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Magnetization of laser-produced plasma in a chiral hollow target

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Abstract

It is demonstrated that targets with a broken rotational symmetry may facilitate the generation of a strong axial (poloidal) magnetic field. An intense laser beam irradiating such a target creates strong electron currents carrying vorticity and producing strong spontaneous magnetic fields. Combined with laser–based acceleration schemes, such targets may be used for generation and guiding of magnetized, collimated particle or plasma beams.

1. Introduction

Laser acceleration of charged particles is a research domain promising interesting applications in fundamental science, medicine and technology [1–5]. Both electrons and ions may be accelerated with lasers, but so far the produced beams unfortunately have characteristics that limit the progress towards the promising envisioned future. The laser accelerated ion beams generated with solid targets have a broad energy and angular distributions, much larger than in conventional charged particle accelerators and many research groups are focused on improving these characteristics [2]. Laser acceleration of electrons is more mature and there has been a significant progress in improvement of laser accelerated electron beams quality in recent years [1, 6–9]. Commercial high power lasers allow to routinely achieve controllable production of electron beams with an energy spread between 1% and 10% and normalized transverse emittance in the sub-micron range. This progress relies entirely on physics-based optimization and control through manipulations of laser and target parameters [10–13]. Nevertheless there is still a progress to be made concerning of the electron beam characteristics. In this manuscript, we propose an additional way to control the laser–target interaction by the plasma magnetization on the relevant spatio-temporal scales. As we show, the considered magnetization scheme may be very effective, which makes the setup extremely interesting in the context of laboratory astrophysics and in many other applications.

Magnetic fields are widely used for collimation and guiding of particle beams, but the common methods of generation of strong magnetic fields with pulse power systems are hardly compatible with laser particle accelerators because of significantly different spatial and temporal scales [14, 15]. In contrast, the laser generated magnetic fields could be much better suited for beam manipulation, but the existing methods of creation of laser-generated magnetic fields are not sufficiently developed and the spatial structure of such fields is usually limited to the azimuthal (toroidal) magnetic field component [16–20]. It is desirable to develop more advanced methods of controllable magnetic field generation with lasers, which would be compatible with the laser generated accelerating fields.

A use of chiral structures allows to affect the polarization state of laser beams. It was suggested in [21] that laser beams reflected from a helically-shaped target may acquire an orbital angular momentum. The recoil orbital momentum is transferred to the target electrons and dissipated. A similar setup was considered in [22]. More generally, by an appropriate choice of the geometry of interaction, target shape and conductivity, one may extend the life time of electron currents and use them for generation of a quasi-static magnetic field. One example of such a structure was suggested in [23] where a laser beam is injected obliquely in a snail-shaped structure thus driving a vortical electron current in the laser propagation direction associated with a strong

\[ \text{References} \]
magnetic field oriented perpendicularly to the laser propagation axis. Here we propose another interaction geometry where the magnetic field is generated in the direction of laser beam propagation. It allows to significantly increase the interaction length and to combine together the processes of laser particle acceleration and magnetic field generation.

Two examples of hollow targets with a broken rotational symmetry (chiral hollow targets) are shown in figure 1. In contrast to toroidal magnetic fields generated by a straight electron current, here an azimuthal current having a structure resembling a charged $\theta$-pinch generates an axial magnetic field inside the target hole. Although the same laser pulse may do both, the magnetic field generation and particle acceleration, we focus here on the formation of the magnetized structure leaving the acceleration physics for a separate publication.

The paper is organized as follows: first, we present in section 2 an approach, which allows to simplify the problem by separating propagation (in section 2.1) and interaction inside a chiral target (in section 2.2). This simplifying approach allows to reduce the numerical problem from the fully three-dimensional geometry to a planar geometry with three velocity components (2D3V). The main elements of our scheme of the axial magnetic field generation are illustrated in section 2.3 with the particle-in-cell (PIC) simulations of a laser pulse interaction with a hollow chiral target. By using the targets shown in figure 1 we analyze typical current and magnetic field structures and their temporal evolution. Then in section 3 we discuss the properties of the vortical structure and the potential applications to the control of particle acceleration and guiding.

2. PIC simulations of a laser pulse interaction with a chiral hollow target

The problem of laser pulse interaction with an azimuthally asymmetric target is fully three dimensional and thus presents serious difficulties for both the theoretical analysis and numerical simulations. It can be simplified by
reduction to a two-dimensional planar geometry considering separately (i) propagation of a laser pulse in a thin fiber (waveguide) having symmetric transverse profile and (ii) laser pulse interaction with one slice of the target in the plane perpendicular to the laser propagation direction $z$.

The results of numerical PIC simulations, presented in this work, were performed with the code PICLS [24] in 2D3V geometry. On the step (i) an internal chirality of the fiber is ignored. The aim of this reduced calculation is to prove the possibility to propagate a laser beam inside a fiber and to heat its walls. The step (ii), being then justified, explains the mechanism of the quasi-static magnetic field generation inside fibers with chiral geometries, shown in figure 1.

### 2.1. Reduction of a 3D problem: 2D cut of a laser beam propagation in a fiber

For the simulations of the propagation of a laser beam inside a fiber, we considered a long simulation box $800 \times 60 \mu m$, see figures 2, 3, 4. The initial plasma density profile and the laser wave front are symmetric in the vertical direction; the laser pulse propagates from the left to the right. It has a Gaussian profile at the entrance plane ($z = 0$) with the full width at half maximum of $\sim 10 \mu m$ and the electric field polarized along $x$-axis. The laser wavelength is 800 nm, the intensity is $10^{20}$ W cm$^{-2}$, the pulse duration is 1 ps. To be able to enter a fiber it is tightly focused up to $\sim 2 \mu m$ at the point $z_f = 25 \mu m$ from the left boundary. Boundary conditions are absorbing for both the fields and the particles. The spatial resolution is $\approx 13.7 \text{ nm}$, so that the simulation box contains $73726 \times 5748$ cells, the temporal resolution is $\approx 0.037$ fs. Three cases are described below: free propagation of the laser pulse (figure 2), propagation in a straight fiber F1 (figure 3), and propagation in a conical fiber F2 (figure 4). The fibers with the walls of 1.2 $\mu m$ are made from Al with ion density $4.6 \times 10^{21}$ cm$^{-3}$ and electron density $6 \times 10^{22}$ cm$^{-3}$. This parameters provide the full laser pulse reflection inside the fibers.

The propagation of the laser pulse in vacuum is shown in figure 2. Three time moments, namely 1.2 ps, 2.7 ps, and 4.7 ps, correspond respectively to the propagation of the laser beam up to $\approx 360 \mu m$, $\approx 800 \mu m$, and $\approx 1400 \mu m$, the latter is already outside the simulation box. These three time moments are chosen also to describe the results for straight and conical fibers below. Already at 1.2 ps the laser pulse is defocused and expands to the side boundaries. Later the intensity is decreasing in the diverging pulse (see figure 2, the middle panel), and at the last shown time moment 4.7 ps the laser pulse has already left the simulation box.

In the case of the straight fiber F1 having the width of 16 $\mu m$, the laser pulse propagation is shown in figure 3. On panel A, the magnetic field $B_z$ component of the laser pulse and the electron density are shown at the time moments 1.2, 2.7, and 4.7 ps. Though the longitudinal position of the laser pulse reproduces the free propagation scenario shown in figure 2, the transverse divergence is suppressed by the fiber walls. As it follows from figure 3, panel (A), the laser beam heats the walls and the plasma expands inside. An important question is whether the expansion occurs during the laser pulse interaction, and what is a relation between the part of the hot electrons propagating along the fiber versus the electrons moving inside the fiber in the transverse direction. To resolve this question, electron spectra for internal electrons (i.e. the electrons, initially located in the half of the wall, which is closer to the fiber center) are shown in panel (B). According to the laser beam position, the electron spectra indicate energy increase both in longitudinal and transversal directions at the front of the laser beam. It follows from figure 3, panel (B), that the major part of the absorbed laser energy is converted in the transverse electron expansion, even in the interaction region. Only a small amount of absorbed energy is transferred into the longitudinal electron motion as it can be seen in figure 3. This is explained by a small number of the ponderomotively accelerated electrons compared to the total number of hot electrons, which are moving in all directions. Conversely, a small number of electrons accelerated to relativistic energies propagates predominantly in the laser propagation direction. Indeed, figure 3, panel (C), shows the spectra of relativistic electrons, with kinetic energy greater, than $m_e c^2$. There, the main peaks correspond to the laser-accelerated electrons in the forward directions, for the time moments 1.2 and 2.7 ps. The spectra for the latest time moment

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**Figure 2.** Results of PIC simulations for an $x$-polarized laser pulse, propagating along $z$-axis ($B_z$ field component) in an empty simulation box at 1.2, 2.7, 4.7 ps from up to down, respectively.
4.7 ps are not shown, because the number of the electrons with relativistic energies in the fiber is so small, that the averaging is not justified. The conclusion which follows from the presented results for the F1 fiber, is that the main part of the absorbed laser pulse energy is transferred to the electrons moving in the transverse direction and the expansion starts to develop already during the interaction process. The velocity of the electron expansion...
from the walls is about 10 μm s⁻¹, so the channel is quickly filled with hot electrons while it remains open for more than 5 ps.

To demonstrate the generality of this conclusion, another fiber geometry was considered for the same laser pulse, see figure 4. This F2 fiber is conical with the transverse size decreasing from 30 μm at the left to 20 μm at the right. Because the cone angle is very small, similar pictures are obtained for the magnetic field component $B_z$.
of the laser pulse, the electron density and for the electron spectra. However, in this case, the laser pulse is confined in the focused geometry, and the intensity increases at 2.7 ps compared to 1.2 ps, see figure 4, panel (A). Because of this intensity variation, the on-axis distribution of the average transverse electron momenta is flatter than in the straight F1 fiber, and the fraction of longitudinally accelerated electrons is greater, see figure 4, panel (B). At the same time, the average energy of relativistic electrons is lower in the case of the F2 fiber, see figure 4, panel (C), due to a larger fiber internal size and correspondingly a lower intensity of the pulse. The comparison between the results shown for the F1 fiber in figure 3 and the F2 fiber shown in figure 4 confirms that in both cases a considerable part of the hot electrons expands transversely inside the fiber.

The laser energy absorbed in the fibers is transmitted first to electrons. From the energy balance in the simulations (not shown) we may conclude, that the rate of laser energy extinction is about 1% per 100 μm. As it is shown in figures 3 and 4, the main electron energy corresponds to their transverse motion, while the most energetic electrons propagates along the fiber axis. The energy spectra are shown for the F1 fiber at 2.7 ps in figure 5. The high-energy part is composed of the longitudinally accelerated electrons, and may be fitted by the composition of the two Maxwell–Juttner distribution functions [25]. The less energetic one in this high-energy part is characterized by the ponderomotive scaling

\[ \epsilon_\perp \approx m_e c^2 \left( \sqrt{1 + a_0^2 / 2} - 1 \right), \]

where \( a_0 = \frac{E_0}{\omega m_e c} \), \( e \) and \( m_e \) are the elementary charge and the electron mass respectively, \( E_0 \) and \( \omega \) are laser electric field amplitude and frequency, \( c \) is the light velocity. The corresponding temperature is \( T_{a_0^{-1}} \approx \epsilon_\perp \approx 2.3 \text{ MeV} \). The other one in the high-energy part has a higher temperature \( T_{a_0^{-1}} \gtrsim 6 \text{ MeV} \). For the transverse electron motion, the temperature is much smaller, \( T_{a_0^{-1}} \sim 0.6 \text{ MeV} \), though these electrons contain the main part of the absorbed energy. The estimation of laser energy absorption rate in complicated geometries is a sophisticated problems, which is out of the scope of this paper, see, i.e. [25], and so for simplified estimates below we use either the results of our numerical simulations for the electrons remaining in the fiber (transversally accelerated) or the ponderomotive scaling for the main part of the directly accelerated electrons along the fiber.

The main goal of the present simulations is to show that along with the direct ponderomotive acceleration of the surface electrons a considerable laser pulse energy is absorbed on fiber walls and transformed to electron heating. By changing the fiber profile one can control the rate of laser absorption. Moreover, the laser absorption in the walls may be increased in the 3D geometry due to the symmetry violations. For example, by introducing a small angle between the laser pulse propagation direction and a fiber axis, providing a special polarization of the laser beam or introducing an additional torsion inside a fiber with a chiral internal shape, shown in figure 1. These and other possibilities may be considered in extended 3D3V PIC simulations or in experiments. Summarizing this subsection, we demonstrate, that it is possible to effectively heat electrons on the internal walls of a thin fiber, while only a limited part of them would be accelerated in the longitudinal direction by the light pressure.

In the forthcoming sections, we introduce a chirality for the internal walls of a fiber and demonstrate that an effective electron heating in such targets results in generation of strong spontaneous magnetic fields. The ablation process in the presence of the magnetic fields leads to formation of a magnetized structure in the target internal volume. Although some plasma appears inside the fiber during the laser pulse heating, the main evolution and relaxation to the quasi-stationary magnetized structure takes a longer time.
2.2. Reduction of a 3D problem: 2D target slice

The results of the simulations presented in the previous section (step (i)) show that the plasma parameters and laser field vary in the axial direction on lengths, which are much larger than the fiber internal size and there is no noticeable backscattered wave for the considered parameters. Therefore, the electron dynamics in the transverse plane can be studied separately by considering one slice of the fiber and prescribing the driver laser field. The particles in this approach are able to leave the simulation box in $x$ and $y$ directions, but remain inside if they are moving along $z$ axis. As the laser pulse absorption length is large compared to the hole size, the processes in consecutive slices are similar. This consideration provides the basis of the step (ii).

Let us consider an electromagnetic wave propagating in the free simulation box along $z$-axis with a frequency $\omega_0$ and a wave vector $k_0 = \omega_0/c$:

$$E(x, t) = E_0 \exp[ik_0 z - i\omega_0 t],$$

where $E_0$ is the laser electric field amplitude in the $(x, y)$ plane and $c$ is the light velocity.

According to Maxwell’s equations, the electric $E$ and magnetic $B$ fields of a plane wave polarized along $x$-axis are related as $B_y = E_x$ (in the CGS units). For a laser beam of a limited aperture $\Delta_x \gg 1/k_0$ and a finite duration $\Delta_t \gg 1/\omega_0$, these fields present a slow dependence in the perpendicular plane $x, y$ and in time defining the beam spatial and temporal profile. Moreover, other field components are also generated but their amplitudes are smaller by a factor $\epsilon = (k_0 \Delta_x)^{-1} \ll 1$. In the numerical examples presented below we consider a laser beam with the transverse width of the order of ten wavelengths. In this case the paraxial approximation is sufficient and other field components are at least $1/\epsilon$ times smaller. This is demonstrated with field maps in a plasma-free simulation in figure 6. Two main components, $E_x$ and $B_y$, dominate with the amplitudes $\sim 0.04$ in the relativistic units $m_e c \omega_0 / \epsilon$. The amplitudes of second-order dipole corrections $E_y$ and $B_x$ are an order of magnitude smaller than the main components. Next corrections of the quadrupole type, $E_z$ and $B_z$, are an order of magnitude smaller than the dipole components.

Figure 6. Electric and magnetic field components of a small amplitude laser beam propagating in the direction ‘into the picture’ in a plasma-free 2D simulation box of a size of $25 \times 25$ laser wavelengths. The beam has a width of 20 wavelengths at half maximum, it is linearly polarized and has $E_x$ and $B_y$ components in a free space. Other field components are generated according to the Maxwell’s equations due to the boundary conditions in the $x, y$ plane: $E_x$ and $B_y$ components are of a dipole type, $E_y$ and $B_x$ components are of a quadrupole type, both have much smaller amplitudes than $E_x$ and $B_y$. 
While propagating inside the hole, the laser beam edges interact with the target walls and create secondary (scattered) waves. Assuming that the laser field is not too much perturbed, we describe these secondary electromagnetic fields with the 2D Maxwell’s equations in the simulation plane and thus neglecting their variation in the laser propagation direction. This approach is certainly simplified as the propagation along z-axis is not accounted for, but its validity is demonstrated in the previous section. Therefore, these fields do not propagate along z-axis and remain in the simulation box even after the laser pulse ends. However, during the time of tens of picoseconds, the amplitude of secondary fields is evaluated correctly and the interaction of these fields with the plasma is accounted for. They induce charge separation electrostatic fields at the target edges and electron currents having a vortical component. The problems of plasma heating and axial magnetic field generation by these currents are the focus of our study.

2.3. Laser interaction with a hollow target

In numerical simulations we consider a fourth order super-Gaussian pulse profile in time and a Gaussian profile in space

$$ E_x(r, t) = E_0(z, t) \exp \left[ -\frac{(r - r_0)^2}{\Delta_r^2} - \frac{(t - t_0)^4}{\Delta_t^4} \right], $$

where $r_0$ is the center position of the target, $t_0$ is a time delay of the pulse, $\Delta_r$ and $\Delta_t$ are the beam width and the pulse duration. The laser wavelength $\lambda_0 = 0.8 \mu\text{m}$ corresponds to the period $\tau_0 = 2\pi/\omega_0 = 2.66$ fs and the beam radius $\Delta_r = 10 \lambda_0$. Three laser pulses are considered: a ‘long’ linearly or circularly polarized pulse of 100 periods, $\Delta_t = 100 \tau_0 = 266$ fs with an intensity $I_t = 10^{40}$ W cm$^{-2}$ and a ‘short’ linearly polarized pulse $\Delta_t = 10 \tau_0 = 26.6$ fs with an intensity $I_t = 2.5 \times 10^{33}$ W cm$^{-2}$.

Two considered targets, shown in figure 1 are made of a solid cylinder with figure cuts, so there are four (in T1) or three (in T2) claws, directed toward the center. The claws are disposed periodically and the shape of their edges is defined as

$$ r(\theta) = r_0 \left( 1 - \frac{\Delta r}{r_0} \frac{\theta}{\theta_0} \right), $$

where $r_0$, $\Delta r$ and $\theta_0$ are parameters. For the internal edge (closer to the center) of the T1 target, $r_0 = 9 \mu\text{m}$, $\Delta r = 7.6 \mu\text{m}$, and $\theta_0$ varies from 0 to $\pi/2$ periodically with a shift of $\pi/2$, while the second edge is straight in the radial direction. Obviously, such a shape brakes the rotational symmetry in the process of interaction. As we prove below, this leads to generation of surface currents of predominant direction, which produce a long-living magnetized plasma structures in the target hollow. For the target T2, the shape of the edges is defined as in equation (4), with the parameters $r_0 = 8.4 \mu\text{m}$ and $\Delta r = 5.2 \mu\text{m}$ for the internal edge, and $r_0 = 10.9 \mu\text{m}$ and $\Delta r = 10.0 \mu\text{m}$ for the external edge. In this case, $\theta_0$ varies from $2\pi/3$ to 0 periodically with a shift of $2\pi/3$. In the case of the target T2, we expect currents directed oppositely to those in the target T1, thus the generated internal magnetic field should also have the opposite direction. A quantitative characterization of the target chirality is discussed in section 3.

Both targets were modeled by homogeneously distributed aluminum ions with a charge $Z = 13$ and a density $n_i = 6 \times 10^{22}$ cm$^{-3}$, and electrons with a density $n_e = 7.8 \times 10^{23}$ cm$^{-3}$ corresponding to a plasma with the density 450 times the critical $n_c = 1.7 \times 10^{21}$ cm$^{-3}$. In all simulations there were 2 ions and 26 electrons per cell, with the initial electron temperature of 50 eV. This temperature is high enough to prevent the initial numerical heating, but much less than the average electron energy during and after interaction with the laser pulse. The simulation box contained 5748 \times 6144 cells, or approximately $25 \times 27 \mu\text{m}^2$. The resolution was 230 points per wavelength and per laser period. The collisions were included in the simulations based on the Takizuka and Abe model [27].

2.4. Examples of magnetic structures

Let us consider first the ‘long’ linearly polarized pulse. Its interaction with both targets is shown in figure 7. The laser field structure inside the target hole is defined by the boundary conditions, which are geometrically different for the targets T1 and T2. However, the pulse polarization does not substantially affect the heating process and magnetic field generation. This can be deduced from the comparison of figures 7 and 9, which present magnetic fields for both targets for the case of linear and circular polarization respectively. At the beginning of interaction, the electrons heated and ejected from the target walls propagate in the azimuthal direction in the target hole. These electron currents produce a fast growing axial magnetic field, which survives till the end of simulation, though its amplitude and spatial distribution vary with time. The direction of magnetic fields for the targets T1 and T2 is different (note the opposite color scales for T1 and T2). In the hole center for both considered targets at the end of the laser pulse the magnetic field amplitude is of the order of 1 in the
relativistic units. Figure 10 shows the distribution of the electron density and the electric current distribution in the target T2 providing an insight on the process of magnetic field generation. The currents are formed firstly near the target internal surface and then move towards the center. Initially, the radial motion dominates, see the time step 200 $\tau_0$ in figure 10. Then, the generated magnetic fields turn electrons in the azimuthal direction thus forming a rotating electron flow defined by the target chirality. Such an annular electron density structure around the central region is shown at the time step 500 $\tau_0$ in figure 10. Later in time, the electron diffusion across the magnetic field makes this structure smoother, as shown in the time step 800 $\tau_0$ in figure 10.

For comparison, figure 8 shows the same targets heated by a linearly polarized ‘short’ laser pulse with an intensity $I = 10^{20}$ W cm$^{-2}$ and time duration of 266 fs. The processes of electron heating and ejection are rather similar to the previous case, but due to higher electron energies, the resulting magnetic field is more structured and has a higher amplitude.

For the case of a circularly polarized long laser pulse the magnetic field distribution for both targets is shown in figure 9. A comparison with the previous case for a linearly polarized laser pulse in figure 7 shows, that the laser circular polarization enhances the magnetic field amplitude and makes the structure more stable. This can be explained by a more symmetric electron heating and also a more symmetric distribution of plasma flows in the hole. In the presented example we used the clock-wise polarization and we did not observe a noticeable difference in the distribution of the magnetic field and electron density by changing polarization direction.

Figure 7. Magnetic field $B_z$ in the target T1 (left column) and T2 (right column) at the time moments 110, 200, 500 and 800 $\tau_0$, for a linearly polarized laser pulse with an intensity $I = 10^{20}$ W cm$^{-2}$ and duration $\Delta t = 266$ fs. The time 140 $\tau_0$ corresponds to the laser pulse maximum.
The efficiency of the laser energy conversion into a quasi-static magnetic field may be analyzed by the energy balance, shown in figure 11 for the T2 target. The total energy deposited by the incident laser pulse in the simulation box grows approximately linearly with time, and it gradually decreases after the pulse ends. The most energetic particles and the radiation leave the simulation box in the transverse directions. Note, that the incident laser pulse energy is not shown in this figure. This explains fast oscillations of the total and electromagnetic energy in the case of a linear polarization. The absorbed laser energy is transferred first to electrons. Then, at the time scale of a hundred of femtoseconds, a part of this energy is transferred to ions due to the charge separation fields. The amount of the absorbed energy depends on the target length in z-direction, and can not be reproduced in the presented 2D3V simulations. Instead, we compare the energy of electrons with the total energy in the simulation box. According to figure 11, electrons and ions have \( \sim 15\% \) and \( \sim 60\% - 70\% \) of the maximum total energy in the simulation box for the shorter and longer pulses correspondingly. The electromagnetic energy during the simulation increases steadily with time, which is explained by the quasi-static magnetic field input. The residual part of the electromagnetic energy remaining in the simulation box after the laser pulse has gone then corresponds to the magnetic energy coupled to vortical plasma structures in the target hole. Estimation of the ratio of the magnetic field energy and the total energy in the simulation box (magnetization factor) gives the level of \( 10\% - 20\% \) depending on the interaction parameters. A more accurate estimate of the efficiency of magnetic field generation, defined as a ratio of the magnetic field energy to the total laser energy, requires a three-dimensional simulation and it depends on the target thickness in the laser.

Figure 8. Axial magnetic field \( B_z \) in the target T1 (left column) and T2 (right column) at time moments 55, 75, 300 and 750 \( \tau_0 \), for a linearly polarized laser pulse with an intensity \( 2.5 \times 10^{21} \, \text{W cm}^{-2} \) and duration 27 fs.
propagation direction. For optimal conditions, where a significant part of laser pulse energy is deposited in a target, it may reach a level of a few percent.

3. Discussion

In this section, we separate the stage of generation of the magnetic field (section 3.1), which is short and lasts as long as the laser pulse does, and the relaxation stage (section 3.2), where the magnetized structure slowly evolves. The first stage is also limited by the time of the internal plasma expansion. Because of this limitation, in our reduced approach, only short pulses are considered. For the case of our ‘long’ pulse of 266 fs (see figures 8, 9), the observed plasma expansion is still small.

3.1. Generation of the magnetic field

Propagating inside the hollow of a chiral target, the laser pulse heats the internal surface of the target. Electrons are gaining energy and ejected from the surface. As a secondary process, discharge currents of thermal electrons are excited, which tend to compensate the lack of electric charge on claw surfaces. The geometry of the claws defines the direction of the currents and in the case of a violated rotational symmetry leads to generation of long-living magnetized plasma structures in the target hole.
The direction of the current of hot electrons is defined by the laser pulse parameters and it has a broad angular distribution. Later, when they are strongly deviated by the generated magnetic field, a self-consistent structure is formed. Opposite to the current of hot electrons, discharging currents are formed at both surfaces of the claws, and directed toward the central part of the void. In figures 10 and 8 at the early time moments, 200 0 \( t \) and 75 0 \( t \), respectively, an increase of electron density and a growth of magnetic field are seen around the sharp edge of the claws, closer to the center. These electrons are accelerated and then deviated by the self-generated magnetic field, as it is seen in figure 10 at later time moment of 500 0 \( t \). The discharge currents have the same direction at both edges of the claws. It can be deduced from figure 7 at the time moment 200 0 \( t \) and figure 8 at the time moment 75 0 \( t \), where the values of the magnetic field, generated near the claws, have the opposite directions at the opposite sides of the claws (compare regions \( \alpha \) and \( \beta \) in figure 12). The spatial scale of the discharge currents is defined by the skin depth at the target claws, which is much smaller than their thickness, both in real solid dense metals and in our numerical approach (several nm). Because of a very sharp geometry of the discharge currents compared to the current of the hot electrons, the former ones are responsible for the generation of the magnetic field. This situation strongly differs from that considered in [23], where hot electrons were more collimated.

We introduce the parameter of chirality \( \chi_{ct} \), defined by a target geometry and characterizing the efficiency of magnetic field generation by the surface currents. These currents are directed along the surface of claws and thus are defined by the target geometry. The magnetic field in the center \( B_z(0) \) of a solenoid-type current distribution is defined by the Biot-Savart law (we use CGS units below)

\[
B_z(0) = \frac{2}{c} \int_{\zeta} \frac{J(\zeta) \times r(\zeta)}{r^2(\zeta)} \, d\zeta, \tag{5}
\]

where \( \zeta \) is a contour of claws, \( J(\zeta) = z^{-1} \int j(r) \, dr \, dz \) is the current on the claw contour per \( z \)-unit length, \( j(r) \) is the claw surface current, \( |j(r)| \sim \delta(r - \zeta) \) and \( r(\zeta) \) is the distance to a surface point. The quantitative calculation of surface currents is a complex problem, which deals with the multiple parameters of laser-target interaction. To classify targets by their geometrical properties, we assume a constant absolute value of \( J(\zeta) \), and normalize \( B_z(0) \) in equation (5) to the magnetic field created by the same current in the center of an ideal solenoid. Then the expression for the chirality parameter reads:

![Figure 10. Current distribution with a background of electron density (left column) and magnetic field (right column) in the target T2 at time moments 200, 500 and 800 \( t_0 \), for a circularly polarized laser pulse with an intensity \( 10^{20} \text{ W cm}^{-2} \) and duration 266 fs.](image-url)
Figure 11. Energy balance in the simulation box, extracted from the PIC simulations for the target T2 and the three types of the laser pulses. The total energy, the electromagnetic energy and the particle energy are calculated in the simulation box and are all the energies per unit length. These energies are scaled to the total energy of the laser pulses.

Figure 12. Electron density (left) and magnetic field (right) in the target T2 at the late time moment 2000 \( \tau_0 \), for a circularly polarized pump laser pulse with an intensity \( 10^{20} \) W cm\(^{-2} \) and duration 266 fs.
$B_z$ in the center (T2)

Figure 13. Magnetic field $B_z$ in the center of target T2, irradiated with a circularly polarized pump laser pulse with an intensity $10^{20} \text{ W cm}^{-2}$ and duration 266 fs, shown as a function of time.

\[ \chi_{igf} = \frac{B_z(0)}{4\pi f/c} = \frac{1}{2\pi} \int_{\zeta} \frac{\zeta \times r(\zeta)}{r^2(\zeta)} \, d\zeta, \]  

where $\zeta$ is a unitary vector along the surface, directed as the current $J(\zeta)$.

Applying this definition to the targets T1 and T2 one finds:

\[ \chi_{igf} = \frac{1}{2\pi} \sum_{\text{claws}} \int_{\theta_{\text{min}}}^{\theta_{\text{max}}} d\theta \left( \frac{\Delta r}{r(\theta) \theta_{\theta}} + 1 \right)^{-1/2}. \]  

After integration we obtain $\chi_{igf}^{T1} = 0.68$ for the target T1 and $\chi_{igf}^{T2} = -1.34$ for the target T2. As expected, the sign of $\chi_{igf}$ for these two targets is opposite, so that the magnetic field in the target center has the opposite direction. Also, from the calculated values of $\chi_{igf}^{T1}$ and $\chi_{igf}^{T2}$ it follows that the geometry of the target T2 is more suitable for the magnetic field generation. Indeed, as one can see in figures 7, 8 and 9, in the target T2 magnetic field structures are more pronounced.

Scaling for the characteristic magnetic field $B_z(0)$ and the current $J$ with the laser intensity may be derived from the energy and charge conservation considerations. The currents of cold electrons on the claw surfaces are excited due to neutralization of a positive charge created under the action of a laser pulse, when the heated discharge surface current per unit length is

\[ J_{\text{dis}} \approx \frac{\Delta r}{r(\theta) \theta_{\theta}} + 1 \]  

where $z_0$ is the laser propagation length along z-axis in the target hole. This current, according to equation (5), produces the magnetic field in the center of a target with the chirality $\chi_{igf}$

\[ B_z^* \sim \chi_{igf} \frac{4\pi f^n}{c} \sim \chi_{igf} \frac{4\pi}{c} \frac{P_{\text{laser}}}{\epsilon_e z_0} \frac{\epsilon_e}{\epsilon_e}. \]  

The mean electron energy may be extracted from the PIC simulations in the chiral targets, as we did before for the fiber F1. The corresponding electron spectra and their fits with Maxwell–Juttner distribution functions at 0.5 ps are presented in figure 5 for the T1 target. The electron distribution is approximately isotropic with the temperature ($T_{\text{rel}} \sim 0.8$–0.9 MeV) which differs two times from the ponderomotive scaling (1), $\epsilon_e \sim 2$ MeV.

For the parameters of figure 7 the mean electron energy is $\epsilon_e \sim 0.8$ MeV and the characteristic field is $B_z^* \sim 1.5 \times 10^9 \chi_{igf} \eta_e /z_0$ Gauss, where $z_0$ is in cm. With the calculated chirality and the values of the magnetic fields, observed in the simulations, the coefficient $\eta_e /z_0$ appears to be $\sim 1$ cm$^{-1}$. This value agrees with the absorption rate observed in the fiber propagation in section 2.1.

A special case presents the interaction of a circularly polarized beam with the T1 target shown in figure 9. In the central region inside the circle limited by the claw’s edges the field has a negative direction, as expected. However, near the central point the field changes the direction, in contrast to the direction dictated by the chirality. This is explained by a smaller heating rate in the case of circular polarization comparing to the linear one (see figure 7). At the beginning of interaction, the electron cloud is located closer to the claws. Discharge currents are sufficiently strong and create compact magnetized structures confined between the neighbor claws. These structures overlap by their edges thus inducing an oppositely directed magnetic field in the fiber center. This however, does not change the direction of the magnetic flux in the central region, which is defined by the target chirality. Our estimates characterize the average magnetic flux. The detailed structure could be more complicated.
3.2. Magnetized plasma structure in a target hole

Let us consider the magnetized structures at later times, when the pulse is gone. This stage corresponds to $t \gtrsim 0.5$ ps for 266 fs pulses, and $t \gtrsim 0.05$ ps for the 27 fs pulse. The magnetic field amplitude changes slowly on hydrodynamic time scale. During this evolution, it reaches the maximum value due to compression by the ablated plasma, and then gradually decreases. In figure 13 the amplitude of the magnetic field in the center of the target T2 is shown as a function of time for the 266 fs circularly polarized laser pulse.

For the case of a 266 ps circularly polarized laser pulse, interacting with the target T2, at the late time moment $t = 2000\tau_0 \approx 5.3$ ps the electron density and the magnetic field distributions are shown in figure 12. They correspond to a quasi-stationary magnetized structure. The profiles averaged over the polar angle for the $z$-component of the magnetic field $B_z(r)$, the electron density $n_e(r)$, the $\theta$-component of the current density $j_\theta$, the charge density $Zn_i(r) - n_e(r)$, and the radial component electric field $E_r(r)$ are shown in figure 14. The charge density and the electric field are relatively low compared to $n_e(r)$ and $B_z(r)$. Figure 12 shows the magnetic field structure very similar to a $\theta$-pinch, which appears to be almost neutral with small radial electric field fluctuations. This special feature is related to non-relativistic velocities of electrons. The average positive current, coupled to the magnetic field, is generated due to the non-zero chirality of the target.

For the central magnetic structure formed inside a chiral target at late times, like the one shown in figure 8 for the T2 target at the time of 800 $\tau_0$, the plasma inside the hole is relatively cold and can be described by the two-fluid hydrodynamic equations:

$$\partial_t n_{e,i} + \nabla (n_{e,i} v_{e,i}) = 0,$$

$$\partial_t p_{e,i} + (v_{e,i} \nabla) p_{e,i} = q_{e,i} E + \frac{q_{e,i}}{c} \nabla \times B,$$

where $p_{e,i} = \gamma m_{e,i} v_{e,i}$ is the particle momentum, $\gamma = 1/\sqrt{1 - v^2/c^2}$ is the relativistic factor and $q_{e,i} = -e$, $Ze$ are the electron and ion charges. The electric and magnetic fields are described by Maxwell’s equations with the current $j = -en_e v_e + eZn_i v_i$ and the charge density $\rho = eZn_i - en_e$.

Let us consider stationary axially symmetric solutions corresponding to the electron vortex with immobile ions, the radial electric field $E_r(r)$, axial magnetic field $B_z(r)$ and the azimuthal electron velocity $v_\theta(r)$. The stationary equations read:

$$m_e \gamma v_\theta^2 / r = eE_r + ev_\theta B_z / c,$$  

(10a)

$$d_r B_z = 4\pi e n_e v_\theta / c,$$  

(10b)

$$r^{-1} d_r (E_r r) = 4\pi e (Zn_i - n_e).$$  

(10c)

These equations describe a collisionless charged $\theta$-pinch. The common nature between charged and neutral (collisional) $\theta$-pinch is the plasma equilibrium at a given $r$, see equation (10a). In a conventional $\theta$-pinch the expanding force is the plasma pressure, while in our situation this role is played by the electric field between the electrons and immobile ions. As in the conventional $\theta$-pinch, there is a freedom to chose two functions due to the fact that the number of equations in (10) is less than the number of variables. A particular solution depends on the scenario of production, it may be found as an initial-value time-dependent problem.
We consider here model solutions by using the profiles, extracted from the PIC simulations for late times when the magnetized plasma structure has zero total charge \( Q_{\text{tot}} = 0 \), where

\[
Q_{\text{tot}} = 2e \int_0^\infty r \, dr (Zn_i - n_e).
\]

We assume then a profile of the magnetic field, similar to that in figure 14

\[
B_z(r) = B_0 \exp \left[ -r^2/r_0^2 \right],
\]

and the velocity profile, which evolves on the same spatial scale

\[
v(r) = -v_0 \frac{r}{r_0} \exp \left[ -r/r_0 \right].
\]

Then, according to equation (10b) the electron density is

\[
n_e(r) = \frac{2B_0}{v_0 r_0} \exp \left[ r/r_0 - r^2/r_0^2 \right],
\]

the electric field and the ion density are then defined from equations (10a) and (10c) (the expressions are too cumbersome to write them explicitly). The condition \( Q_{\text{tot}} = 0 \) is fulfilled as long as \( E(r) r_0^{-\infty} = 0 \).

The parameters in equation (12), \( B_0, r_0 \) and \( v_0 \) are obtained from the profiles, shown in figure 14. The characteristic scale length \( r_0 \) is of the order of 3 \( \lambda_0 \), and the amplitude \( B_0 \approx 1.7 \) in the relativistic units. The adjusted profiles are shown in figure 15, they reproduce qualitatively the general features of the numerical results from figure 14.

The model presents in general a relativistic magnetized structure. To reach a relativistic regime, the charge separation should be sufficient to produce a high electric field, which can bind relativistic electrons. For this situation, the profiles are shown in figure 16, where we use another set of parameters \( B_0 = 2 \) relativistic units, \( r_0 = 3 \lambda_0 \) and \( v_0 = 0.9c \). As it follows from figure 16, for the higher electron velocity the electron density is smaller. This is explained by equation (10b), which binds the current and the magnetic field. As a result of a higher velocity the centrifugal electrical force and the charge separation become greater.

4. Conclusions and perspectives

We presented a new scheme for the production of strong quasi-stationary magnetic field, bound in plasma structures, in interaction of intense laser pulses with chiral targets. The scheme has some common features with the previous studies dealing with the laser generation of quasi-static magnetic fields in a specific interaction geometry [23, 28]. Here, we make a step in this direction toward more complex three-dimensional target geometries, which possess high symmetry and are more suitable for particle acceleration and astrophysical applications. Another advantage of the presented approach is overcoming the restriction of the snail-shaped targets, caused by the relativistic transparency [29].
Magnetic fields in plasma are widely studied in the context of particle acceleration mechanisms [30, 31]. The magnetic field structure produced in the laser pulse interaction with a chiral target may be used for the guiding of a charged particle beam. It would also be possible to accelerate charged particles directly from the magnetized plasma using a secondary laser pulse [32]. Another extension of the considered approach may be developed for recently proposed MeV-photon generation with ultrarelativistic laser pulses, interacting with an overdense structured targets [20]. The azimuthal (toroidal) magnetic fields, allowing the high-energy photon generation in [20], with a chirality of the targets, may be enriched by the axial component, leading to an additional control of the electron beam and generated high-energy photon flux properties. As we showed, by choosing the 3D target structure one may control a spatial distribution of a magnetic field and use it to obtain certain conditions for particle or plasma beam production.

In the considered approach, the plasma density may be varied by adjusting the size and material of a chiral target and introducing a secondary laser pulse with a controllable delay time. In the case of low density magnetized structures, electrons could be accelerated in the bubble regime [33] and ions could be accelerated in collisionless shocks [34, 35]. In the case of higher density plasmas it would be possible to accelerate ions through charge-separation fields due to hot electrons exiting the target, like in the case of target normal sheath acceleration (TNSA) [36], or directly by the laser field, like in the case of radiation pressure acceleration (RPA) [37, 38]. In all cases, the fact that the plasma where particles are accelerated is magnetized can lead to the generation of higher quality electron beams [39] and to more energetic and collimated ion beams thanks to a better confinement of the hot electrons and higher accelerating fields. Laser ion acceleration through TNSA or RPA in such geometries would also offer new possibilities to study magnetized plasma collisions, which are of great interest for laboratory astrophysics.

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