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Inverse Faraday effect driven by radiation friction

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Abstract

A collective, macroscopic signature to detect radiation friction in laser–plasma experiments is proposed. In the interaction of superintense circularly polarized laser pulses with high density targets, the effective dissipation due to radiative losses allows the absorption of electromagnetic angular momentum, which in turn leads to the generation of a quasistatic axial magnetic field. This peculiar ‘inverse Faraday effect’ is investigated by analytical modeling and three-dimensional simulations, showing that multi-gigagauss magnetic fields may be generated at laser intensities $>10^{23}$ W cm$^{-2}$.

1. Introduction

The development of ultrashort pulse lasers with petawatt power has opened new perspectives for the study of high field physics and ultra-relativistic plasmas [1, 2]. In this context, the longstanding problem of radiation friction (RF) or radiation reaction has attracted new interest. RF arises from the back-action on the electron of the electromagnetic (EM) field generated by the electron itself and plays a dominant role in the dynamics of ultra-relativistic electrons in strong fields. A considerable amount of work has been devoted both to revisiting the RF theory [3, 4] and to its implementation in laser–plasma simulations [5–9], as well as to the study of radiation-dominated plasmas in high energy astrophysics, see e.g. [10–12].

While RF is still an open matter both for classical and quantum electrodynamics [2], RF models have not been discriminated experimentally yet. This circumstance led to several proposals of devoted experiments providing clear signatures of RF, e.g. in nonlinear Thomson scattering [13–18], Compton scattering [19], modification of Raman spectra [20], electron acceleration in vacuum [21–23], radiative trapping [24–26] or $\gamma$-ray emission from plasma targets [27, 28]. Most of these studies are based on single particle effects, and RF signatures are found in modifications of observables such as emission patterns and spectra when RF is included in the modeling. Detecting such modifications may require substantial improvements in reducing typical uncertainties in laser–plasma experiments. At very high intensities, RF losses may affect the collective dynamics, e.g. by modifying the spectra of accelerated ions in the radiation pressure dominated regime [29, 30, and references therein] or the dynamics of magnetic field generation by the filamentation instability in laser-generated colliding pair plasmas [31]. However, also in this case the modifications are quantitative, rather than qualitative, and relatively modest so that it may be difficult to discriminate RF effects.

Instead, in this paper we identify a collective, macroscopic effect induced by RF, namely the generation of multi-gigagauss, quasi-steady, axial magnetic fields in the interaction of a circularly polarized (CP) laser pulse with a dense plasma. This is a peculiar form of the inverse Faraday effect (IFE) [32–35] and may be more accessible experimentally than single-particle effects. In fact, the IFE has been previously studied in different regimes of laser–plasma interactions [36–42, and references therein]. By using three-dimensional (3D) particle-in-cell (PIC) simulations, we find that at laser intensities foreseeable with next generation facilities producing multi-petawatt [43] or even exawatt pulses [44, 45], the magnetic field created by the RF-driven IFE in dense...
plasma targets reaches multi-gigagauss values with a direction dependent on the laser polarization, which confirms its origin from the ‘photon spin’. The magnetic field is slowly varying on times longer than the pulse duration and may be detected via optical polarimetry techniques [46–50]. This would provide an unambiguous signature of the dominance of RF effects, since the axial magnetic field disappears in the absence of RF. The effect might also be exploited to create strongly magnetized laboratory plasmas in so far unexplored regimes (see e.g. [51]).

2. Role and modeling of radiative losses

The IFE is due to absorption of EM angular momentum [6], which in general is not proportional to energy absorption. As an example of direct relevance to the present work, let us consider a mirror boosted by the radiation pressure of a CP (with positive helicity, for definiteness) laser pulse. From a quantum point of view, the laser pulse of frequency $\omega$ propagating along $\hat{x}$ corresponds to $N$ incident photons with total energy $N\hbar \omega$ and angular momentum $N\hbar \hat{\mathbf{x}}$. If the mirror is perfect, $N$ is conserved in any frame. If the mirror moves along $\hat{x}$, the reflected photons are red-shifted leading to EM energy conversion into mechanical energy (up to 100% if the mirror velocity $\sim c$) but there is no spin flip for the reflected photons, hence no absorption of angular momentum. However, if the electrons in the mirror emit high-frequency photons, a greater number of incident low-frequency photons must be absorbed with their angular momentum. From a classical point of view, absorption of angular momentum requires some dissipation mechanism [42] which, in our example, implies a non-vanishing absorption in the rest frame of the mirror.

In the case here investigated, effective dissipation is provided by the RF force which makes the electron dynamics consistent with the radiative losses. In order to demonstrate IFE induced by RF, we consider a regime of ultra-high laser intensity $I_L > 10^{23}$ W cm$^{-2}$ and thick plasma targets (i.e. with thickness much greater than the evanescence length of the laser field) where the radiative energy loss is a significant fraction of the laser energy as shown by simulations with RF included [29, 30, 54–56]. We use a simple model to account for such losses and provide a scaling law with the laser intensity. The power radiated by an electron moving with velocity $v_e$ along the propagation axis of a CP pulse of amplitude $E_L = (m_e c/\omega) a_0 \equiv B_0 a_0$ (with $\omega$ the laser frequency) is

$$P_{\text{rad}} = \frac{2e^2 \omega^2 \gamma^2 a_0^2}{3c} \left(1 - \frac{v_e}{c}\right)^2. \quad (1)$$

The classical RF force on an electron is defined in order that the work done per unit time equals $P_{\text{rad}}$. In our simulations the Landau–Lifshitz (LL) expression for the RF force has been used [57]. Consistently with equation (1) the LL force in a plane wave vanishes for $v_e = c$, and has a maximum for $v_e = -c$. The spectrum of the emitted radiation peaks at frequencies $\omega_{\text{rad}} \simeq \gamma^2 \omega$, with $\gamma$ the relativistic factor of electrons which can be estimated as $\gamma \simeq a_0 \omega$. At such frequencies the radiation from the plasma is incoherent (see also the discussion in section 3), thus the total radiated power by $N$ comoving electrons will be $NP_{\text{rad}}$. For thin targets accelerated by the CP laser pulse (‘light sail’ regime), all electrons move with the foil at $v_e \simeq c$, and there is no high-frequency oscillation driven by the $\mathbf{v} \times \mathbf{B}$ force. Thus the radiation is strongly suppressed by the factor $(1 - v_e/c)^2 \ll 1$, as observed in simulations [7, 58]. In contrast, RF losses become much larger for thick targets [29, 30, 54] (‘hole boring’ regime) because the acceleration of the plasma surface has a pulsed nature [55, 59, 60] with a dense bunch of electrons being periodically dragged towards the incident laser pulse, i.e. in a counterpropagating configuration ($v_e < 0$).

In order to estimate the number of radiating electrons per unit surface we consider the dynamic picture of hole boring [39, 61]. As illustrated in figure 1, at the surface of the plasma the radiation pressure generates a positively charged layer of electron depletion (of thickness $d$) and a related pile-up of electrons in the skin layer (of thickness $\ell_\gamma$), i.e. the evanescent laser field region. Ions are accelerated in the skin layer leaving it at a time $\tau$ at which an ion bunch neutralized by accompanying electrons is formed. At this instant, the equilibrium between ponderomotive and electrostatic forces is lost and the excess electrons in the skin layer will quickly return back towards the charge depletion region. The number per unit surface of returning electrons is $N_\gamma = (n_{\text{po}} - n_0) \ell_\gamma$ where $n_{\text{po}}$ is the electron density in the skin layer at the beginning of the acceleration stage. Using the model of [39, 61], $N_\gamma$ may be estimated from the balance of electrostatic and radiation pressures: $eE_{\text{d}} n_{\text{po}} \ell_\gamma/2 = 2H/\gamma$, where $E_{\text{d}} = 4\pi e \varepsilon_0 d$ is the peak field in the depletion region, and $n_{\text{po}} \ell_\gamma = n_0(d + \ell_\gamma)$ because of charge conservation. Eliminating $d$ from these equations yields for the density compression ratio in the skin layer

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6 Here we consider only the absorption of intrinsic angular momentum or photon ‘spin’. For studies on orbital angular momentum absorption and IFE in laser–plasma interaction see, e.g., [52, 53].
where \( n_i = m_e \omega^2 / (4 \pi e^2) \). In our conditions the field laser is evanescent, thus \( (\omega / \omega_c) > 1 \), and we consider a range of parameters such that \( a_0 n_i / n_0 > 1 \). We thus obtain \( n_{\text{ef}} / n_0 \approx 2 a_0 (n_i / n_0)^{1/2} \gg 1 \), in agreement with a detailed theory of nonlinear wave propagation (see e.g. [62]) and numerical simulations. Thus, by writing \( I_\ell = m_e \omega^2 a_i^2 / (4 \pi c) \) (where \( r_c = e^2 / m_e c^2 \)) we obtain \( N_\ell \approx a_0 / r_c \lambda \) (where \( \lambda = 2 \pi c / \omega \) is the laser wavelength), independently on the initial density\(^7\).

The total radiated intensity is \( I_{\text{rad}} = P_{\text{rad}} N_\ell \). In order to compare the radiated energy with the laser pulse energy we take into account that the radiation is emitted as bursts corresponding to the periodic return of electrons towards the laser, i.e. for a fraction \( f_\ell \approx \tau_c / (\tau_e + \tau_i) \) of the interaction stage where \( \tau_c \) is the time interval during which the electrons move backwards. Analysis of laser piston oscillations in [53] suggests that \( \tau_c \approx \tau_i \) so we take \( f_\ell \approx 1/2 \) for our rough estimate. Since \( (1 - \nu_i / \nu_e) \approx 1 \) for returning electrons, we obtain for the fraction of radiated energy to the laser pulse energy

\[
\eta_{\text{rad}} \approx \frac{4 \pi r_c}{3 \lambda} a_0 \gamma^2.
\]

If the energy of electrons is mainly due to the motion in the laser field, then \( \gamma \approx (1 + a_0^2)^{1/2} \approx a_0 \) for \( a_0 \gg 1 \) and \( \eta_{\text{rad}} \propto a_0^3 \). For \( \lambda = 0.8 \mu \text{m}, \eta_{\text{rad}} \approx 1 \) for \( a_0 \approx 400 \), corresponding to \( I_\ell \approx 7 \times 10^{23} \text{ W cm}^{-2} \). This order-of-magnitude estimate implies that for such intensities a significant part of the laser energy is lost as radiation, strongly affecting the interaction dynamics. A more precise estimate would require to account both for the energy depletion of the laser and for the trajectory modification of the electrons due to the RF force.

### 3. Simulation results

A 3D approach is essential to model the phenomena of angular momentum absorption and magnetic field generation, thus we rely on massively parallel PIC simulations in which RF is implemented following the approach described in [7], based on the LL equation (see [9] for a benchmark with other approaches). We remark that the inclusion of the radiation loss as a dissipative process via the RF force requires the following assumptions: (i) the dominant frequencies in the escaping radiation are much higher than the highest frequency that can be resolved on the numerical grid, (ii) the radiation at such frequencies is incoherent, (iii) the plasma is transparent to such frequencies. Since, as also stated above, for radiation in the field of a plane CP wave of ultrarelativistic intensity the radiation spectrum peaks at frequencies of the order of \( \omega_{\text{rad}} \approx a_0^3 \omega \), all the above assumptions are well-satisfied in our conditions.

The laser pulse is initialized in a way that at the waist plane \( x = 0 \) (coincident with the target boundary) the normalized amplitude of the vector potential \( \mathbf{a} = e \mathbf{A} / m_e c^2 \) would be

\(^7\)The scaling \( N_\ell \propto a_0 / \lambda \) is a consequence of the balance between electrostatic and radiation pressures before electrons return towards the laser pulse. In fact, since the electron density of the compressed skin layer \( n_{\text{el}} \gg n_0 \) (the initial electron density), the excess number \( N_\ell \) of electrons is almost proportional to the charge-displacement field \( E_a \). Since the electrostatic pressure \( P_{\text{el}} \propto E_a^2 \), by posing \( P_{\text{el}} = P_{\text{rad}} \propto k_i / \epsilon \propto a_i^2 / \lambda^2 \) we obtain \( N_\ell \propto E_a \propto a_0 / \lambda \).

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**Figure 1.** Cartoon showing the electron dynamics during the ‘hole boring’ stage. Frame (a) shows the approximate profiles of the densities of ions \( n_i \) and electrons \( n_e \) and of the electrostatic field \( E_d \) at the early time \( t = 0 \) when ions have not moved yet and electrons from the depletion region \( 0 < x < d \) pile up in the skin layer \( d < x < d + \ell_s \); the number of excess electrons \( N_\ell \) is proportional to the shaded area. Frame (b) corresponds to the time \( t \approx \tau_i \) when the ions have reached the \( x = d + \ell_s \) position and formed a quasi-neutral bunch [59]; the excess electrons return towards the depletion region and radiate predominantly at the angle \( 3\pi / 4 \) with respect to the propagation direction.

\[
n_{\text{el}} = \frac{1}{2} + \left( \frac{1}{4} + 4 \left( \frac{n_e c}{n_0 \omega \ell_s a_0} \right)^2 \right)^{1/2},
\]
where \( r = (y^2 + z^2)^{1/2} \). Both radial profiles with \( n = 2 \) (Gaussian, G) and \( n = 4 \) (super-Gaussian, SG) have been used in the simulations. For all the results shown below, we take \( \eta_1 = 3 \lambda \) and radius \( R_0 = 3.8 \lambda \). The plus and minus sign in the expression for \( \Phi \) correspond to positive and negative helicity, respectively. The pulse energy is given by: \( E = \frac{c n_a}{h \omega} a_0 \), where \( A = \Gamma(1/4)2^{-17/2} \approx 0.19 \) and \( A = \pi^{1/2}\Gamma(1/4)2^{-19/2} \approx 0.24 \) for the G and SG pulse cases, respectively. The target is a plasma of thickness 10 \( \lambda \) and electron density \( n_e = 90n_c \) and charge-to-mass ratio for ions \( Z/A = 1/2 \). The range of laser amplitudes investigated in the simulation is \( a_0 = 200-600 \). Assuming \( \lambda = 0.8 \mu m \), the density \( n_0 = 1.55 \times 10^{17} \) cm\(^{-3} \), the pulse duration (full-width-half-maximum of the intensity profile) is 14.6 fs and the range for the peak laser intensity \( I_e = mc^2n_0a_0^2 \) is \( (1.9-16.7) \times 10^{19} \) W cm\(^{-2} \) corresponding to a pulse energy \( U_e \approx (0.38-3.4) \) kJ for the G pulse and \( U_e \approx (0.48-4.3) \) kJ for the SG pulse. The numerical box had a 30 \( \times \) 25 \( \times \) 25 \( \lambda \) size, with 40 grid cells per \( \lambda \) and 64 particles per cell for each species. The simulations were performed on 480 cores of the JURECA supercomputer at Forschungszentrum Jülich.

Figures 2(a)-(c) show the magnetic field \( B_x \) (normalized to \( B_0 = 1.34 \times 10^8 \) G) along the propagation direction at time \( t = 27\lambda/c \) for a simulation (a) where RF is not included and for two simulations (b) and (c) including RF and having positive and negative helicity, respectively; the laser profile was super-Gaussian and \( a_0 = 600 \) in all the three simulations. Only with RF included an axial magnetic field of maximum amplitude \( B_{max} \approx 22B_0 = 2.9 \times 10^8 \) G, extending over several microns and a polarity inverting with the pulse helicity is generated. The comparison of figures 2(d) and (e) shows that \( B_x \) has similar values and extension for a Gaussian pulse. The field is slowly varying over more than a ten laser cycles (\( \sim 30 \) fs) time, with no sign of rapid decay at the end of the simulation, as shown in figure 2(f).

The fraction \( \eta_{tot} \) of the laser energy dissipated by RF reaches values up to \( \eta_{tot} \approx 0.24 \) for \( a_0 = 600 \) as shown in figure 3(a). A fit to the data gives \( \eta_{tot} \approx a_0^{3.1} \), close to the \( \eta_{tot} \approx a_0^3 \) prediction of our model. Figure 3(a) also shows the peak magnetic field \( B_{max} \) scaling as \( a_0^{3.8} \) up to the highest value \( B_{max} \approx 28B_0 = 3.75 \) G for \( a_0 = 500 \). The decrease down to \( B_{max} \approx 22B_0 \) for \( a_0 = 600 \) is related to the early interruption of the hole boring stage due to the breakthrough of the laser pulse through the target as observed in this case. Notice that we do not show simulations for \( a_0 < 200 \) since in such case the RF losses become too close to the percentage of energy which is lost due to numerical errors (<1%). However, the inferred scaling would predict \( B_{max} \approx 8 \) MG for \( a_0 = 100 \), which may be still detectable making an experimental test closer.
Analytical model for IFE

To sketch an analytical model for IFE, let us first observe that the density of angular momentum of the laser pulse \( I_{\text{L}}(r) \) vanishes on axis and has its maximum at the edge of the beam. We thus consider angular momentum absorption to occur in a thin cylindrical shell of radius \( R \), thickness \( dR \), and length \( h \). The temporal growth of the axial field \( B_x \) induces an azimuthal electric field \( f_E \), which in turn allows the absorbed angular momentum to be transferred from electrons to ions. Assuming that the electron and ion shells rotate with angular velocities \( \Omega_e, \Omega_i \), respectively, we may write for the angular momenta \( M_{\text{ee}} = \mathcal{I}_e \Omega_e \) and \( M_{\text{ii}} = \mathcal{I}_i \Omega_i \), where \( \mathcal{I}_e = 2\pi R^2 h m_e n_e \) and \( \mathcal{I}_i = (A m_i/Z m_e) \mathcal{I}_e \) are the momenta of inertia for electrons and ions, respectively.

The global evolution of the angular momenta of electrons and ions may be described by the equations

\[
\mathcal{I}_e \frac{d\mathcal{I}_e}{dt} = M_{\text{abs}} - M_{\text{E}}, \quad \mathcal{I}_i \frac{d\mathcal{I}_i}{dt} = M_{\text{E}},
\]

where \( M_{\text{abs}} \) is the torque due to angular momentum absorption (related to the absorbed power \( P_{\text{abs}} \) by \( M_{\text{abs}} = P_{\text{abs}}/\omega \)) and \( M_{\text{E}} \) is the torque due to \( E \) fields:

\[
M_{\text{E}} = \int e E_\phi(r) r n_e d^3r \simeq \frac{e E_\phi(R)}{m_e R} \mathcal{I}_e.
\]

The rotation of the electrons induces a current density \( j_\rho \approx -en_e \Omega_e R \). Neglecting the displacement current, in the limiting case \( h \gg R \) the field \( B_x \simeq 4\pi \omega_p \delta/c \) and it is uniform as in a solenoid. In the opposite limit \( h \sim \delta \ll R \), the current distribution may be approximated by a thin wire of cross-section \( \sim h\delta \), and \( E_\phi(R) \) can be obtained via the self-induction coefficient of a coil [63]. We thus obtain

\[
M_{\text{E}} \simeq \mathcal{F} \frac{\omega_p^2 R^5}{2\epsilon^2} \frac{d\mathcal{I}_e}{dt} \equiv \mathcal{I}_e \frac{d\mathcal{I}_e}{dt},
\]

where \( \omega_p = \frac{(4\pi n_e e^2/m_e)^{1/2}}{\sqrt{2}} \) is the plasma frequency. The geometrical factor \( \mathcal{F} \approx 1 \) if \( h \gg R \), and \( \mathcal{F} \approx (h/R) \ln(8R/\sqrt{h\delta}) \) if \( h \sim \delta \ll R \). Therefore

\[
\Omega_e(t) = \frac{1}{\mathcal{I}_e + \mathcal{I}_e t} \int_0^t M_{\text{abs}}(t') dt',
\]

which shows that the electron rotation follows promptly the temporal profile of \( M_{\text{abs}}(t) \), and that effect of the inductive field on electrons is equivalent to effective inertia. Since in our conditions \( \mathcal{I}_e \simeq (\omega_p^2/\omega^2) \mathcal{I}_e \simeq (n_e/n_i) \mathcal{I}_e \gg \mathcal{I}_e \), and therefore \( M_{\text{E}} \simeq \mathcal{I}_e (d\mathcal{I}_e/dt) \gg \mathcal{I}_e (d\mathcal{I}_e/dt) \), the lhs term in equation (5) can be neglected and \( M_{\text{E}} \simeq M_{\text{abs}} \) holds. Thus, from equation (5) we obtain...
\[ L_{\text{abs}} = U_{\text{abs}} / \omega \]

i.e. the total angular momentum of ions is much larger than that of electrons, in agreement with the simulation results (figure 3(b)).

In turn, posing \( M_{\text{abs}} \approx M_{\text{abs}} \) in equation (6) and using \( E_{\phi}(R) \approx - (R / 2 \alpha) G \delta B_{0} (r = 0, t) \) (where \( G = 1 \) for \( h \gg R \) and \( G \approx (2 / \pi) \ln (8 R / \sqrt{\hbar \delta}) \) for \( h \sim \delta \ll R \)) we obtain for the final value of the magnetic field on axis \( B_{\text{cm}} = B_{0} (r = 0, t = \infty) \)

\[
\frac{\pi e}{c} n_{e} h R^{3} \delta G B_{\text{cm}} \approx \int_{0}^{t_{\text{f}}} M_{\text{abs}}(t') \, dt' \approx \frac{T_{\text{f}}}{L_{e}} \gg L_{e}.
\]

The total angular momentum absorbed \( L_{\text{abs}} = U_{\text{abs}} / \omega \) where the absorbed energy is \( U_{\text{abs}} \approx \eta_{\text{lead}} U_{L} \), assuming RF as the main source of dissipation. We thus estimate the final magnetic field as

\[
\frac{B_{\text{cm}}}{B_{0}} \approx \frac{A}{\pi G} \frac{\eta_{\text{lead}}}{n_{e} h} \frac{\eta_{\lambda} \ell_{\lambda}}{\alpha_{0} \omega}.
\]

The product \( n_{a} \) is the surface density of the region where dissipation and angular momentum absorption occur. Thus, with reference to figure 1 we may estimate \( n_{e} h \approx n_{p0} \ell_{\lambda} \approx (\ell_{\lambda} / \pi e c)^{1 / 2} = 2 n_{e} \alpha_{0} \omega / \omega \) (for \( n_{p0} \gg n_{e} \)). Noticing that \( B_{0} / e n_{e} c = 2 \lambda \) we eventually obtain

\[
\frac{B_{\text{cm}}}{B_{0}} \approx \frac{A}{\pi G} \frac{\eta_{\text{lead}}}{\eta_{\lambda} \ell_{\lambda}} \frac{\alpha_{0}}{B_{0}}.
\]

If \( \eta_{\text{lead}} \approx \alpha_{0}^{2} \) then \( B_{\text{cm}} \approx \alpha_{0}^{4} \), in good agreement with the observed scaling in figure 3(a). If we pose \( R \approx r_{0} \), the laser initial beam radius, and \( \delta \approx \lambda \), the radial width of the angular momentum density, for \( a_{0} \approx 500 \), \( \eta_{\text{lead}} = 0.16 \) and \( G = 1 \) equation (12) yields \( B_{\text{cm}} \approx 4.8 B_{0} \). The discrepancy with the observed value of \( \approx 28 B_{0} \) may be attributed to the nonlinear evolution and self-channeling of the laser pulse in the course of the hole boring process. For instance, figure 2 shows that the magnetic field is generated in a region of radius \( \sim 2 \lambda \).

Further analysis of the simulation data shows both a slight increase (by a factor \( \sim 1.2 \)) of the laser amplitude on the axis and a localization of the densities of both EM and mechanical angular momenta in a narrow layer of \( \sim 0.5 \lambda \) width. Posing \( R \approx 2 \lambda \), \( \delta \approx 0.5 \lambda \) and an effective \( a_{0} \approx 600 \) in the above estimate yields \( B_{\text{cm}} \approx 23 B_{0} \), which is in fair agreement with the simulation results considering the roughness of the model.

5. Discussion

In our simulations, RF is the dominant dissipative mechanism (if not the only one at all) allowing for the IFE, i.e. angular momentum absorption and magnetic field generation; collisional absorption is suppressed already at intensities \( \approx 10^{18} \) \( \text{W cm}^{-2} \) even for solid targets, and the comparison between simulations with and without RF shows that collisionless mechanisms (included in the numerical modeling) do not produce any noticeable IFE.

Our simulations use a classical model of RF, based on the LL equation, and do not include quantum electrodynamics (QED) effects. This completely classical approach rises two questions: (a) would a quantum model of RF, which is in principle needed at extremely high intensities, significantly affect the radiative losses and (b) could QED effects including production of electron–positron pairs contribute considerably into absorption of laser radiation, competing with RF as the dominant source of energy dissipation?

Quantum effects on RF are important when the characteristic frequency of emitted photons is comparable to the electron energy, so that the photon recoil is significant. In our case the radiation spectrum peaks at \( \omega_{\text{rad}} \approx a_{0} \omega \) while the electron energy \( \approx a_{0} m_{e} c^{2} \), so that quantum effects are important at \( a_{0} \approx 500 \). However, the simulations of the hole boring process performed in [64] under conditions quite similar to our case show that using a quantum corrected model of RF leads at most to a \( \approx 10\% \) reduction of the conversion efficiency into high-energy radiation, without changing the laser–plasma dynamics qualitatively.

For what concerns the role of other QED effects, in the recent simulations of [31] at intensities close to \( 10^{24} \) \( \text{W cm}^{-2} \) it is found that radiative losses are the dominant energy loss mechanism in solid targets (corresponding to 65% of the pulse energy over a 90% total absorption, to be compared with 5% conversion efficiency into electron–positron pairs). For relatively tightly focused pulses as we use in our simulations, intensities of the order of \( 10^{23} \) \( \text{W cm}^{-2} \) (i.e. about one order in magnitude higher than in our case) are required to develop QED cascades of pairs and photons leading to a considerable depletion of a laser pulse (see e.g., [65]).

We are thus quite confident that RF remains the dominant dissipative mechanism even for the highest intensity applied in the simulations, and that the use of a quantum corrected RF model instead of the classical LL equation would not change the results qualitatively. On the other hand, quantitative modifications on the magnetic field may provide a signature of quantum RF.
6. Conclusions

In conclusion, we showed in 3D simulations that in the interaction of superintense, circularly polarized laser pulses with thick, high density targets the strong radiation friction effects lead to angular momentum absorption and generation of multi-giga gauss magnetic fields via the inverse Faraday effect. Simple models for the efficiency of radiative losses, the transfer of angular momentum to ions and the value of the magnetic field are in fair agreement with the simulation results for what concerns both the scaling with intensity and order-of-magnitude estimates. With the advent of multi-petawatt laser systems, the investigated effect may provide a laboratory example of radiation-dominated, strongly magnetized plasmas and a macroscopic signature of RF, providing a test bed for related theories.

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