Magnetic field generation in a laser-irradiated thin collisionless plasma target by return current electrons carrying orbital angular momentum

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Abstract
Magnetized high energy density physics offers new opportunities for observing magnetic field-related physics for the first time in the laser–plasma context. We focus on one such phenomenon, which is the ability of a laser-irradiated magnetized plasma to amplify a seed magnetic field. We performed a series of fully kinetic 3D simulations of magnetic field amplification by a picosecond-scale relativistic laser pulse of intensity $4.2 \times 10^{18}$ W cm$^{-2}$ incident on a thin overdense target. We observe axial magnetic field amplification from an initial 0.1 kT seed to 1.5 kT over a volume of several cubic microns, persisting hundreds of femtoseconds longer than the laser pulse duration. The magnetic field amplification is driven by electrons in the return current gaining favorable orbital angular momentum from the seed magnetic field. This mechanism is robust to laser polarization and delivers order-of-magnitude amplification over a range of simulation parameters.

1. Introduction
High energy density physics (HEDP) emerged as a new sub-field only about two decades ago, but has already substantially advanced our understanding of materials under extreme conditions and led to the development of several applications, such as those involving energetic particle beams. Much of this progress is due to breakthroughs in laser technology which have enabled drivers capable of depositing energy on a picosecond time scale and creating the high energy density state of matter [1]. Until recently, quasi-static magnetic fields have been of relatively low importance in HEDP research due to the technological challenges associated with generating sufficiently strong macroscopic fields at the laser facilities used for HEDP research.

The recent development of open-geometry, all-optical magnetic field generators [2–4] which are portable to any high-energy laser facility has opened up new regimes relevant to magnetized HEDP to exploration. It is now feasible to experimentally probe laser–plasma interactions with an embedded magnetic field that reaches hundreds of Tesla in strength. As a result, there has been an increased interest in laser-driven, high-energy-density systems embedded in strong magnetic fields. Such systems may deliver advances in inertial confinement fusion [5–7], particle sources [8, 9], and atomic physics [10].

Concurrently, limitations on the strength of externally applied magnetic fields have also stimulated research into mechanisms of magnetic field generation and amplification by the laser-irradiated plasma itself. One such well-known method for generating quasi-static magnetic axial fields in an underdense plasma is the inverse Faraday (IF) effect enabled by circularly polarized light [11–14]. An underdense plasma with net orbital angular momentum (OAM) can also be created using intense twisted laser beams rather than circularly polarized lasers [15–17]. Magnetic field generation has also been observed in solid
Table 1. 3D PIC simulation parameters. $n_e = 1.8 \times 10^{21}$ cm$^{-3}$ is the critical density corresponding to the laser wavelength. The target is initialized without any preplasma. We confirmed that the resolution is adequate by conducting partial simulations with 4 particles per cell and 50 cells/μm. We have also conducted simulations with a smaller problem size with up to 20 particles per cell. These simulations further support the adequacy of our resolution.

| Laser parameters                                      |  |
|-------------------------------------------------------|---|
| Peak intensity                                        | $4.2 \times 10^{18}$ W cm$^{-2}$ |
| Normalized field amplitude                            |  |
| Right-hand circular polarization                      | $a_{\phi} = a_d = 1.0$ |
| Wavelength                                            | $\lambda = 0.8$ μm |
| Pulse duration (sin$^2$ electric field)               | $\tau = 600$ fs |
| Focal spot size (1/e electric field)                  | 3 μm |
| Location of the focal plane                           | $x = -2$ μm |
| Laser propagation direction                           | +x |

| Other parameters                                      |  |
|-------------------------------------------------------|---|
| Target thickness                                      | 0.4 μm |
| Electron density                                      | $n_e = 10 n_e$ |
| Seed magnetic field strength                          | $B_{0x} = 0.1$ kT |
| Transverse size of simulation box                     | 30 μm × 30 μm |
| Boundary conditions                                   | Open |
| Spatial resolution                                    | 40 cells/μm |
| Macroparticles per cell for each species              | 2 |

| Position and time reference                           |  |
|-------------------------------------------------------|---|
| Location of the front of the target                   | $x = 0$ |
| Time when peak of the laser is at $x = 0$             | $t = 0$ |

density targets, including the generation of azimuthal surface and bulk magnetic fields due to the propagation of relativistic electron beams [18–20], and is beneficial for hot electron transport and electron beam collimation [11, 13, 21–24]. Magnetic field amplification has also been observed in laser-produced plasma, for example in shocks [25], colliding flows [26], implosions [27–29], and with twisted light [30].

In this paper, we introduce a novel method for the amplification of a seed axial magnetic field in the interaction of a picosecond-scale, relativistic intensity laser pulse with a thin overdense target. This mechanism persists for both linear and circular polarization and can amplify a seed axial magnetic field of 20–100 T by a factor of 10. This paper is organized as follows: in section 2, we demonstrate magnetic field amplification for three choices of laser polarization, which distinguishes our amplification mechanism from the inverse Faraday effect. In section 3, we show that the magnetic field amplification is associated with the azimuthal current we observe in simulations, which in section 4 we identify as originating from the favorable orbital angular momentum electrons gain in the return current. In section 5, we explore the robustness of the amplification mechanism to the choice of simulation parameters.

2. Observation of axial magnetic field amplification

We conduct 3D simulations of a picosecond-scale, relativistic intensity laser pulse interacting with a thin target of fully ionized hydrogen with an imposed uniform axial seed magnetic field $B_{0x}$. We take the laser pulse to be either circularly or linearly polarized while keeping the same peak intensity. Simulations were carried out using the fully relativistic particle-in-cell code EPOCH [31]. Detailed parameters for the simulation are given in table 1. Energy was well conserved in these simulations.

We observe magnetic field amplification for three different laser polarization configurations, right-hand circularly polarized, left-hand circularly polarized, and linearly $y$-polarized. The relative strengths of the amplified magnetic field are given in table 2. In the right-hand circularly polarized case, the seed magnetic field is amplified from the initial $B_{0x} = 0.1$ kT to a peak amplitude of $B_x = 1.5$ kT over a volume of 4 μm$^3$ in approximately 300 fs. Figures 1(b) and (c) show the axial magnetic field in the $x$–$z$ and $y$–$z$ planes. The dashed lines in figure 1(b) correspond to the initial position of the target. In contrast with the dramatic axial magnetic field amplification shown in table 2, no substantial change is seen in the azimuthal magnetic field upon adding the axial seed magnetic field (see appendix D).

By probing different laser polarizations, we find that the observed magnetic field amplification clearly goes beyond the inverse Faraday (IF) effect. In the IF effect, the spin angular momentum of a circularly
Table 2. Amplified magnetic field strength for three laser polarizations with initial seed magnetic field $B_{x0} = 0.1 \, \text{kT}$ and for right-hand circular polarization with no seed magnetic field ($B_{x0} = 0$).

| Polarization                       | Peak magnetic field | Averaged magnetic field |
|------------------------------------|---------------------|-------------------------|
| Right-hand circular                | 1.5 kT              | 0.9 kT                  |
| Left-hand circular                 | 0.6 kT              | 0.38 kT                 |
| Linear                             | 0.8 kT              | 0.45 kT                 |
| No seed magnetic field ($B_{x0} = 0$) | 0.4 kT              | 0.2 kT                  |
| Right-hand circular                | 0.4 kT              | 0.2 kT                  |

Figure 1. Magnetic field amplification in a thin laser-irradiated target. (a) Schematic of axial magnetic field amplification. Initially, the laser pondromotively expels electrons creating charge separation. Later, electrons in the return current gain favorable OAM in the seed magnetic field, creating an azimuthal current which amplifies the field. (b) and (c) Amplified axial magnetic field for the right-hand circularly polarized case (table 1) at $t = 340 \, \text{fs}$ in the (b) $x$–$y$ plane ($z = 0 \, \mu\text{m}$, dashed lines denote the original target position), and (c) $y$–$z$ plane ($x = 0.5 \, \mu\text{m}$). The black contours denote $B_x = B_{x0}$. $B_x$ is temporally (16 fs) and spatially averaged in the plane with stencil size $0.25 \, \mu\text{m} \times 0.25 \, \mu\text{m}$. (d) Electron (left) and proton (right) density at $t = 340 \, \text{fs}$. The black contours denote the critical density $n_c$.

polarized laser [12] or the orbital angular momentum of a twisted laser [14] is transferred to electrons in the plasma, driving the generation of a magnetic field. In the absence of a seed magnetic field, the right-hand circularly polarized laser pulse generates a peak magnetic field of $B_x = 0.4 \, \text{kT}$. A left circularly polarized laser would generate $B_x = -0.4 \, \text{kT}$. In all cases with a seed magnetic field, we see peak magnetic fields in the same direction as the seed significantly above this level. We additionally observe substantial magnetic field amplification with linear polarization keeping the same peak intensity, albeit with broken symmetry in the $y$–$z$ plane (see appendix A), and with a left-hand circularly polarized laser pulse, where the IF effect generates a magnetic field in the opposite direction from the initial seed. As shown in table 2, the strongest magnetic field is generated with right-hand circular polarization, followed by linear polarization and then left-hand circular polarization. In light of these simulations, we find it likely that the high magnetic field amplitude we observe in the right-hand circularly polarized simulation represents a combination of magnetic field generation via the IF effect and magnetic field amplification via the novel mechanism we describe in this work. The remainder of this work will elucidate the magnetic field amplification process incorporating analysis of the right-hand circularly polarized case.

3. Relationship between magnetic field amplification and $j_\theta$

We find that magnetic field amplification is driven by electrons in the return current that arises after the peak of the laser pulse hits the target surface. When the laser pulse interacts with the target, it expels electrons from the laser spot, creating charge separation which later induces a return current of electrons relaxing back towards the laser interaction volume to neutralise the space charge. This return current obtains orbital angular momentum (OAM) from the seed axial magnetic field (figure 1(a)). The corresponding azimuthal current $j_\theta$ (figures 2(a) and (b)) drives the amplification of the seed magnetic field.
In our simulations, magnetic field amplification is driven by the generation of an azimuthal current by electrons which gain favorable OAM. First, in this section, we show that the observed magnetic field amplification can be explained quasi-statically, and is tied to the azimuthal current. Then, in section 4, we demonstrate how this azimuthal current can be generated by the electron return current and can persist over hundreds of femtoseconds.

We find that the azimuthal current \( j_\theta \) is responsible for the magnetic field amplification. We demonstrate this in two ways. First, we perform an order of magnitude estimate for the maximum axial magnetic field based on the \( j_\theta \) we observe in simulations. We see that a positive azimuthal current density \( j_\theta \) averaged over the cylindrical volume with radius and length of 1 \( \mu m \) and 0.5 \( \mu m \) averaged over all \( \theta \) and averaged in \( x \) over 0.4 \( \mu m < x < 0.9 \) \( \mu m \). Solid blue line: \( B_z \) measured from simulation. Dotted blue line: \( B_z \) calculated from Ampère’s law using the shown \( j_\theta \). The calculation was performed for \( r < 2.5 \) \( \mu m \) with \( B_z(r = 2.5 \mu m) = 0 \).

(c) Time evolution of axial magnetic field strength (blue curves, left axis), with reference laser amplitude \( a \) at the target surface (red line, right axis). Solid blue line: \( B_z \) averaged over the cylinder \( r < 1 \mu m, 0 < x < 1 \mu m \) (the same volume as table 2). Dotted blue line: average over a longer cylinder with the same radius \( r < 1 \mu m, 0 < x < 2 \mu m \), to capture the initial diamagnetic plasma response. The red dashed line denotes \( t = 0 \), which is when the peak of the laser pulse would pass the front target surface in the absence of the plasma. The simulation parameters are as given in table 1.

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We estimate \( R = \Delta x \sim 1 \mu m \), which predicts a maximum magnetic field strength of \( B_z^{max} \sim 1 \) kT, close to the peak magnetic field strength \( \langle B_z \rangle = 0.9 \) kT shown in figure 2(b), which was averaged over a cylindrical volume with radius and length of 1 \( \mu m \).

Second, we calculate the magnetic field profile as a function of radius based on Ampère’s law. The azimuthal current density and amplified magnetic fields we observe are fairly azimuthally symmetric (figures 1(c) and 2(a)). In the limit of perfect azimuthal symmetry, the axial magnetic field depends only on the azimuthal current density. We find good agreement between the axial magnetic field calculated from \( j_\theta \) and the radial magnetic field profile in simulations (figure 2(b)). Thus, we find that the magnetic field amplification is magnetostatic and driven by \( j_\theta \).

We can also use the azimuthal current density to calculate the OAM density produced in this simulation. The OAM density of electrons as a function of position \( r \) can be written as \( l_\theta = rm_v j_\theta \) where \( m_v \) is the electron mass, \( n_e \) is the number density, and \( v_\theta \) is the effective azimuthal velocity. The effective azimuthal velocity is related to the current density by \( v_\theta = -j_\theta / (|e| n_e) \), which allows us to calculate the OAM density as \( l_\theta = -r m_v j_\theta / |e| \). Using \( j_\theta \) as given above and the electron density from simulations, \( n_e \approx 10^{21} \) cm\(^{-3} \), we find that the OAM density is \( l_\theta = -0.02 \) kg m\(^{-1} \) s\(^{-1} \) at \( r = 1 \mu m \) and \( v_\theta \approx 0.02c \). In terms of the energy content, we find that the energy in the magnetic field (\( \epsilon_B = \int \mathbf{B}^2 / (2\mu_0) dV \approx 1 \) µJ) remains small compared to the kinetic energy of electrons around the amplifying area (\( \approx 50 \) µJ).
We now illustrate how the azimuthal current is produced by electrons which gain favorable OAM. When the laser pulse interacts with the target, the ponderomotive force expels electrons from the laser spot. Initially, electrons in the area in which the magnetic field is amplified (0.5 μm < x < 1 μm, |y| < 3 μm, |z| < 3 μm) at the sampling time \( t_i = 290 \, \text{fs} \). The simulation parameters are as given in table 1.

The averaged trajectory of 344 electrons selected at time \( t_i = 290 \, \text{fs} \) is given in figure 3. By analyzing the individual (e.g. the movie in appendix C) and averaged transverse position of these electrons \( r_i(t) \), the blue line in figure 3(a), we see that the majority of these electrons originate from radius \( r_e > 10 \, \mu m \), well outside the laser spot, and move inward to small radius during the falling edge of the laser pulse \( (t > 0) \). These observations identify these electrons as belonging to the return current.

We further see that inward motion of electrons can generate a negative average OAM \( \langle l_{ex} \rangle < 0 \) due to the axial seed magnetic field, which drives magnetic field amplification. The effect of the azimuthal magnetic field is already captured in the trajectory \( r(t) \) and does not exert an azimuthal force on the electrons (see appendix D). The average radial velocity, \( v_{er} = \frac{dr_e}{dt} \), of the test electrons is shown in the red line in figure 3(a). In order to illustrate how the sign of the OAM can be favorable for magnetic field amplification, we consider the azimuthal velocity this radial velocity produces in conjunction with the seed field \( \langle v_{\theta i} \rangle = \int (|v_{e}B_{0y}|/m_e dt \rangle \) and its corresponding OAM \( \langle l_{ex} \rangle \). Figure 3(b) shows that the negative inward radial velocity associated with the return current can drive \( v_{\theta i} < 0 \) and \( l_{ex} < 0 \), which represents a net positive azimuthal current \( \langle j_\theta > 0 \rangle \) for this group of electrons. The net rotation of electrons in the \( \gamma-z \) plane is also immediately visible in the trajectories of the electrons, e.g. the movie in appendix C.

We replicate the above analysis for groups of electrons chosen at different sampling times, \( (t_{s1} = 140 \, \text{fs}, t_{s2} = 190 \, \text{fs}, t_{s3} = 240 \, \text{fs}, t_{s4} = 290 \, \text{fs}) \) and find a negative azimuthal velocity \( \langle v_{\theta i} < 0 \rangle \), figure 4) can be maintained in the region of magnetic field amplification over a long time, consistent with our observation that a positive azimuthal current density at \( r = 0.5 \, \mu m \) is maintained in the simulation over hundreds of femtoseconds (red line in figure 4). Looking earlier in time, we see that the azimuthal current density is negative around the time when the peak of the laser pulse impacts the target \( (t = 0, \text{red dashed line in figure 4}) \). This is consistent with our observation that it is the return current that drives magnetic field amplification.

The net cycle-averaged inward force on an electron starting near rest (i.e. neglecting the magnetic force) encodes the competition between the ponderomotive force of the laser pushing the electron outward and the force due to charge separation pulling the electron inward. During the rising edge of the laser pulse \( (t < 0) \), electrons are pushed predominantly outward, and we see a small negative azimuthal current density, corresponding to the usual diamagnetic plasma response. We find that the diamagnetic response predominantly affects the axial magnetic field at the rear target surface, but does not have much effect on...

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**Figure 3.** Average trajectory of 344 electrons participating in the amplification process. (a) Averaged radial position of electrons \( r_i \) (blue line, left axis) and radial velocity \( v_{er} = \frac{dr_e}{dt} \) (red line, right axis). (b) Azimuthal velocity which would be expected based on interaction with the seed magnetic field \( v_{\theta i} = \int v_eB_{0y}/m_e dt \) (blue line, left axis) and corresponding OAM per electron \( l_{ex} = m_e v_{\theta i} \) (red line, right axis). Red dashed lines indicate \( t = 0 \), the time when the laser pulse is at its peak at \( x = 0 \).
Figure 4. Azimuthal velocity for different sets of sampled electrons and azimuthal current density in the plasma. (a) Azimuthal velocity $v_\theta$ (left axis) calculated as described in figure 3 for different sampling times ($t_{s1} = 140$ fs, $t_{s2} = 190$ fs, $t_{s3} = 240$ fs, $t_{s4} = 290$ fs), and azimuthal current density (red line, right axis) at location $x = 0.4 \mu m$ and averaged over the circle defined by $r = 0.5 \mu m$. The red dashed line denotes the time when the laser intensity at the front target surface ($x = 0$) is maximum. The simulation parameters are as given in table 1.

Figure 5. Net amplification of the axial magnetic field. (a) Scan over the seed magnetic field strength $B_{x0}$. (b) Scan over laser pulse duration. The parameters not scanned over are as given in table 1. Only the net amplification, $\langle B_x \rangle - B_{x0}$, is shown. $\langle B_x \rangle$ is averaged over the cylinder $r < 1 \mu m$, $0 < x < 1 \mu m$. Dashed lines indicate the growth rate of the magnetic field during the initial, roughly linear phase.

The magnetic field within the target [see the two blue curves in figure 2(c), which are averaged over different ranges in $x$]. During the subsequent falling edge of the laser pulse [to the right of the red dashed line in figure 2(c)], the force associated with charge separation can overcome the ponderomotive force of the laser and a net return current is produced. We find that the return current is able to drive magnetic field amplification during the falling edge of the laser pulse from $\langle B_x \rangle \approx B_{x0} = 0.1 kT$ to $\langle B_x \rangle = 0.9 kT$. The amplification occurs over approximately 250 fs, which is comparable to both half the pulse duration ($\tau_g/2 = 300$ fs) and the cyclotron period for an electron in the initial 0.1 kT seed field ($\tau_B = 340$ fs). The relationship between amplification and these timescales is discussed in more detail in section 5. After $t = 240$ fs, by which time the laser pulse has mostly reflected from the target (see figure 2(c)), the magnetic field begins to slowly decay. Over the course of the amplification, the charge separation (which peaks around $t = 0$) substantially decreases, but does not return to zero (see appendix B) and we find that both the positive azimuthal current density (e.g. figure 4) and the amplified magnetic field (figure 2(c)) can persist for hundreds of femtoseconds.

5. Parameter scan

We now investigate the robustness of the amplification mechanism described in sections 3 and 4 to the choice of simulation parameters. First, we scan the seed magnetic field strength $B_{x0}$ holding all the other parameters the same as in table 1 and investigate the balance of magnetic field amplification via our mechanism versus magnetic field generation via the IF effect. For a seed magnetic field strength below 20 T, the axial magnetic field becomes indistinguishable from the field produced in the absence of any seed magnetic field, as shown in figure 5(a). In other words, below 20 T no net magnetic field amplification is seen and the observed magnetic field can be attributed to magnetic field generation via the IF effect. In contrast, for a seed magnetic field strength above 20 T, we see that the seed magnetic field can be amplified by a factor of 10. The growth rate of the magnetic field during the initial linear rise increases as the seed field strength is increased (dashed lines in figure 5).

Second, we scan over the laser pulse duration $\tau_g$ at the original seed magnetic field strength ($B_{x0} = 0.1 kT$) to probe the impact of pulse duration on the magnetic field amplification. We see that the net amplification becomes weaker as the driving laser pulse duration becomes shorter and that the growth...
rate of the field during the initial linear rise is maximum for the 600 fs case, as shown in figure 5(b). We can obtain an average magnetic field as high as 1.1 kT for a pulse duration of 1 ps.

In conjunction with these first two parameter scans, we note that the period of electron Larmor precession in $B_{0y} = 0.1$ kT is around $\tau_B = 340$ fs and that for $B_{0y} = 20$ T, we have $\tau_B = 1700$ fs. This suggests that the amplification process requires the laser pulse duration to be sufficiently long that $\tau_f \gtrsim 0.5\tau_B$. This requirement on the pulse duration is consistent with our observation that the magnetic field amplification is driven by the return current. The return current rises and decays with the laser-induced charge separation, which decreases rapidly on the timescale of the pulse duration during the falling edge of the laser pulse (see appendix B). Thus, the primary return current timescale is the pulse duration. Over this time, electrons must undergo substantial momentum rotation to gain the orbital angular momentum needed to amplify the magnetic field, which sets the requirement that the pulse duration be at least on the order of the cyclotron period.

Third, we consider the robustness of the magnetic field amplification to the target thickness. The use of a thin target (0.4 $\mu$m thickness) in our original simulations maintains the feasibility of having the seed axial magnetic field penetrate the target on a reasonable time scale. However, for thicker targets the experimental time scales for generating the seed field and allowing it to penetrate into the target must be accounted for. For example, experiments have shown that the capacitor coil target can produce an axial magnetic field in excess of 0.1 kT with a sub-nanosecond rise and slow decay in excess of 10 nanoseconds [2, 3, 33], which corresponds to approximately a 10 $\mu$m penetration depth into a copper-like conductive material. Without performing fully self-consistent simulations of the magnetic field penetration into the target, we are therefore limited to studying $\mu$m-scale target thickness. For the sake of demonstrating that magnetic field amplification is feasible in a thicker target, we consider a second case with 2 $\mu$m thickness. We see that kilotesla-level magnetic field amplification is still present, albeit over a smaller spatial scale and with a somewhat reduced amplitude (see appendix A).

6. Summary and discussion

We demonstrate a novel mechanism for magnetic field amplification by a short pulse laser interacting with a thin overdense target capable of amplifying a 0.1 kT seed to 1.5 kT over a spatial extent of several cubic microns and persisting for hundreds of femtoseconds longer than the laser pulse duration. We find that magnetic field amplification is driven by the return current arising during the falling edge of the laser pulse. Electrons in the return current gain orbital angular momentum in the presence of the seed magnetic field, driving an azimuthal current with favorable sign for magnetic field amplification. This amplification process is robust to the choice of simulation parameters and occurs for both linear and circular polarization. For a right-hand circularly polarized pulse, we find that a seed magnetic field above 20 T delivers order-of-magnitude amplification from a 600 fs pulse and increasing the pulse duration from 150 fs to 1 ps increases the amplified magnetic field by a factor of 5.

We have neglected electron collisions when examining the generation of the magnetic field. In order to demonstrate that this simplification is appropriate, we estimate the electron collision frequency based on our simulation results. Electron tracking shows that the characteristic energy of the return-current electrons is $\epsilon \approx 15$ keV. The electron–electron collision frequency is $\nu_{ee} \approx 5.8 \times 10^{-6} e^{-3/2}n_e \lambda_{ee}$, where $n_e$ is the characteristic electron density in cm$^{-3}$, $\epsilon$ is the characteristic electron energy in eV, and $\lambda_{ee} \approx 10$ is the Coulomb logarithm. The magnetic field is generated over 400 fs, so the collisions can be considered as frequent only if $1/\nu_{ee} \ll 400$ fs. This corresponds to $n_e \gg 44n_c$. In our simulations, the initial target density is $n_t = 10n_c$, which indicates that it is reasonable to neglect electron collisions and explore the magnetic field generation using collisionless kinetic simulations. The target design considered in our work is experimentally feasible. Flat $\mu$m-thick hydrogen jets can be generated using cryogenic cooling [34]. The typical density is in the range of 20$n_c$.

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Appendix A. Magnetic field amplification by a linearly polarized laser and in a thick target

In this appendix A, we show magnetic field amplification driven by two additional configurations. First, we consider a thin target irradiated by a linearly y-polarized laser pulse. Second, we demonstrate amplification in a thick target driven by a right-hand circularly polarized pulse.

In the linear polarization case, we take $a_0 = 1.414$ to obtain the same peak intensity as the circularly polarized case. All other simulation parameters are the same as table 1. Figure A1 shows the amplified magnetic field we obtain at $t = 340$ fs. Compared to the right-hand circularly polarized case (figures 1(b) and (c)), the magnetic field is weaker in the linearly polarized case. We also see that there is asymmetry in the $y-z$ plane (figure A1(b)), which may be related to the polarization direction.

In the thick target case, we again consider a right-hand circularly polarized laser, which is now incident on a 2 $\mu$m thick target. All other simulation parameters are as given in table 1. The amplified magnetic field at $t = 340$ fs is shown in figure A2. Compared to the thin target case (figures 1(b) and (c)), the peak magnetic field strength is reduced with the thick target and the magnetic field is amplified over a smaller volume. However, the magnetic field is still amplified to the kilotesla level.

Appendix B. Charge density, electron current density, and ion current density

In this appendix B, we present additional properties of the charge density and current present in the thin target in the right-hand circularly polarized case with the simulation parameters given in table 1. The blue curve in figure B1(d) shows the charge density distribution $\rho(t)/|e|$ averaged over $0.2 \mu m < x < 0.4 \mu m$ and $1 \mu m < r < 4 \mu m$. During the rising edge of the laser pulse, the charge density increases, indicating a
net ponderomotive expulsion of electrons from the laser spot. After the peak of the laser pulse [vertical dashed line in figure B1(d)], the ponderomotive force on electrons in the target, $f_p(t) \propto -\partial a^2(t)/\partial r$, decreases, and the charge density $\rho(t)$ also decreased, consistent with a net inward return current. The charge separation initially decreases substantially over the falling half of the laser pulse ($0 < t \lesssim 300$ fs), then transitions to a much slower decay at later times. In this simulation, the timescale for the charge separation to decrease is roughly comparable to the cyclotron period ($\tau_B = 340$ fs) and is far longer than the plasma period ($\sim 0.8$ fs for a $0.1$ kT plasma).

Figure B1(b) shows the azimuthal and radial electron current densities (blue and red curves, respectively) as a function of radius at time $t = 340$ fs, which is after the laser pulse has been fully reflected by the target. Figure B1(c) similarly shows the ion current densities. These densities have been averaged over $0.2 \mu m < x < 0.4 \mu m$. We see that the ion current density is much smaller than the electron current density, which suggests the ion motion in the transverse ($y$–$z$) plane can be ignored. We also note that the total current density is positive, i.e. there is a net inward electron return current.

**Appendix C. Movie: trajectories of traced electrons**

Figure C1 shows the trajectory of one of the electrons we traced for the right-hand circularly polarized case with the simulation parameters given in table 1. Electrons oscillate longitudinally through the target ($x$-direction) while being pulled radially inward (left plot). The inward motion also corresponds to rotation in the transverse ($y$–$z$) plane (right plot). The movie further shows the trajectories of many traced electrons.

The movie shows clear cyclotron rotation of electrons, which is consistent with our expectations based on the electron cyclotron period ($\tau_B = 340$ fs for a 0.1 kT field). In this way, we can see directly that the return current electrons are magnetized. Ions, for which the cyclotron period is a factor of more than 1000 longer, are unmagnetized. We also see that as electrons approach small radius, the cyclotron motion can reverse direction. This is expected as electrons can overshoot the axis and begin to move radially outward (see, for example the dark blue group of electrons in figure 4). However, what we are showing in the movie is only a small subset of the electrons in the simulation and the overall azimuthal current density can still be positive (e.g. $j_\theta$ in figure 4).
Appendix D. Azimuthal magnetic field

In this appendix D, we show the azimuthal magnetic field $B_\theta$ generated by the hot electron current in the right-hand circularly polarized case both with and without the axial seed magnetic field. The simulation parameters are as given in table 1. Figure D1 shows a slice of the azimuthal magnetic field in the $y$–$z$ plane at $t = 340$ fs and $x = 0.4$ $\mu$m (the same time and roughly the same axial location as figure 1(c)). The strong azimuthal magnetic field likely plays a role in the return current dynamics. As return currents are drawn radially inward, they oscillate longitudinally through the target (shown in the trajectories in appendix C), which may be in part driven by $B_\theta$.

However, the azimuthal magnetic field is expected to have the same effect on the return current both with and without the axial seed field. As shown in figure D1, the difference in the azimuthal magnetic field with and without a 0.1 kT axial seed field is only 10%–20%. The effect of the azimuthal field on the return current behavior is encapsulated in the $r(t)$ we extract from simulations in analysis given in section 4. $B_\theta$ is aligned with the azimuthal current $j_\theta$ responsible for amplifying the axial magnetic field and does not affect $j_\theta$ directly. The biggest change in the return current dynamics is therefore associated with the axial magnetic field, which increases by a factor of 5 upon adding a 0.1 kT seed field.

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