interestingly, when comparing Fig. 5.3(a) with the simulation parameters, the peak magnetic field coincides with the laser maximum. During the laser maximum, maximum electron heating takes place. Thus, this is the time when the target experiences the largest charge separation across the tip, inducing the largest surface current and consequently, the largest magnetic field amplitude. Additionally, the magnetic field sharply drops once the laser has disappeared. This results in a relatively short magnetic field lifetime that is comparable to the laser's FWHM. However, using two paisley targets will extend the magnetic field's lifetime [Fig. 5.4(c)]. For the double paisley case, the two lasers hit two of paisley targets from the $+x$ and $-x$ -direction. The generated magnetic field covers a much larger area [Figs. 5.4(a) & (b)]. Additionally, the

magnetic field lifetime is longer than the laser pulse duration. The electrons form a current loop [Fig. 5.4(f)] in the central region. The current loop sustains the magnetic field for more than a picosecond after the laser pulse ended. This stability is similar to the “escargot” target, where the magnetic field displays the same characteristics as a θ -pinch magnetic field [76]. A θ -pinch magnetic field is stable, so the defining characteristic for the timescale will be target expansion. The confined region formed by the two paisley targets delays the expansion of the current loop. From the electron collision frequency [79], the characteristic time scale of the electron collision, ν_e^{-1} , is several picoseconds long. Thus, dissipation through Coulomb collisions are negligible for the duration of the magnetic field.

5.5 Simplified Design

The complex design of the paisley targets is challenging to fabricate. However, a simpler design can have comparable magnetic field generation. The main magnetic field generation mechanism is to create a highly charged separated region, which induces a surface current to flow along a curved path. Thus, as long as we have a relatively thick target with a thin tip along a curved surface, the mechanism and plasma dynamics should be similar upon laser irradiation. Thus, a target such

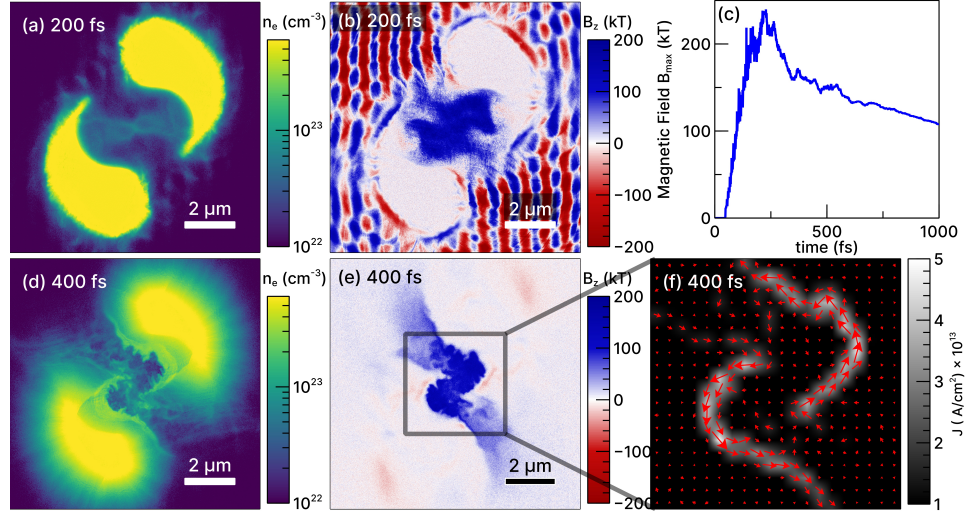


Figure 5.4: At $t = 200$ fs, (a) the electron density profile and (b) the magnetic field profile, B_z . (c) Maximum magnetic field as a function of time. At $t = 400$ fs (d) the electron density profile, (e) the magnetic field profile, and (f) the current vector diagram.

as Figure 5.5(a) can be used as a simplified paisley target because it satisfies the requirements. Although the design differs from the original one, the core concept of utilizing the thickness gradient to guide the surface current remains. Despite a major change in its appearance, the magnetic field strength produced by the simplified target is comparable to the paisley design [Fig. 5.5].

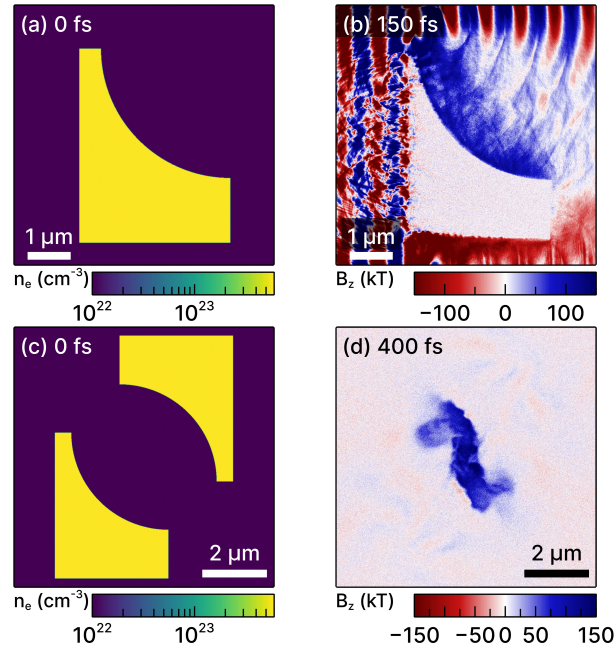


Figure 5.5: A simplified single paisley design: (a) Initial electron density profile and (b) magnetic field profile at $t = 150$ fs. A simplified double paisley design: (c) Initial electron density profile and (d) magnetic field profile at $t = 400$ fs.

5.6 Four-sided simplified paisley setup

The paisley target can generate 100-kT-level magnetic fields without a seed magnetic field. However, it is also possible to use the paisley targets in a configuration wherein it can utilize the same principles of MTI and generate megatesla-level magnetic fields. Figure 5.6 shows the initial configuration using four simplified paisley targets. Four 10^{21} W cm⁻² lasers will irradiate the target, creating a surface current around the cavity wall. Because the electrons will flow toward the tip of the targets, the generated magnetic field will be along the +z-axis. This induced magnetic field will act as a seed field for the imploding electrons and ions, bending their trajectories to produce an amplified field similar to MTI.

Figure 5.7 shows that the amplified magnetic is ~ -0.7 MT. Interestingly, the amplified magnetic field is opposite the magnetic field generated by the surface currents, indicating that polarity switching may have occurred (Section 2.8). Thus, the electrons are the primary source of the amplified field. Figure 5.8 further confirms the polarity switching mechanism because it shows the dense electron density around the center associated with electrons orbiting around an ion core. Finally, the net charge distribution [Fig. 5.9] shows a highly positive region at the center of the target, indicating that the electrons could be orbiting around the imploded ion core.

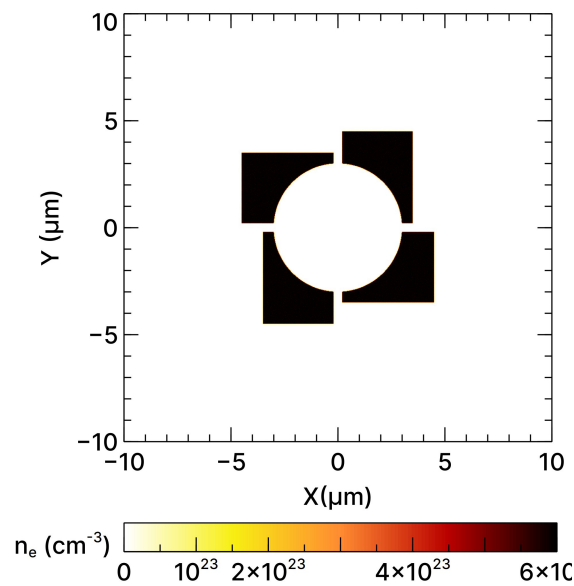


Figure 5.6: Initial electron density profile. The cavity radius is 3 μm .

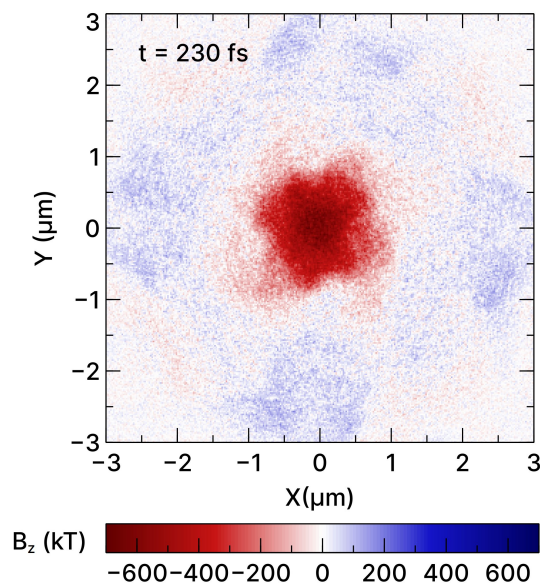
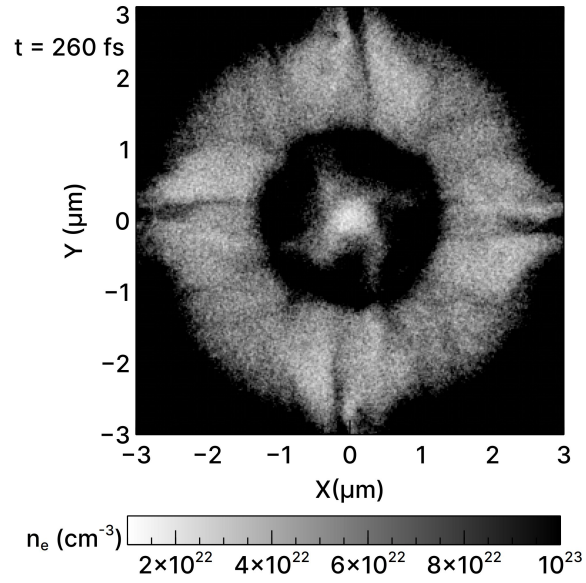
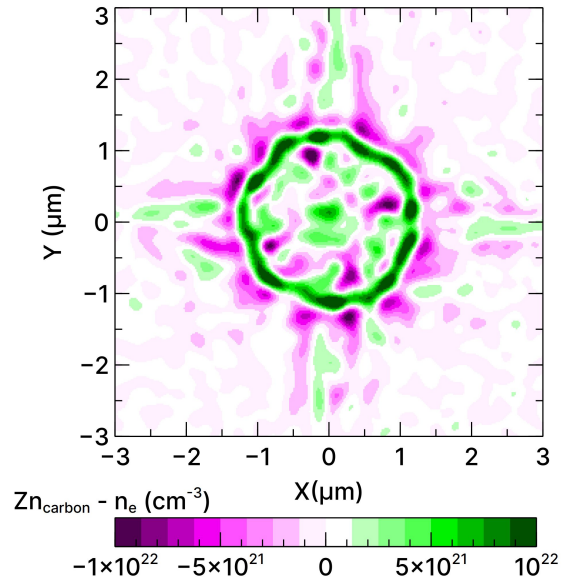


Figure 5.7: The magnetic field profile of the 4 paisley configuration at $t = 260$ fs. The maximum magnetic field amplitude is roughly -700 kT.

Figure 5.8: Electron density profile of the 4 paisley configuration at $t = 260$ fsFigure 5.9: The net charge profile of the 4 paisley configuration at $t = 260$ fs

5.7 Section remarks

The paisley target has multiple applications because of its ability to generate a magnetic field without a seed. Additionally, its modular design gives it more flexibility and further expands its applications. The double paisley target, can generate intense 100-kT magnetic fields that are sustained for more than a picosecond after the laser pulse has ended. Although the current double paisley setup is used to enhance the magnetic field lifetime, inverting one of the paisley targets results in two opposing magnetic fields. This configuration is suitable for studying magnetic reconnection. Additionally, having a four paisley configuration leads to an amplified MT-level magnetic field similar to those generated in MTI.

In summary, a brief introduction to the pertinent regimes were discussed in Chapter 1 and Chapter 2. The CE of hollow nanospheres were investigated and optimized in Chapter 3. The results showed that using a hollow nanosphere with a heavy ion inner layer and an outer layer composed of a heavy and light ion mixture yielded the best coupling efficiency. In this case, the heavy ion inner layer was able to provide more energy to the innermost light ions. Thus, for spherically symmetric CE, adding a heavy ion inner layer improves the energy distribution profile by giving more energy to the slower moving innermost light ions.

In Chapter 4, the MBI of finite nanospheres were investigated and a model to calculate the electron distribution of a hollow ion sphere was formulated. The results in this chapter have shown that MBI implosion occurs for hollow nanospheres and that there is an optimal thickness that would yield the best results for ion flashing. Moreover, thicker

targets will yield much better compression. From these results, it can be seen that the dynamics of finite nanospheres are different from the ideal case of MBI. In an ideal MBI, the outermost layers of the target barely move because the system is considered quasi-neutral. Without the quasi-neutrality condition, the outer layers of the MBI target will expand due to CE. However, when properly optimized, the inner wall ions will still implode toward the center because they will be attracted to the bound electrons within the cavity. This is in stark contrast with hollow nanospheres that will not experience a violent implosion due to their lack of cavity electrons. Thus, target optimization is very crucial for MBI once the quasi-neutrality condition no longer holds.

Finally, Chapter 5 discussed how an intense magnetic field is generated through the novel paisley design. In this design a magnetic field is generated through the surface currents produced by the target structure. The surface currents are heavily influenced by the large charge separation in the tip of the target because of its very thin structure. This causes electrons to flow towards the tip, generating an intense sub-megatesla magnetic field. This paisley design combines the different scales used in the previous chapters. The tip is thin enough to have most of the electrons stripped by the laser. In the paisley design, the thickness gradually increases which means that the charge separation also gradually decreases. In the main body of the target, only the surface electrons can be stripped from the target. The rest of the electrons

remain. Because the tip is highly charge separated, the surface current will disappear once it starts to expand because of CE. However, if two paisley targets are utilized, they would synergize with each other to create a current loop that enhances the lifetime of the magnetic field. With four paisley targets, the magnetic fields produced by the surface currents acts as a seed magnetic field. The fields bend the electron and ion trajectories, initiating an MTI process, amplifying the magnetic field to the megatesla level.

The dynamics of hollow nano- and micro-plasmas is not easily defined. The dominant process in hollow nanospheres is Coulomb explosion because they cannot fill their cavities with enough electrons to initiate a violent ion implosion. Also, targets that are close to the laser skin depth will be significantly more charge separated, making hollow nanosphere even less likely to demonstrate MBI. However, when heated optimally, hollow microspheres will have enough electrons filling their cavities to initiate MBI. Additionally, because the hollow microspheres are also charge separated, the outer layers will Coulomb explode. Thus, while the behavior of the electron distribution can be expressed in terms of dimensionless parameters that scale with target size [Fig. 4.3], some phenomena associated with such behaviors, such as MBI, will only initiate if a critical size is reached (Section 2.5). However, once a target is large enough, the electrons in the cavity can attract the ions to rapidly implode. Conversely, it is also possible to use ions to direct the electron

motion. In the paisley target, a highly charge separated tip is used to direct the flow of the bulk electrons and generate a surface current, resulting in a magnetic field. With four paisley targets, the generated magnetic field is amplified to the megatesla level through the implosion of the electrons and ions. Thus, the four paisley set-up demonstrates that it is possible to utilize charge separation, surface current generation, and implosion to reach megatesla magnetic fields without needing an initial seed magnetic field.

In conclusion, this research demonstrates that hollow targets have various applications determined by their size, electron Debye length, topology, and symmetry. However, the dynamics and phenomena displayed by hollow targets may not necessarily scale with size. Thus, it is important to carefully consider and optimize the target design to achieve the desired phenomenon.



First Author Publications

- **M.A. H. Zosa**, Y. J. Gu, and M. Murakami, 100-KT Magnetic Field Generation Using Paisley Targets by Femtosecond Laser–Plasma Interactions, Appl. Phys. Lett. 120, 132403 (2022).
- **M. A. H. Zosa** and M. Murakami, Microbubble Implosions in Finite Hollow Spheres, Physics of Plasmas 29, 013105 (2022).
- **M. A. H. Zosa** and M. Murakami, Generation of Quasi-Monoenergetic Ions Using Optimized Hollow Nanospheres, Physics of Plasmas 27, 033103 (2020).

Co-authored Publications

- M. M. Rosete, **M.A. H. Zosa**, and R. V. Sarmago, Improved Fabrication of Electrophoretically Deposited $\text{Bi}_2\text{Sr}_2\text{CaCu}_2\text{O}_8$ Films via Supporting Electrolyte Optimization in Ethanol, *J Supercond Nov Magn* 34, 55 (2020).
- M. Murakami, A. Arefiev, **M. A. Zosa**, J. K. Koga, and Y. Nakamiya, Relativistic Proton Emission from Ultrahigh-Energy-Density Nanosphere Generated by Microbubble Implosion, *Physics of Plasmas* 26, 043112 (2019).
- M. Murakami, A. Arefiev, and **M. A. Zosa**, Generation of Ultra-high Field by Micro-Bubble Implosion, *Sci Rep* 8, (2018).
- M. M. Rosete, **M.A. H. Zosa**, and R. V. Sarmago, Effect of Potassium Chloride as a Supporting Electrolyte on the Dispersion Towards the Fabrication of Films by Electrophoretic Deposition of $\text{Bi}_2\text{Sr}_2\text{CaCu}_2\text{O}_8$ in Ethanol, *J Supercond Nov Magn* 30, 1833 (2017).

International Conferences

- **M.A. H. Zosa** and M. Murakami, Ultra-high compression and ion acceleration in finite-sized hollow spherical targets via microbubble implosion, OPIC-HEDS, Yokohama, Japan (2022). Oral

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- **M.A. H. Zosa** and M. Murakami, Enhanced Proton Acceleration Generated via Spherical Symmetry Breakdown During Microbubble Implosion, IFSA, Osaka, Japan (2019). Poster
 - **D. Pan**, **M.A. H. Zosa** and M. Murakami, Coulomb Explosion of Ultrathin Foils, IFSA, Osaka, Japan (2019). Poster
 - **M.A. H. Zosa** and M. Murakami, Quasimonoenergetic proton production through coulomb explosion of spherical nanostructures, OPIC-LANSA, Yokohama, Japan (2017). Oral
 - **M.A. H. Zosa**, R. Sarmago, and M. Murakami, Increased velocity collimation of protons inside a capped nanotube using a Bi-sepcies proton source, First ICMRE, Chengdu, China (2016). Poster

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