

Spin-dependent axion generation with controllable emission angles in strong-laser fields

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We investigate axion production in the collision between a spin-polarized relativistic electron beam and an ultraintense laser pulse. A spin-resolved Monte Carlo framework is developed to model axion-electron and axion-photon couplings in arbitrary electromagnetic fields, using quantum emission probabilities under the local constant field approximation. Owing to spin-dependent asymmetries in the radiation probability, the emitted axions acquire a characteristic angular deflection correlated with the initial electron polarization, enabling control of the axion emission via the polarizations of the electron beam and laser field. Simulations demonstrate that a dense ($\sim 10^{10} g_{ae}^2$) and highly collimated axion beam with a tunable milliradian-scale deflection can be produced within tens of femtoseconds using current laser technology. The polarization-tagged angular asymmetry, combined with femtosecond time gating, provides enhanced signal-background discrimination. These results establish a spin-dependent mechanism for manipulating axion trajectories and open a promising route toward laboratory-based searches for the axion-electron coupling.

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The QCD axion, originally proposed by Weinberg and Wilczek as a solution to the strong CP problem, provides a well-motivated extension of particle physics [1–4] and is also a leading candidate for dark matter [5–7]. In addition to the QCD axion, similar particles called axionlike particles (ALPs) also emerge in string theory [8,9] and supersymmetric extensions [10]. Various experimental searches have been conducted and are ongoing to detect ALPs and impose stricter limits on ALP models. These limits can be broadly categorized into three types: cosmological [11–23], astrophysical [24–47], and laboratory based [48–73]. While cosmological and astrophysical limits generally provide stronger constraints than laboratory-based searches, they are highly model dependent and can be significantly weakened [74,75]. The vulnerability in the astrocosmological constraints provides a strong motivation to pursue well-controlled laboratory-based searches.

Laboratory-based searches for axionlike particles include light-shining-through-walls (LSW) experiments

such as OSQAR [50] and ALPS [49], which rely on photon-axion conversion and subsequent regeneration across an opaque barrier [76]. The PVLAS Collaboration [51] probes axion-induced vacuum dichroism by measuring the polarization rotation of a laser beam propagating through a magnetic field. However, the sensitivity of these setups is fundamentally limited by the achievable magnetic field strength ($|B| \sim 10^5$ G), the finite interaction length ($L \sim 1$ km), and the requirement of phase coherence between the photon and axion waves. Furthermore, in PVLAS, the tree-level axion signal is superimposed with the QED-induced vacuum birefringence, making signal isolation particularly challenging [77]. These limitations motivate the development of alternative approaches capable of overcoming coherence constraints, suppressing QED backgrounds, and accessing unexplored regions of the axion parameter space.

The advent of ultraintense laser systems enables electromagnetic field strengths reaching $B \sim 10^{10}$ G on ultrashort timescales, exceeding those attainable in conventional static-field setups by several orders of magnitude [78,79]. These extreme conditions, together with high photon fluxes and short temporal scales, offer new opportunities to explore ALP interactions in previously inaccessible regimes. In particular, the large momentum transfer in laser-beam collisions extends sensitivity to higher ALP masses [80–84]. To date, theoretical investigations have

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predominantly focused on ALP-photon couplings [85–88]. However, several well-motivated extensions of the Standard Model predict direct couplings to electrons—such as the Arion [89], Dine-Fischler-Srednicki-Zhitnitsky (DFSZ)-like scenarios [90], and string-inspired ALPs [91], which remain largely unexplored. These couplings, in principle, allow for direct ALP production through collisions between relativistic electrons and intense laser pulses, opening a complementary channel for laboratory-based ALP searches.

The potential of intense laser pulses for probing electron-ALP coupling has been investigated in the low-energy and coherent limit. Previous studies have shown that nonlinear Thomson scattering of light scalars in a monochromatic electromagnetic background can extend the sensitivity to ALP masses in the range of 10–100 eV [92]. Electron-seeded ALP production and decay in a monochromatic, circularly polarized electromagnetic background have been studied previously [93–95]. However, the corresponding probabilities are expressed as sums over photon number n , whose evaluation becomes prohibitively expensive in the strong-field regime, limiting their applicability. This difficulty can be overcome using the local constant field approximation (LCFA) [96], which is widely employed in strong-field QED and is reliable at sufficiently large laser amplitudes $a_0 \equiv A_0/m_e \gg 1$, where m_e is the electron mass and A_0 is the amplitude of the laser vector potential. While the radiation probability of ALPs has been obtained by averaging over the initial electron spin and summing over the final spin states [97], recent studies have demonstrated that spin effects in strong-field QED processes can have significant consequences for the polarization of final particles and the electron kinetic motion [98–110]. The spin degree of freedom, although fundamental in strong-field QED, has not yet been leveraged in axion production models, leaving a promising direction largely unexplored.

In this work, we study the production of axions via nonlinear Compton scattering from polarized electrons in intense laser pulses (hereafter “axions” refers to ALPs for brevity); see Fig. 1. To this end, we have developed a spin-resolved Monte Carlo code that self-consistently incorporates the couplings of axions to electrons and photons. We find that spin effects, neglected in previous studies, can induce a notable deflection of the produced axions; this deflection can be controlled by tuning the laser ellipticity and the electron spin polarization. The asymmetry in the angular distribution could serve as a distinctive signal of the axion-electron coupling [Fig. 1(c)], clearly separated from its QED counterpart, i.e., photon emission [Fig. 1(b)], as well as from dark counts. In addition, the high single-shot yield and femtosecond-scale emission duration jointly contribute to an enhanced signal-to-noise ratio. Our scheme utilizes the polarization freedom to advance the control over laboratory-based searches of the axion and may

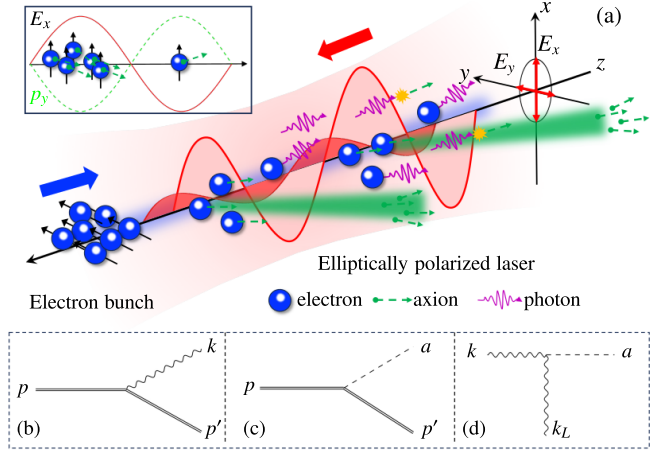


FIG. 1. (a) Illustration of axion generation schemes in strong laser fields. The electrons with initial polarization of $\zeta_y = 1$ collide with an elliptically polarized laser pulse. The inset demonstrates that the axions are selectively generated at the positive half-cycle ($E_x > 0$) and emitted toward $p_y < 0$. Feynman diagrams illustrating photon emission (b) and axion production via electron-axion coupling (c) and the photon-photon coupling (d).

provide useful complementary bounds to those obtained in other laboratory-based and astrophysical experiments.

In our simulations, the electron trajectory is determined by classical dynamics governed by the Lorentz force and the Bargmann-Michel-Telegdi equation, while incorporating quantum effects such as stochastic photon emission, spin polarization due to radiative spin-flip processes, and radiative corrections from the one-loop electron self-energy [111]. The influence of axion generation on the dynamics is statistically negligible due to the smallness of the coupling constant, $g_{ae} \ll \alpha \sim \mathcal{O}(10^{-2})$ [111]. Consequently, axion production from electrons via Compton scattering and from photons via the Primakoff effect [Fig. 1(d)] is evaluated by integrating the spin-resolved emission probability [Eq. (2)] and the photon-axion conversion probability [111] along the particle trajectories, enabling studies of QED-axion interplay in complex strong-field configurations. We consider the pseudoscalar axion-electron interaction

$$\mathcal{L}_{\text{int}} = -ig_{ae}\phi\bar{\psi}_e\gamma^5\psi_e, \quad (1)$$

where ϕ is the axion field and ψ_e is the electron field. The spin-dependent axion emission probabilities in LCFA are employed with the leading-order contribution with respect to $1/\gamma$ and are derived with the quantum electrodynamics operator method of Baier-Katkov. The probability of emitting an axion with energy of ω_a reads [131]

$$\begin{aligned} \frac{dW_a}{d\omega_a dt} = C_a \left\{ -\left[\frac{\varepsilon\varepsilon'}{2\omega_a^2} \delta_m^2 \right] \text{Int}K_{\frac{1}{3}}(z_q^a) + \left(1 + \frac{\varepsilon\varepsilon'}{\omega_a^2} \delta_m^2 \right) K_{\frac{2}{3}}(z_q^a) \right. \\ \left. + \left(1 + \frac{\varepsilon\varepsilon'}{\omega_a^2} \delta_m^2 \right)^{\frac{1}{2}} (\vec{\zeta} \cdot \vec{b}) K_{\frac{1}{3}}(z_q^a) \right\}. \end{aligned} \quad (2)$$

Here, $C_a = \frac{\sqrt{3}}{24\pi^2} \frac{g_{ae}^2 \omega_a^2 m_e^2}{\epsilon^3 \epsilon'}$, $z_q^a = \frac{2}{3} \frac{\omega_a}{\chi_e \epsilon'} (1 + \frac{\epsilon \epsilon'}{\omega_a^2} \delta_m^2)^{\frac{3}{2}}$, and $\delta_m = \frac{m_a}{m_e}$, where m_a and m_e denote the masses of the axion and the electron, respectively. The electron energies before and after axion emission are denoted by ϵ and ϵ' , respectively. The vector $\vec{\zeta}$ is the polarization vector of electron before emission. The unit vectors $\vec{v}$ and $\vec{s}$ correspond to the velocity and acceleration directions, respectively, with $\vec{b} = \vec{v} \times \vec{s}$. The integral of the modified Bessel function of the second kind is defined as $\text{Int}K_{\frac{1}{3}}(z_q^a) \equiv \int_{z_q^a}^{\infty} dz K_{\frac{1}{3}}(z)$, where K_n denotes the modified Bessel function of the second kind of order n . The parameter χ_e is the nonlinear quantum parameter controlling the strength of quantum effects. Upon averaging over the initial electron spin, the spin-unresolved radiation probability derived in previous literature is recovered [97]. In contrast to photon emission, the axion production rate in Eq. (2) is free of infrared divergences [131].

The initial spin polarization of electrons significantly influences the angular distribution of axions, as shown by the simulations in Fig. 2, which were performed using the following laser parameters: a peak intensity of $I_0 \approx 1.38 \times 10^{22}$ W/cm² ($a_0 = 100$), a wavelength of $\lambda_0 = 800$ nm, a pulse duration of $\tau = 10T_0$ (where T_0 is the laser period), a focal radius of $w_0 = 4$ μ m, and an ellipticity of $\epsilon = |E_y|/|E_x| = 0.05$. The cylindrical electron bunch has an angular divergence of 0.3 mrad and an initial kinetic energy of $\epsilon_0 = 2$ GeV with a relative energy spread of $\Delta\epsilon_0/\epsilon_0 = 0.05$. The maximum quantum parameter reaches $\chi_{\max} \approx 1.5$. The bunch has $N_e = 1 \times 10^6$ electrons, radius $w_e = \lambda_0$, and length $L_e = 5\lambda_0$, with a transversely Gaussian and longitudinally uniform distribution, which can be obtained by current laser wakefield accelerators. The axion yield scales with the coupling parameter g_{ae} and electron number N_e as $N_a \approx 3.03 N_e g_{ae}^2$, indicating that the axion yield can reach $N_a = 3.03 \times 10^{-10}$ for $N_e = 1 \times 10^{10}$ and $g_{ae} = 10^{-10}$. The corresponding single-shot axion flux can be estimated as $\Psi = \frac{N_a}{\pi w_e^2 L_e} \approx 1.15 \times 10^{32} g_{ae}^2 \text{ cm}^{-2} \text{ s}^{-1}$ [132]. For comparison, solar axions produced via the Compton process yield a flux at Earth of $\Psi \approx 2.16 \times 10^{34} g_{ae}^2 \text{ cm}^{-2} \text{ s}^{-1}$ [133,134]. Our scheme can access a broader axion-mass range, and its scalability with laser intensity and electron energy allows the flux to approach the solar level [111]. Electrons with $\zeta_y = 1$ predominantly generate axions with $p_y^f < 0$ [Fig. 2(a)], while those with $\zeta_y = -1$ produce axions with $p_y^f > 0$ [Fig. 2(b)]. Thus, for electrons with initial spin polarization $\zeta_y = 1$, the majority of axions are deflected to $\theta_y > 0$, accounting for 74.83% of the axions generated within the angular window $|\theta_x| < 20$ mrad and $|\theta_y| < 1$ mrad. This spin dependence arises from the explicit dependence of the axion radiation probability on the initial electron spin,

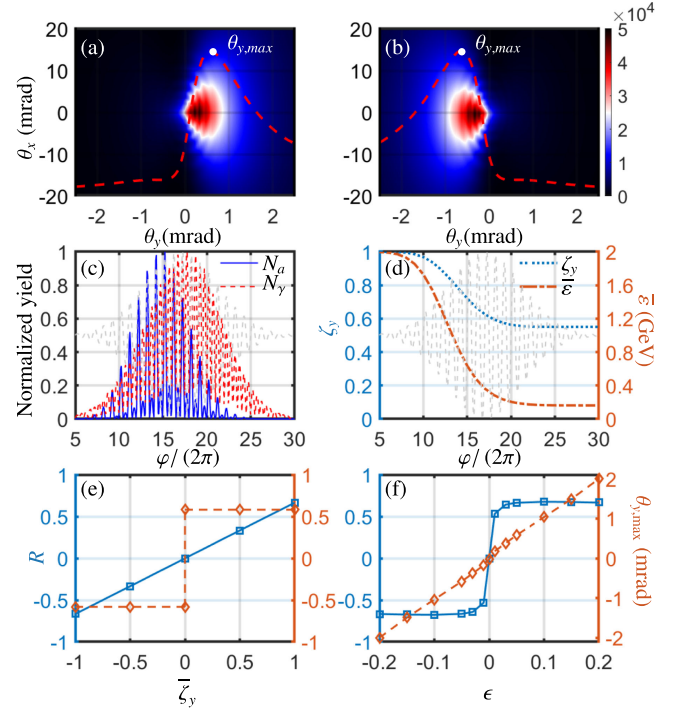


FIG. 2. The angular distribution of axion density $d^2(N/g_{ae}^2)/d\theta_x d\theta_y$ (mrad⁻²) for initial electron spin $\zeta_y = 1$ (a) and $\zeta_y = -1$ (b), plotted against the deflection angles $\theta_x = \arctan k_{ax}/k_{az}$ and $\theta_y = \arctan k_{ay}/k_{az}$. The red dashed lines indicate the angular distribution of axion density $d(N/g_{ae}^2)/d\theta_y$ vs θ_y for $\zeta_y = 1$ [in (a)] and $\zeta_y = -1$ [in (b)]. (c) Normalized axion number N_a (blue solid line) and γ photon number N_γ (red dashed line) at production points, vs laser phase $\varphi/(2\pi)$. (d) Left y axis: polarization ζ_y (blue dotted line) vs laser phase $\varphi/(2\pi)$; Right y axis: electron average energy $\bar{\epsilon}$ (red dashed-dot line) vs laser phase $\varphi/(2\pi)$. The superimposed gray lines in (c) and (d) represent the x component of the electric fields. The remaining parameters are $\epsilon = 0.05$, $m_a = 0$, and $\bar{\zeta}_y = 1$. The angular asymmetry R (blue-solid line) and deflection angle $\theta_{y,\max}$ (red-dashed line) vs (e) average polarization of electron $\bar{\zeta}_y$ and (f) ellipticity of laser ϵ .

$$w(\vec{\zeta}) = C_a [K_{\frac{2}{3}}(z_q^a) + (\vec{\zeta} \cdot \vec{b}) K_{\frac{1}{3}}(z_q^a)], \quad (3)$$

where $\vec{b}$ is the unit vector aligned with the local magnetic field direction in the considered setup. For electrons with $\zeta_y > 0$, the radiation probability during the half-cycle when $B_y > 0$ is significantly higher than during the half-cycle when $B_y < 0$, resulting in a greater axion yield when $B_y > 0$. The final momentum of the axion is determined by the electron momentum at the emission time, i.e., $\vec{k}_a^f = \vec{p}(t_r)$. In an elliptically polarized laser field, the y -component of the electron momentum p_y exhibits a π -phase delay relative to the magnetic field B_y . Consequently, axions predominantly generated during $B_y > 0$ acquire

negative transverse momentum p_y , corresponding to a positive emission angle $\theta_y = \arctan(p_y/p_z)$ [Fig. 2(a)]. Conversely, for electrons with $\zeta_y < 0$, the emission angle becomes $\theta_y < 0$ [Fig. 2(b)]. The central deflection angle is approximately $|\theta_{y,\max}| \approx 0.5$ mrad. Compared to γ photon emission, axion production exhibits a much more pronounced cycle-dependent asymmetry [Fig. 2(c)], leading to a significantly enhanced angular asymmetry in axion-electron coupling compared to its QED counterpart, photon-electron coupling. Meanwhile, as electrons propagate through the laser pulse, they undergo substantial energy loss due to radiation reaction and depolarization caused by spin-flip transitions induced by photon emission [Fig. 2(d)]. Consequently, the cycle-dependent asymmetry gradually diminishes with increasing laser phase and eventually vanishes at the pulse tail [Fig. 2(c)]. This indicates that the asymmetry in the axion angular distribution is primarily governed by axion production in the early phase of the laser pulse, where electrons are highly energetic and polarized.

More interestingly, our scheme enables axion generation with a tunable angular asymmetry by adjusting either the initial electron spin polarization [Fig. 2(e)] or the ellipticity of the laser field [Fig. 2(f)]. The spin-related asymmetry can be defined as $R = (N_+ - N_-)/(N_+ + N_-)$, where $N_{\pm}$ denote the numbers of axions emitted at $\theta_y > 0$ and $\theta_y < 0$, respectively. When the incident electrons are in a pure spin state with $\zeta_y = -1$, most axions are deflected toward $\theta_{y,\max} = -0.58$ mrad [Fig. 2(e)]. As the average polarization degree $\bar{\zeta}_y$ increases from -1 to 1 , the number of axions deflected to $\theta_y > 0$ increases, leading to a nearly linear increase of the asymmetry R [Fig. 2(e)]. However, the central deflection angle $\theta_{y,\max}$ of the density peak remains unchanged until $\bar{\zeta}_y$ flips sign. In contrast, as the laser ellipticity increases from $\epsilon = -0.2$ to 0.2 , the central deflection angle increases monotonically from negative to positive values, owing to the dependence of the transverse momentum p_y on the field strength E_y . The asymmetry R vanishes when the laser is linearly polarized but starts increasing rapidly as ellipticity is introduced [Fig. 2(f)]. With increasing absolute value of the ellipticity parameter $|\epsilon|$, the impact of E_y on the electron motion becomes dominant over other effects (e.g., ponderomotive force and radiation reaction), leading to a stronger correlation between p_y and B_y and, consequently, a larger R . As $|\epsilon|$ continues to increase, the growth of $|R|$ slows down and eventually reaches saturation for $|\epsilon| \geq 0.1$. Thus, tuning the laser ellipticity allows precise control over the central peak of the deflection angle, whereas the density asymmetry is more strongly influenced by the electron polarization.

The produced axions can be detected through their conversion into photons in a magnetic field. For representative LSW parameters, $B = 5$ T and $L = 4.21$ m, the photons converted from axions inherit the spin-induced

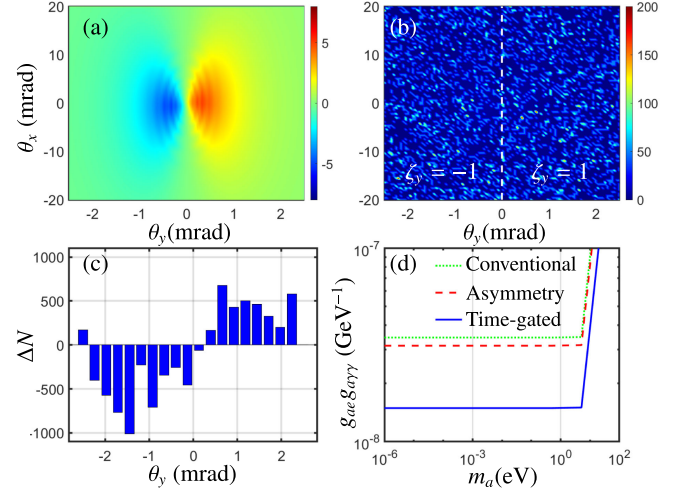


FIG. 3. (a) Difference of the regenerated photon yield defined as $\Delta N = d^2N_\gamma(\zeta_y = +1)/d\theta_x d\theta_y - d^2N_\gamma(\zeta_y = -1)/d\theta_x d\theta_y$, versus θ_x and θ_y . (b) Angular distribution $d^2N_\gamma/d\theta_x d\theta_y$ (mrad⁻²) including background noise, shown for electron spin states $\zeta_y = -1$ (left) and $\zeta_y = +1$ (right). (c) Difference of the regenerated photon yield ΔN versus θ_y , including background noise. $g_{ae}g_{\gamma\gamma} = 1.45 \times 10^{-8}$ GeV⁻¹ for (a)-(c). (d) Projected 95% C.L. exclusion limits on the coupling $g_{ae}g_{\gamma\gamma}$ (GeV⁻¹) as a function of the axion mass m_a . The calculation assumes one week of data taking with a laser repetition rate of 1/60 Hz, integrating over the angular region $|\theta_x| < 20$ mrad and $|\theta_y| < 2.5$ mrad. The curves correspond to: conventional analysis with continuous detection (green dotted line), asymmetry-based analysis with continuous detection (red dashed line), and asymmetry-based analysis with a 13.3 fs time-gated detection window (blue solid line).

θ_y asymmetry [Fig. 3(a)]. Conventional searches are limited by the Poisson fluctuations in a nearly symmetric background rate ($n_b \sim 10^{-4}$ s⁻¹ [49]) [Fig. 3(b)]. In contrast, our scheme reveals a dipolelike structure in the difference of the photon distributions obtained with opposite incident electron spin polarizations [Fig. 3(c)], providing strong discrimination against symmetric backgrounds. Below, we illustrate how this dipolelike structure can lead to improved sensitivity; a more detailed analysis is summarized in Ref. [111]. First, we compare a conventional single-region analysis, based solely on the total count $k = k_1 + k_2$, with an asymmetry-based likelihood that uses the independent Poisson counts k_1 for $\theta_y > 0$ and k_2 for $\theta_y < 0$. For a one-week run at a laser repetition rate of 1/60 Hz, the single-region method yields the bound $g_{ae}g_{\gamma\gamma} \lesssim 3.49 \times 10^{-8}$ GeV⁻¹, whereas the asymmetry-based likelihood improves this to $g_{ae}g_{\gamma\gamma} \lesssim 3.16 \times 10^{-8}$ GeV⁻¹, corresponding to an enhancement of approximately $\sim 10\%$ [Fig. 3(d)]. A further advantage is that the two-region likelihood provides a built-in null test: for a fixed total count $k_1 + k_2$, different partitions (k_1, k_2) yield different likelihood ratios for the asymmetric signal

relative to the symmetric background, enabling an internal check of whether any observed excess follows the expected pattern. Such a self-consistency test is not available when using the total count alone. Furthermore, axion emission in the strong-field laser interaction occurs within an ultrashort ($\sim \mathcal{O}(10)$ fs) time window, allowing for a matched detector gate that suppresses accumulated background by many orders of magnitude. With temporal gating and a 10 petawatt (PW) laser operating at 1/60 Hz, the substantial single-shot yield ($N_a \gtrsim 10^{10} g_{ae}^2$) leads to a projected one-week sensitivity $g_{ae} g_{a\gamma\gamma} \lesssim 1.49 \times 10^{-8} \text{ GeV}^{-1}$, representing an improvement by a factor of ~ 2 over continuous detection, as shown in Fig. 3(d). Therefore, our method is intrinsically more robust against background noise, yielding more reliable limits under realistic experimental conditions.

We recast the photon-based sensitivities as bounds on the axion-electron coupling g_{ae} and compare them with existing laboratory constraints [61,62,135–139]. For parameters characteristic of present 10 PW laser facilities [140,141], one week of data taking gives $g_{ae} \lesssim 3 \times 10^{-5}$, comparable to the electron $g-2$ constraint, although stronger laboratory bounds exist in part of the relevant parameter space,

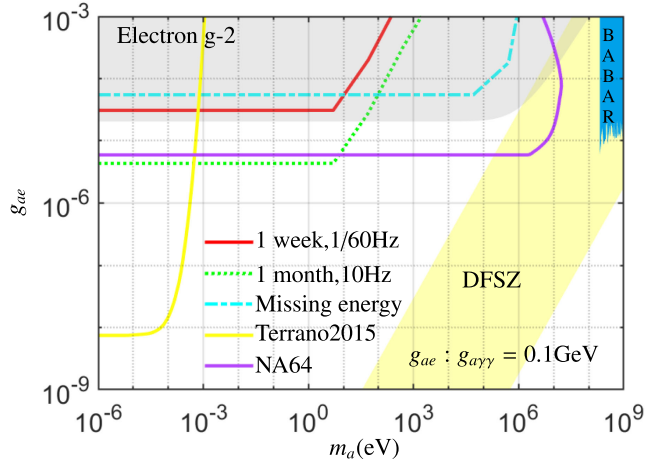


FIG. 4. Projected laboratory bounds on the axion-electron coupling. The red solid (green dotted) curve shows the photon-signal projection for a repetition rate of 1/60 Hz with one week of data taking (10 Hz with one month), adopting the illustrative DFSZ-motivated benchmark $g_{ae} \simeq (0.1 \text{ GeV}) g_{a\gamma\gamma}$ [142,143]. The cyan dot-dashed curve shows the projected exclusion limit from the axion-induced missing-energy signal, assuming a repetition rate of 1/60 Hz and one week of data taking. The yellow solid curve represents the existing constraint from the torsion-pendulum experiment [62], while the purple solid curve denotes the existing limits from the NA64 missing-energy search [135,136]. The blue shaded region shows the exclusion limits derived from BABAR data [137–139]. The gray band indicates the existing constraint from the electron $g-2$ [61], and the yellow band marks the DFSZ axion region [120,144]. All projected limits are quoted at 95% C.L.

e.g., from torsion-pendulum experiments [62] (Fig. 4). This photon-signal projection assumes the illustrative DFSZ-motivated benchmark $g_{ae} \simeq (0.1 \text{ GeV}) g_{a\gamma\gamma}$ [111] and should therefore be interpreted together with existing constraints on both the electron and photon couplings.

The reach can be improved in principle by increasing the laser intensity or repetition rate, or the electron-beam density. As an illustrative future-facility benchmark, a $E \simeq 10^2 \text{ J}$, 10 Hz laser system [145] operating for one month could reach $g_{ae} \lesssim 4.35 \times 10^{-6}$. Nevertheless, the main purpose of this work is not to claim the strongest exclusion limit but to identify a spin-reversible angular-asymmetry signature of strong-field ALP emission, whose sign flips with the initial electron spin and thus provides a distinctive observable against spin-independent backgrounds. In addition to the regenerated-photon signal, an axion-induced missing-energy search in the electron beam provides a complementary detection channel [111], extending the accessible mass range as shown in Fig. 4.

In conclusion, we propose a spin-controlled scheme for axion generation via nonlinear Compton scattering in strong laser fields. By tuning the electron polarization or laser ellipticity, the axion emission angle becomes continuously adjustable, leading to an asymmetric angular distribution. This built-in directional asymmetry not only constitutes a distinctive experimental signature but also acts as a self-referencing null test, enabling robust discrimination against isotropic noise, astrophysical γ -ray backgrounds, and conventional QED signals. Beyond the angular-asymmetry signature, the ultrashort pulsed strong-field interaction provides temporal gating that can help suppress uncorrelated backgrounds and improve the signal-to-noise ratio. Under present 10 PW, low-repetition-rate conditions, the projected sensitivity is comparable to selected existing laboratory bounds, while future improvements in laser repetition rate and electron-beam density could further enhance the accessible reach. By accessing larger values of χ_e , the axion yield from nonlinear Compton scattering could increase by several orders of magnitude compared to the currently investigated regime [131]. In even stronger fields, electron-positron pair-production processes may also become relevant and could provide an additional channel for axion production. Similar ideas may be extended to strong laser-solid and electron-solid interactions, where the much higher electron density could lead to further yield enhancement.

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Data availability. The data that support the findings of this article are not publicly available upon publication because it is not technically feasible and/or the cost of preparing,

depositing, and hosting the data would be prohibitive within the terms of this research project. The data are available from the authors upon reasonable request.

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