Retrieving transient magnetic fields of ultrarelativistic laser plasma via ejected electron polarization

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Interaction of an ultrastrong short laser pulse with non-prepolarized near-critical density plasma is investigated in an ultrarelativistic regime, with an emphasis on the radiative spin polarization of ejected electrons. Our particle-in-cell simulations show explicit correlations between the angle resolved electron polarization and the structure and properties of the transient quasistatic plasma magnetic field. While the magnitude of the spin signal is the indicator of the magnetic field strength created by the longitudinal electron current, the asymmetry of electron polarization is found to gauge the island-like magnetic distribution which emerges due to the transverse current induced by the laser wave front. Our studies demonstrate that the spin degree of freedom of ejected electrons could potentially serve as an efficient tool to retrieve the features of strong plasma fields.

Magnetic fields play a crucial role in various plasma collective phenomena and nonlinear quantum electrodynamic processes in extreme environments of laboratory and universe [34]. The astrophysical magnetic fields can govern the internal structure of interstellar shocks [4], mediate the radio wave emission nearby neutron stars [5], induce baryon inhomogeneities [6], and catalyse the dark matter formation [7]. Self-generated fields with strength \( \sim 10^3 \) Tesla have been produced in high-intensity plasma experiments [8–12], and the guidance of jet flows by laboratory magnetic fields helps interpret the evolution of young stellar objects [13–15]. With recent advancement of ultrastrong laser techniques [16–22] more extreme conditions and larger fields are expected in ultrarelativistic laser plasma interaction [23–29].

Generally, detection of plasma magnetic fields requires an external probe beam, where the field information is imprinted in the velocity space of charged particles [30–35] or the rotated polarization vector of the optical beam [36–39]. However, these conventional methods are inapplicable for scenarios with unprecedented field strength, ultrashort time scale (\( \sim \text{fs} \)), and overcritical plasma density [10]. Furthermore, the spin, an intrinsic property of particles, offers a new degree of freedom of information, which is widely utilized in exploring magnetization of solids [40], nucleon structure [41], and phenomena beyond the standard model [42]. In extreme laser fields there is a strong coupling of the electron spin to the laser magnetic field [44–48], which may yield radiative spin polarization (SP) [49–52], i.e., polarization of electrons due to spin-flip during photon emissions. Even though in the oscillating field the electron net SP is suppressed, fast polarization of a lepton beam with laser pulses becomes possible when the symmetry of the monochromatic field is broken, such as in an elliptically polarized, or in two-color laser pulses [53–57]. Because of collective effects, more complex spin dynamics occurs in strong laser field interaction with plasma. Consequently, the question arises if it is possible to employ the spin signal of spontaneously ejected particles from plasma to retrieve information on transient plasma fields.

In this Letter, based on particle-in-cell (PIC) simulations, we investigate the ultrarelativistic dynamics of a short strong pulse interacting with a non-prepolarized near-critical density plasma, see Fig. 1. Special attention is devoted to describing the spin dynamics of plasma electrons, being strongly disturbed by the radiative spin-flips modulated by the quasistatic plasma magnetic field (QPMF). The latter is commonly transient with a time scale as short as the driving pulse duration while being quasistatic with respect to the fast oscillating laser field. We show that the angle dependent SP of ejected electrons carries signatures of the inhomogeneous QPMF. The signal of SP of ejected electrons can be used to predict the strength of the leading order antisymmetric QPMF created by the longitudinal current. A more detailed analysis reveals that the asymmetry of SP of two outgoing divergent electron bunches characterizes the secondary QPMF, which is induced by a transverse transient current and generally neglected in previous studies [58–62]. The sum of these two part QPMFs gives rise to a nonlinear island-like magnetic structure [see Fig. 2(a)(b)]. Our results demonstrate that the spin degree of freedom of ejected electrons from ultrarelativistic plasmas can be employed in principle as a tool to retrieve information on the QPMF structure and properties.

In 2D PIC simulations, a near-critical density target is irradiated by a linearly polarized pulse (with the transverse electric field along \( y \)). Our main example adopts a peak intensity of \( 1.7 \times 10^{23} \) W/cm\(^2\), equivalent to the normalized field amplitude \( a_0 = 350 \) given the laser wavelength \( \lambda_0 = 1 \mu \text{m} \). The pulse has a 2.6 \( \mu \text{m} \) focal spot size and 18 fs duration (FWHM intensity measure). The target has thickness \( l_0 = 10 \mu \text{m} \) and electron (carbon) density \( n_e = 5n_i \) (\( n_i = n_e/6 \)), where \( n_i = m_\omega c^2/4\pi e^2 \) is the plasma critical density for a laser frequency \( \omega_0 = c\kappa_0^0 \); \( m_e \) (\( e \)) the electron mass (charge); \( c \) the speed of light. The dynamics of spin precession is governed by the Thomas-Bargmann-Michel-Telegdi equation and spin-dependent photon emissions have been implemented in the EPOCH simulation code.
θ_1 > 0 (θ_1 < 0). The spatial evolution of SP in Fig. 1(b) manifests that two groups of electrons are firstly polarized and confined inside the channel, and then intersect with each other towards the opposite transverse direction. This procedure is also unveiled by the evolution of SP ⟨s_z⟩ in the transverse phase space (y, θ) in Fig. 1(c), where θ = arctan(p_y/p_x) denotes the direction of electron momentum. The clockwise rotation of ⟨s_z⟩ indicates that the QPMF not only generates spatial dependent SP but also deflects the electrons to form an angle dependent polarization distribution of ejected electrons. In Fig. 1(d), asymmetry exists for both electron SP and number angular distributions. Specifically, the averaged SP (final angle) with a positive θ_1 is ⟨s_z⟩ ≈ −3.3% ((θ_+)) ≈ 11.4°), whereas ⟨s_−⟩ ≈ 4.0% ((θ−) ≈ −16.5°) for θ_1 < 0. The magnitude of the SP signal is characterized by the parameter δ⟨s_z⟩ ≡ ⟨s_+⟩ − ⟨s_−⟩. According to Fig. 1(e), SP is insignificant for low-energy electrons because of damped radiative spin-flips. Therefore, the criterion of ε_0 > 4a_0m_0c^2 is adopted here to filter out the low-energy noise. To reveal more subtle features of QPMF B_z, we introduce also the spin (angle) asymmetry characteristics via the absolute difference: Δ ⟨s_z⟩ = |⟨s_+⟩| − |⟨s_−⟩| and Δ (θ) ≡ |⟨θ_+⟩| − |⟨θ_−⟩|, which will be discussed below.

The QPMF B_z is determined by electric currents via ∂B_z/∂y = μ_0J_x and ∂B_z/∂x = −μ_0J_y (with the vacuum permeability μ_0). In general, inside a laser-driven plasma channel, the current is dominated by the longitudinal one J_x and the transverse current J_y is neglected [58–62]. However, the magnetic field in our simulation shows an irregular structure, with multiple islands associated with the current kinks and vortices, see Fig. 2(b). The latter indicates that the transverse current J_y is important in characterizing the exact form of B_z. Let us divide QPMF into two parts B_z = B_z,1 + B_z,2, where the leading part B_z,1 is induced by J_x, while the secondary B_z,2 by J_y: ∂B_z,1/∂y = μ_0J_x and ∂B_z,2/∂x = −μ_0J_y. The leading part B_z,1 ∝ −μ_0|e|n_c e/pyc with antisymmetric feature B_z,1(y) = −B_z,1(−y) is ubiquitously utilized in previous studies [58,59]. Now, we focus on the secondary B_z,2. Considering the electron velocity v_0 = p_y/(γ_0 m_e c) and momentum p_0 ∼ A_y = a_0 cos(ξ + φ_0) where ξ = γ_0 t − k_0 x and φ_0 the carrier envelop phase (CEP), we obtain J_y ∼ −|e| ∫ n_0 δ(x/v_y − t)v_y dt ≈ |e|n_0 c^2 cos(ωt/v_y − x/v_p) + φ_0. The δ(t−x/v_y) function indicates that the transverse current is predominantly contributed by the electron density n_0 δ(t − x/v_y) piled up at the front edge of the plasma channel nearby the region x ∼ v_y t, where the electron’s transverse velocity is significant. Here, v_p (v_pn) is the laser group (phase) velocity in plasma, and the Lorentz-factor γ ∼ a_0 is assumed. With ∂B_z,2/∂x = −μ_0J_y, the secondary magnetic field can be estimated:

$$B_{z,2} \sim -\frac{μ₀|e|n_0}{k_2} \sin(k_2 x + φ_0)$$

where k_2 = k_0(v_p - v_y)/v_y. The analytically predicted

FIG. 1. (a) The interaction scheme: the laser pulse impinges on unpolarized plasma (the electron density is shown in gray shades); accelerated and radiatively polarized electrons due to spin-flips form outgoing polarized bunches. The red (blue) dots represent the electrons with spin s_z = 1 (−1) and the lines show their movement tendency. The green line shows a typical electron trajectory with a spin-flip marked by a pentagram. (b) and (c) show snapshots of the electron SP distribution in spatial (x, y) coordinates and transverse (y, θ) phase space, respectively. The green lines in (b) profile the laser field E_y at slice y = 0. (d) Angular distribution of electron number dN/dθ_1 and SP ⟨s_z⟩. (e) δ ⟨s_z⟩ and dN/dδε vs electron energy ε_0. All parameters are indicated in the text.

When the pulse impinges on the target, a fraction of bulk electrons is expelled outwards by the laser ponderomotive force to form a plasma channel [65]. Meanwhile, the peripheral electrons are prone to be injected [66] and subsequently polarized inside the channel due to spin-flips during photon emissions, see Fig. 1(b). Since the ion reaction partially compensates the transverse charge separation [67], the quasistatic electric field E_y is negligible in this scenario. Thus the deflection of the accelerated electrons in transverse direction is predominantly governed by the azimuthal QPMF B_z, which is presumably sustained by the longitudinally forward moving electron current J_x. The simulation results in Fig. 1(a) show that the electrons with a positive (negative) final angle θ_1 mainly originate from the plasma region of y < 0 (y > 0). As the magnetic field B_z ∝ −μ_0J_0|y| is antisymmetric, created by the nearly uniform current J_x ∝ −|J_0|, the electrons exiting the plasma area with a final angle θ_1 > 0 mostly experience a positive B_z [see Fig. 1(a)] and vice versa. This leads to oppositely SP ejected electron bunches: ⟨s_z⟩ < 0 (⟨s_z⟩ > 0) for the electron bunch of
when the electron spin-flip is dominated by the plasma
spin-flip as $\chi_\text{ph}$ when emitting an energetic photon with $\chi_\text{ph}$ of the electron or photon, respectively. As Fig. 2(c) il-

The electron quantum invariant parameter is

The blue line profiles the photon emission probability $d^2N_{ph}/d\chi_\text{ph}dt$ with a bandwidth accounting for
the influence of electron spin. (d) The number distribution of all emitted photons $N_{ph}$ (grey) and emission associated with
spin-flips $N_{sf}$ (yellow). (e) The spatial dependent SP difference $ds/dz$ contributed by the cases of $\kappa \leq 1$ for electrons
with final angle $\theta_f > 0$ and $\theta_f < 0$, respectively. (f) The angular dependence of spin-flip occurrence, where the result
of condition $\kappa > 1$ is multiplied by 10 for better visibility.

$B_z = B_{z,1} + B_{z,2}$ is shown in Fig. 2(a), which agrees qualitatively with the simulated $B_z$ in Fig. 2(b). The
asymmetric periodic island-like structure of QPMF $B_z$ stems from the nontrivial current vortex $(\nabla \times j)_{x,y} \neq 0$
generated by the transverse current of electrons plough away by the laser beam front.

As we are interested in the relation of the electron SP to the magnetic field structure, and considering the polarization attributable to the spin-flip during a photon emission, we analyze the probability of this process $P(\chi_\text{ph})$
in Fig. 2(c) for typical parameters of our PIC simulations. Here, the electron with an initial $\gamma_e = 2000$ normally crosses the uniform magnetic field $B_0 = 10^4T$, and the electron quantum invariant parameter is $\chi_e \sim 0.1$, with $\chi_{e,\text{ph}} = (e\hbar/c^3)|F_{\mu 0}|^2$ and the momentum $p''$ of the electron or photon, respectively. As Fig. 2(c) il-
nuates, the electron spin-flips exclusively take place when emitting an energetic photon with $\chi_\text{ph}$ close to $\chi_e$, while the photon emission probability is peaked at $\hbar \omega_e \sim \chi_e \gamma_e m_e c^2$ (at $\chi_e < 1$), i.e., the peak of the spin-flip process is shifted with respect to the photon emission to higher $\chi_e$'s, see Fig. 2(d). Both the laser magnetic field $B_0$ and QPMF $B_z$ can cause the electron spin-flip as $\chi_e \sim \gamma_e[(1 - \cos \theta)B_0 + B_{z,1}]/B_c$ with the Schwinger limit $B_c \sim 4 \times 10^7 T$. We introduce the parameter $\kappa = |B_{z,1}/[(1 - \cos \theta)B_0]|$, defining two regimes, when the electron spin-flip is dominated by the plasma
($\kappa > 1$) or by the laser field ($\kappa < 1$). The evolution of SP in Fig. 2(c) demonstrates that the laser field dominated regime ($\kappa < 1$) mostly contributes to the final electron SP. A distinguishing feature between the $\kappa \leq 1$ regimes is the angle $\theta$ of the electron’s instantaneous momentum when the spin-flip occurs. As the angular dependent spin-flip shows in Fig. 2(f), the $\kappa < 1$ regime applies at backward emissions, while $\kappa > 1$ for forward ones.
The detailed particle tracking further confirms these conclusions. In the \( \kappa > 1 \) regime Figs. 3(a), (b), the position of spin-flip with \( \kappa > 1 \) is closely correlated with the spatial distribution of QPMF \( \overline{B}_2 \). The time evolution of \( p_x \) illustrates that the spin-flip happens after the electron starts an efficient acceleration and its velocity aligns longitudinally \( \theta \ll 1 \), resulting in \( (1 - \cos \theta) B_l < \overline{B}_2 \). For the laser dominant regime \( \kappa < 1 \), the electron trajectory and momentum evolution [in Fig. 3(c),(d)] demonstrate that the typical spin-flip occurs at the electron’s temporarily backward motion, when \( (1 - \cos \theta) \sim 1 \) and \( B_l > \overline{B}_2 \).

It should be noted that even in the laser dominant regime, the QPMF \( \overline{B}_2 \) is still the key factor for the SP. The reason is that the laser field has oscillating character. Although it can cause spin-flips, its net contribution to the final SP is negligible. The laser magnetic field \( B_l \) acts as a catalyst to enhance the electron spin-flips by increasing \( \chi_e \) and net spin-flips contributing to the final SP are still determined by \( \overline{B}_2 \) [64]. We may estimate \( s_z \approx -\int \overline{B}_2(|B_l|,|\chi_e|)dt \) (at \( B_l \gg \overline{B}_2 \)), with \( \chi_e = \sqrt{\alpha_0 f m_e^2 \chi_e}/(h \gamma_e)A_2^*(\chi_e) \), and \( A_2^*(\chi_e) \approx 0.18 \chi_e \) (at \( 0.01 < \chi_e < 0.4 \)) [64]. The electrons with final angle \( \theta_f < 0 \) are mainly expelled to the QPMF \( \overline{B}_2 > 0 \) at the region \( y < 0 \) (\( y > 0 \)), and the overall SP with \( \theta_f > 0 \) would be \( s_z < 0 \) (\( s_z > 0 \)) which are illustrated as the solid black lines in Fig. 3(e),(f).

Thus, we calculate the electron’s SP magnitude \( \delta (s_z) \) being correlated with the leading order QPMF \( \overline{B}_{z,1} \):

\[
\delta (s_z) \sim -\frac{|e|}{m_e \omega_0} a_0^2 \frac{\sim \eta}{\overline{B}_{z,1}} \approx \eta \frac{|e|}{m_e \omega_0} a_0^2 \frac{\sim \eta}{\overline{B}_{z,1}} \quad (2)
\]

where \( \gamma_e \sim a_0 \) is used, and \( \eta \approx 4 \times 10^{-8} \) accounts for the deviations from the radiative spin evolution. With \( \overline{B}_{z,1} \approx \sqrt{a_0/4\pi}(n_e/n_c)(m_e \omega_0/|e|) \), we find the SP scaling \( \delta (s_z) \propto a_0^{-5/2} \), as well as the relation \( \overline{B}_{z,1} \approx \sqrt{a_0/(4\pi n_e)}(m_e \omega_0/|e|) \). In Fig. 4(a),(b),(c), the analytically predicted scalings of \( \delta (s_z) \) and \( \overline{B}_{z,1} \) are in good accordance with the 2D simulation results.

Finally, we show how with the help of the SP asymmetry signal \( \Delta (s_z) \) defined above, the secondary QPMF can be retrieved. In the \( \Delta (s_z) \) signal the contribution of the \( \overline{B}_{z,1} \) is cancelled, and \( \Delta (s_z) \approx 2 \int \overline{B}_{z,2}(|B_l|,|\chi_e|)dt \). Since \( \overline{B}_{z,2} \sim b_2 \sin(k_2x + \phi_0) \) is oscillating along the longitudinal position (along the laser CEP), the overall effect of \( \overline{B}_{z,2} \) imprinted on the signal of \( \Delta (s_z) \) is oscillating as well. Taking into account the results for \( \delta (s_z) \) and \( \overline{B}_{z,2} \), we find for the asymmetry signal

\[
\Delta (s_z) \sim \delta (s_z) \frac{b_2}{|\overline{B}_{z,1}|} k_0 \cos(k_2l_0 + \phi_0), \quad (3)
\]

where \( b_2 \) is the amplitude of \( \overline{B}_{z,2} \). The oscillating dependence of \( \Delta (s_z) \) on the laser CEP \( \phi_0 \) is reproduced by the simulation results in Fig. 4(d).

We see that the amplitude of the SP asymmetry signal \( \Delta (s_z) \) is directly related to the secondary QPMF \( b_2 \):

\[
\Delta (s_z) \sim \frac{\delta (s_z)}{|\overline{B}_{z,1}|} b_2 k_0 \cos(k_2l_0 + \phi_0).
\]

FIG. 5. (a) The \( \overline{B}_2 \) predicted by the model based on the sign of \( \Delta (s_z) \) and \( \Delta (\theta) \). The valid range of the model is illustrated in the white region of the parameter space in (b) \( a_0 = 350 \) and (c) \( n_c = 5n_e \), where the circles (triangles) mark the simulation results of \( \overline{B}_2 \) with no more (more) than two islands on each side of \( y = 0 \). (d) The \( \overline{B}_2 \) obtained from simulation for the case of \( a_0 = 350 \), \( n_c = 5n_e \), and \( l_0 = 20\mu m \).
tigated the role of experimental imperfections and uncertainties, in particular, the asymmetry in the driving laser pulse, and the ramp-up and -down of the plasma density profile [64]. The simulation results indicate that the presented scheme is robust to moderate imperfections of such practical issues. It should be noted that distinguishing more complex field structures, e.g., the three-island structure like that in Fig. 3(d), could be achievable with modifications of the retrieval method, see an example in [64], which however needs further exploration.

In conclusion, the ejected electron spin provides a new degree of freedom to extract information on the structure and magnitude of different components of the transient plasma fields. Our results open a new avenue for the electron spin-based plasma diagnostics in extreme conditions, which are prevalent in astrophysical environments and are expected in near future laser facilities.

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