

Estimate of multishot laser-induced polarization for high energy electrons

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The use of an intense ultrashort laser pulse to induce electron polarization has been proposed previously in the literature. We used Python to reproduce the local constant crossed-field approximation to calculate the transverse polarization obtained from repeated laser interactions, starting from nonzero initial polarization, as well as the associated energy kicks. We also studied how the collision-geometry affects both the magnitude of χ_e and the effective field asymmetry experienced by the electrons, identifying several possible ways to generate useful asymmetry even when the laser wavelength does not admit pulse lengths that directly reach the plane-wave optimum. We show that, over multiple laser shots, lower values of the quantum efficiency parameter χ_e are associated with higher transverse polarization output but require a greater number of shots to reach saturation. We also show that the mean energy loss scales approximately linearly with χ_e . Based on this framework, we developed a preliminary Better MAD (methodical accelerator design) (BMAD) laser module to simulate the full 6D dynamics in a proposed Future Circular Collider Z pole lattice and to identify the threshold χ_e value above which particle losses begin to occur.

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I. INTRODUCTION

The polarization of high energy electrons in a storage ring is an active area of interest for accelerators. For example, several planned storage rings will require polarized electrons. These include the electron-ion collider (EIC), China's proposed Circular Electron Positron Collider [1], and CERN's proposed Future Circular Collider (FCC-ee), where polarization is important for energy calibration [2]. For the FCC-ee, having a polarized electron beam is important for energy calibration, and the required level is on the order of 5%–10%. This is necessary so that resonant depolarization can be used to observe the energy at which a spin resonance occurs and thus determine the energy of the stored beam [3].

A related long-term motivation is the possibility of extending such methods to positrons. While the present paper focuses only on electrons, a laser-based spin-polarization mechanism could, in principle, also be relevant for future efforts to prepare polarized positron beams, where source options are more limited. In that sense, understanding the beam-dynamics tolerances and

energy-loss constraints of laser-induced polarization for electrons is also a useful preliminary step toward assessing whether analogous ideas could ever be practical for stored positron beams.

The EIC has been designed to investigate the strong force by colliding intense beams of polarized electrons with protons and nuclei to reveal the arrangement of quarks and gluons within the atomic nucleus. Polarization of the electron beam is essential for the EIC, and the current plan is to transport highly polarized electrons from a polarized source through a linear accelerator, and then accelerate them to the final energy in a rapid cycling synchrotron using a design that minimizes the polarization loss due to spin resonances during the acceleration process [4,5]. It is also known that electrons circulating in a storage ring will self-polarize due to the Sokolov–Ternov (ST) effect, a type of radiative spin polarization whereby a particle's change in momentum caused by motion along a bent path results in the ejection of a photon (i.e., synchrotron radiation) and subsequent spin flip. Since transitions to the spin-down state are more likely, up to 92.4% transverse polarization antiparallel to the magnetic field may be achieved over the course of minutes or hours—a relatively slow process [6]. For the EIC, relying on ST self-polarization would be impractical not only because of the long polarization time but also because the collider requires controlled bunch-by-bunch polarization patterns and subsequent spin manipulation to deliver the desired longitudinal polarization at the interaction points.

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The idea of using a circularly polarized electromagnetic wave to polarize stored electrons was discussed by Derbenev, Kondratenko, and Saldin [7]. More recently, it has been shown that an intense laser pulse can be used to polarize electrons within a few femtoseconds [8]. The interaction between an unpolarized beam and an intense laser field results in the transfer of energy to an outgoing gamma-ray photon through nonlinear Compton scattering. As in the ST effect, polarization is obtained through the emission of a photon and a resultant spin flip. Experimental results and theoretical work have confirmed that the degree of polarization due to nonlinear Compton scattering varies with laser parameters as well as initial polarization [9]. Laser-induced polarization for high-energy polarized machines like the EIC and the FCC-ee is worth considering for several reasons: not only could polarization potentially be attained on a much shorter time scale, but a laser-based approach may also provide additional flexibility in how the spin orientation is prepared and manipulated before the beam is delivered to the experimental interaction point. At the EIC, the experimentally relevant polarization is longitudinal at the interaction point, whereas the natural Sokolov–Ternov polarization axis in the arcs is vertical. In the baseline EIC design, spin rotators near the interaction regions are used to rotate the electron spin from vertical in the arcs to longitudinal at the collision point. In that sense, a laser-based scheme could complement the existing spin-manipulation strategy rather than relying solely on self-polarization in the storage ring. The present work therefore explores a complementary mechanism that could supplement, but not replace, the baseline polarized source and spin-rotation design. For energy calibration purposes in the FCC-ee, a laser approach could potentially target a single pilot bunch, avoiding polarization of the colliding bunches. Certain experiments are sensitive to beam polarization, which can bias results [3]. Also, at the FCC-ee Z pole, the Sokolov–Ternov polarization time is extremely long (about 256 h, or roughly 14–29 h to reach 5%–10% polarization), so a laser-based approach could in principle help avoid the need for wigglers that are otherwise planned to shorten the polarization time, thereby potentially reducing cost and complications (i.e., lowering beam aperture, impedance, and radiation damage).

The magnitude of polarization (5%) reported in these studies [10] is much smaller than the 70% average polarization expected for the EIC program, although it is in the range relevant for the FCC-ee. This motivated us to study the repeated application of a laser pulse to a circulating high-energy beam and the possibility of building up sizable average polarization over many shots. In this paper, we use the method developed in Ref. [10], based on the local constant crossed-field approximation (LCFA), to study the viability of polarizing high energy electron beams with a high-power laser in a machine such as the FCC-ee or EIC and to develop an initial qualitative

understanding of the polarization response. Typically, the LCFA is used in cases where ξ , the nonlinearity parameter, is much greater than 1. Here the nonlinearity parameter is given by:

$$\xi = 8.55 \sqrt{\lambda^2 [\mu\text{m}] I \left[10^{20} \frac{\text{W}}{\text{cm}^2} \right]}, \quad (1)$$

where λ is the wavelength and I is the intensity. For example, Ref. [10] showed that LCFA calculations at $\xi = 20$ and $\xi \approx 10$ reproduce exact QED results to within about 5% and 10%, respectively, for a quantum efficiency parameter χ_e between 10^{-3} and 10. The quantum efficiency parameter, χ_e , is roughly a measure of the field strength in the electron’s rest frame. In traditional radiative polarization, where the electromagnetic field is weak, $\chi_e = e \sqrt{(F^{\mu\nu} p_\nu)^2} / m^3 \ll 1$ [11], where e is the fundamental electric charge, F is the electromagnetic tensor, p is the momentum, and m is the electron mass. The parameter χ_e can also be written as the product $\chi_e = b\xi$, with

$$b = \frac{k^\mu p_\mu}{m^2},$$

where k^μ and p^μ are the laser photon and electron four-momenta, respectively, and $k^\mu p_\mu$ denotes the Lorentz-invariant four-vector product.

In this paper, we present results for laser intensities compatible with these values of ξ . Unless otherwise noted, laser parameters such as wavelength, frequency, and pulse duration are given in the laboratory frame. We consider the types of lasers, their pulse energies, focus sizes, repetition rates, and the overall geometry of the laser-electron interaction that could achieve these parameters. We also examine the relative positions and overlap of the laser and electron beam in an accelerator, including the mismatch between the laser spot size and the electron bunch dimensions. In addition, we investigate the energy kick delivered to an electron by such lasers and show that it is a function of the quantum efficiency parameter χ_e . Both the spin and energy-kick effects due to the laser interaction in the LCFA regime have been implemented in the general accelerator platform BMAD [12] via a laser-interaction module that simulates a periodic interaction between an electron bunch and a laser, including the kinematics of scanning along the longitudinal extent of a given electron bunch in an accelerator rf bucket. Using this code, we simulated an FCC-ee Z pole electron beam in an rf bucket and studied the effects on spin, energy, and beam loss over repeated laser shots while scanning across the bunch.

II. POLARIZATION MODEL

Two methods have been previously proposed for calculating the degree of polarization in final state electrons.

The first uses a density matrix to calculate the quantum state of the electrons. The alternative method, known as the LCFA, estimates the results of the density matrix for a linearly polarized laser. According to the LCFA, the electron polarization vector, describing the degree of polarization in the transverse plane, can be represented by:

$$\Xi_y = \frac{\int d\tau dt \left[\Xi_y \text{Ai}_1(z(t, \tau)) + \Xi_y \frac{2\text{Ai}'(z(t, \tau))}{z(t, \tau)} + \frac{t}{1-t} \frac{\text{Ai}(z(t, \tau))}{\sqrt{z(t, \tau)}} \text{sgn}(\dot{h}(\tau)) \right]}{\int d\tau dt \left[\text{Ai}_1(z(t, \tau)) + \left(1 + \frac{t^2}{2(1-t)}\right) \frac{2\text{Ai}'(z(t, \tau))}{z(t, \tau)} + \Xi_y \text{sgn}(\dot{h}(\tau)) t \frac{\text{Ai}(z(t, \tau))}{\sqrt{z(t, \tau)}} \right]}. \quad (2)$$

Here, Ξ_y describes the initial polarization of the beam in the vertical plane in a coordinate system with the storage ring and laser, as shown in Fig. 1. The integration variables t and τ are the fractional momentum transfer and proper time, respectively. The argument z of the Airy functions $\text{Ai}(z)$, their derivatives $\text{Ai}'(z)$, and their integrals $\text{Ai}_1(z) = \int_z^\infty dx \text{Ai}(x)$ is given by $z = z(t, \tau) = \left[\frac{t}{1-t} \frac{1}{|\dot{h}(\tau)|\chi_e} \right]^{\frac{2}{3}}$, where the pulse shape function $h(\phi)$ describes the vector potential and can be written as:

$$h(\phi) = \cos(\phi + \phi_{\text{CE}}) \cos^2\left(\frac{\pi\phi}{2\Delta\phi}\right) \Theta(\Delta\phi - |\phi|), \quad (3)$$

where Θ is the Heaviside function. Thus, the time derivative $\dot{h}$ describes the shape of the electric field. When the initial polarization is zero (i.e., $\Xi_y = 0$), the last term in Eq. (2) is responsible for the buildup of polarization. In general, this buildup is related to the shape of $\dot{h}$. The shape and asymmetry of a magnetic field produced by a laser are limited due to the inability to produce “unipolar” fields (where $h(-\infty) \neq h(\infty)$). In the semiclassical limit, the polarization is proportional to the asymmetry $\mathcal{A}$ of $\dot{h}$:

$$\Xi_y = -\frac{3\chi_e}{10} \frac{\int d\tau |\dot{h}(\tau)| \dot{h}(\tau)}{\int d\tau |\dot{h}(\tau)|} \equiv \frac{3\chi_e}{10} \mathcal{A}. \quad (4)$$

The asymmetry is maximal at $|\mathcal{A}| \approx 0.17 \left(\frac{3.27\lambda}{2\pi\tau_c} \right)^4$ when the carrier-envelope phase $\phi_{\text{CE}} = \pi/2$ and the pulse duration is $\Delta\phi = 3.27$, and it effectively becomes zero for $\Delta\phi > 8$ [10]. The parameter $\Delta\phi$ determines the pulse length for a given wavelength according to $\Delta\phi \approx \frac{\pi c}{0.88\lambda} \tau_{\text{FWHM}}$. Thus, to reach $\Delta\phi = 3.27$ for an 800 nm laser, one needs a pulse of about 2.44 fs, while for a 9.20 μm laser, one would need a pulse length of about 28 fs. For Ti:sapphire lasers at 800 nm, such pulse lengths have been demonstrated; however, for CO₂ lasers at 9.20 μm , 28 fs is currently out of reach. In such cases, there may be other possible approaches to achieving useful pulse asymmetries, for example, quantum-cascade-laser-seeded schemes or plasma shutters; see Refs. [13,14]. In the present problem, we must also consider the orientation of the electron-laser collision, since it affects both the field

asymmetry experienced by the electron and the magnitude of χ_e .

III. GEOMETRY AND EFFECTIVE ASYMMETRY

Figure 1 shows the layout of how the laser would interact with the electron bunch.

Here we denote θ as the collision angle. The quantum efficiency parameter in a plane-wave laser field is

$$\chi_e \simeq \xi \frac{k^\mu p_\mu}{m^2} = \xi \frac{\gamma\omega}{m} (1 - \beta \cos \theta), \quad (5)$$

where p^μ is the electron four-momentum, k^μ is the laser four-wavevector, ω is the laser photon energy, and θ is the angle between the electron momentum and the laser propagation direction. This shows explicitly that χ_e can be controlled by the collision geometry. In the head-on case ($\theta = \pi$),

$$\chi_e \simeq \xi \frac{\gamma\omega}{m} (1 + \beta) \approx 2\xi \frac{\gamma\omega}{m}, \quad (6)$$

whereas in the copropagating collinear limit ($\theta = 0$),

$$\chi_{e,\text{min}} = \xi \frac{\gamma\omega}{m} (1 - \beta) = \xi \frac{\omega/m}{\gamma(1 + \beta)} \approx \xi \frac{\omega}{2\gamma m}. \quad (7)$$

Thus, for fixed beam energy, the collinear limit does not make χ_e exactly vanish, but it does reduce it to a very small minimum value that scales as $1/\gamma$. For $\xi = 10$, this minimum is already extremely small; for all cases

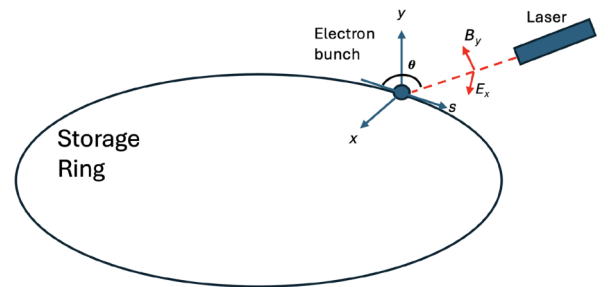


FIG. 1. Layout with a linearly polarized laser incident at angle θ on an electron bunch stored in an accelerator.

considered in this paper, its minimum value would still be greater than 10^{-9} .

The geometric configuration must be discussed somewhat carefully because it influences the polarization dynamics in two different ways. First, it sets the magnitude of χ_e through Eq. (5). Second, once a finite laser waist and a finite electron bunch are considered, it determines which portion of the laser pulse is actually sampled by a given electron. The latter effect can be described by a geometric weighting function G , which may generate an *effective* asymmetry in the sampled field even when the underlying pulse is not strongly asymmetric in the ideal plane-wave limit.

To derive G , consider the electron trajectory:

$$\mathbf{r}_e(t) = \mathbf{r}_{e0} + \beta ct \hat{\mathbf{z}}, \quad (8)$$

and let the laser propagate in the x - z plane with unit vector:

$$\hat{\mathbf{n}} = (\sin \theta, 0, \cos \theta). \quad (9)$$

If the laser axis passes the bunch with offsets (x_0, y_0, z_0) relative to the chosen electron trajectory, then the squared transverse distance from the electron to the laser axis is

$$\rho^2(t) = y_0^2 + [x_0 \cos \theta - (z_0 + \beta ct) \sin \theta]^2. \quad (10)$$

Neglecting longitudinal diffraction effects over the short interaction region, a Gaussian beam gives the overlap factor:

$$G(t) = \exp \left[-\frac{\rho^2(t)}{w_0^2} \right]. \quad (11)$$

The sampled field can then be written as:

$$E_{\text{eff}}(t) \propto \dot{h}(\phi(t) - \phi_{\text{off}})G(t), \quad (12)$$

where ϕ_{off} is a timing offset between the bunch and the laser pulse. In the centered case $x_0 = y_0 = z_0 = 0$, Eq. (11) reduces to:

$$G(t) = \exp \left[-\frac{(\beta ct \sin \theta)^2}{w_0^2} \right]. \quad (13)$$

The phase sampled along the electron trajectory evolves as:

$$\phi(t) = \omega(1 - \beta \cos \theta)t, \quad (14)$$

so while oblique incidence suppresses χ_e through the factor $(1 - \beta \cos \theta)$, it does not by itself change the intrinsic pulse shape as a function of phase. What it does change, once a finite waist is included, is the duration

for which the electron remains inside the focal region. Combining Eqs. (13) and (14) gives a characteristic sampled phase width:

$$\sigma_\phi = \frac{\omega(1 - \beta \cos \theta)w_0}{\beta c \sin \theta} \approx \frac{2\pi w_0}{\lambda} \tan \frac{\theta}{2}, \quad (15)$$

where the last form assumes $\beta \simeq 1$. Correspondingly, a characteristic phase interval sampled during the crossing is

$$\Delta\phi_\perp \sim 2\sigma_\phi = \frac{4\pi w_0}{\lambda} \frac{1 - \beta \cos \theta}{\beta \sin \theta} \approx \frac{4\pi w_0}{\lambda} \tan \frac{\theta}{2}. \quad (16)$$

In the centered Gaussian model, the field history sampled by the electron can therefore be represented in phase space by an effective pulse derivative:

$$\dot{h}_{\text{eff}}(\phi) = G(\phi)\dot{h}(\phi), \quad G(\phi) = \exp \left[-\left(\frac{\phi}{\sigma_\phi} \right)^2 \right], \quad (17)$$

which reduces to the usual plane-wave pulse shape when $G(\phi) = 1$. This is the natural quantity to insert into the LCFA integrator when geometry, timing offsets, or finite overlap are included.

In the ideal plane-wave pulse model, $\Delta\phi = 3.27$ identifies the pulse-duration parameter that maximizes the intrinsic asymmetry. Once finite-waist and collision-geometry effects are included, however, the relevant quantity is the field history actually sampled by the electron, and that need not coincide with the plane-wave optimum. Geometry can therefore produce an effective asymmetry larger than the plane-wave value even when the sampled phase interval differs substantially from $\Delta\phi = 3.27$.

This motivates a generalized effective-asymmetry model,

$$\mathcal{A}_{\text{eff}} = -\frac{\int d\tau |\dot{h}_{\text{eff}}(\tau)| \dot{h}_{\text{eff}}(\tau)}{\int d\tau |\dot{h}_{\text{eff}}(\tau)|}, \quad (18)$$

which reduces to Eq. (4) when $\dot{h}_{\text{eff}}(\tau) = \dot{h}(\tau)$. For the pulse model used in this paper with $\Delta\phi = 3.27$ and $\phi_{\text{CE}} = \pi/2$, Eq. (4) gives the plane-wave value $\mathcal{A} \simeq 0.174$. By contrast, the generalized model of Eq. (18) permits substantially larger effective asymmetries whenever the geometric weighting preferentially selects one sign of the laser magnetic field over the other. Thus, in the geometry-gated case, the relevant asymmetry need not be limited by the plane-wave value $\mathcal{A} \approx 0.17$, although the larger asymmetry is generally accompanied by a reduced total interaction strength. At the same time, excessively strong geometric clipping can shorten the sampled interaction to the point that finite-pulse and field-gradient effects become important. In that regime, the sampled field may no longer be well represented by a locally constant crossed field over the relevant formation interval, so the LCFA

description should be treated with caution even if the peak field still corresponds nominally to $\xi \sim 10$.

To characterize the smoothness of the sampled waveform relative to the benchmark pulse used throughout this paper, it is useful to compare the geometry-gated history to the reference pulse of Eq. (3) with $\Delta\phi = 3.27$ and $\phi_{\text{CE}} = \pi/2$, for which earlier work found good LCFA agreement at $\xi \gtrsim 10$ [10]. Thus we can use

$$F(\phi) = \dot{h}_{\text{eff}}(\phi), \quad (19)$$

to characterize how rapidly the sampled waveform varies in phase, and thus define the rms phase-bandwidth,

$$\mathcal{B}[F] = \left(\frac{\int d\phi |F'(\phi)|^2}{\int d\phi |F(\phi)|^2} \right)^{1/2}, \quad (20)$$

and the corresponding geometry-distortion ratio:

$$R_{\text{geom}} = \frac{\mathcal{B}[F]}{\mathcal{B}[F_{\text{ref}}]}, \quad (21)$$

where $F_{\text{ref}}(\phi) = \dot{h}(\phi, \Delta\phi = 3.27, \phi_{\text{CE}} = \pi/2)$. By construction, $R_{\text{geom}} < 1$ means that the geometric weighting has made the sampled waveform smoother in phase than the benchmark pulse, while $R_{\text{geom}} > 1$ means that the geometry has introduced additional rapid variation through clipping or distortion. This definition captures cases in which a long intrinsic pulse can actually become smoother after moderate Gaussian gating because the weighting suppresses the sharp outer structure of the original waveform. Accordingly, R_{geom} should be interpreted here as a benchmark-relative waveform-smoothness diagnostic rather than as a standalone proof of LCFA validity. The separate requirement that the field itself remain in the nominal LCFA regime is still represented by the usual condition $\xi \sim 10$.

Equations (10), (11), (15), (17), (18), (20), and (21) also show that there are several distinct experimental knobs that can be used to change the field history sampled by the electrons: (i) Carrier-envelope phase and pulse shaping. This is the intrinsic method emphasized in Ref. [10]. It changes the waveform itself and therefore directly modifies the asymmetry parameter $\mathcal{A}$ in Eq. (4). (ii) Collision angle. For a finite spot size, oblique incidence limits the time that an electron spends inside the focal region. This can truncate the interaction to only part of the laser pulse and therefore convert a longer pulse into a more asymmetric effective field sample, while simultaneously reducing χ_e through Eq. (5). (iii) Timing offset. If the laser arrives early or late relative to the bunch center, then the bunch no longer samples the laser envelope symmetrically in time. In practice, this means that different slices of the bunch may interact preferentially with the leading or trailing part of the pulse. This is likely the most direct

geometric way to generate an effective asymmetry without changing the laser waveform itself. (iv) Transverse offset in the crossing plane. An offset in the plane defined by the beam direction and the laser propagation direction changes the spatial part of the overlap in a way that can map into a timing bias when the collision is oblique. This can again weight the front and back of the pulse differently along the bunch trajectory. (v) Vertical offset y_0 . A pure offset orthogonal to the crossing plane mainly reduces the peak field sampled by the electron through the Gaussian overlap factor in Eq. (12). By itself, this does not create a strong sign asymmetry in the sampled waveform; instead, it primarily lowers the effective field strength and therefore lowers χ_e . However, in a realistic machine, it can still matter because it changes the bunch-to-bunch and slice-to-slice weighting of the interaction and could combine with timing or angular effects to modify the net polarization buildup.

For the purposes of accelerator implementation, it is useful to distinguish between these mechanisms. Carrier-envelope phase control and pulse shaping modify the laser pulse itself. In contrast, collision angle, timing offset, and transverse offsets do not necessarily make the laser more asymmetric in the strict plane-wave sense; rather, they change which portion of the pulse is sampled by the electron bunch. In an accelerator, this latter class of effects may be easier to tune shot-to-shot and can be incorporated naturally into tracking studies by scanning the laser timing and spatial overlap relative to the bunch. The main cases are illustrated schematically in Fig. 2.

IV. METHODOLOGY

A Python program [15] was initially used to find final polarization amplitudes given input values of $\Delta\phi$, ϕ_{CE} , and χ_e . The LCFA, Eq. (2), was used to create a function taking input values of the quantum efficiency parameter χ_e and returning values of transverse polarization Ξ_y . This was accomplished by splitting Eq. (2) into four different integrals and then integrating them separately. Some of the integrands diverge as $t^{-2/3}$, as t approaches zero, but this can be avoided with the substitution $u = t^{1/3}$. In order to confirm the validity of the code, values for transverse polarization Ξ_y were plotted as a function of the quantum efficiency parameter χ_e for zero initial polarization. The results were consistent with those presented in existing literature [10], with a peak of 0.006 occurring at $\chi_e = 0.5$ (Fig. 3).

V. RESULTS OF REPEATED-SHOT LCFA STUDY

To observe the effect of multiple laser shots, the calculated transverse polarization resulting from each laser shot was used as the starting value for the subsequent shot over 1000 laser shots, assuming an initial polarization of 0 and χ_e values of 0.01 [Fig. 4(a)], 0.1–0.7 [Fig. 4(b)], and

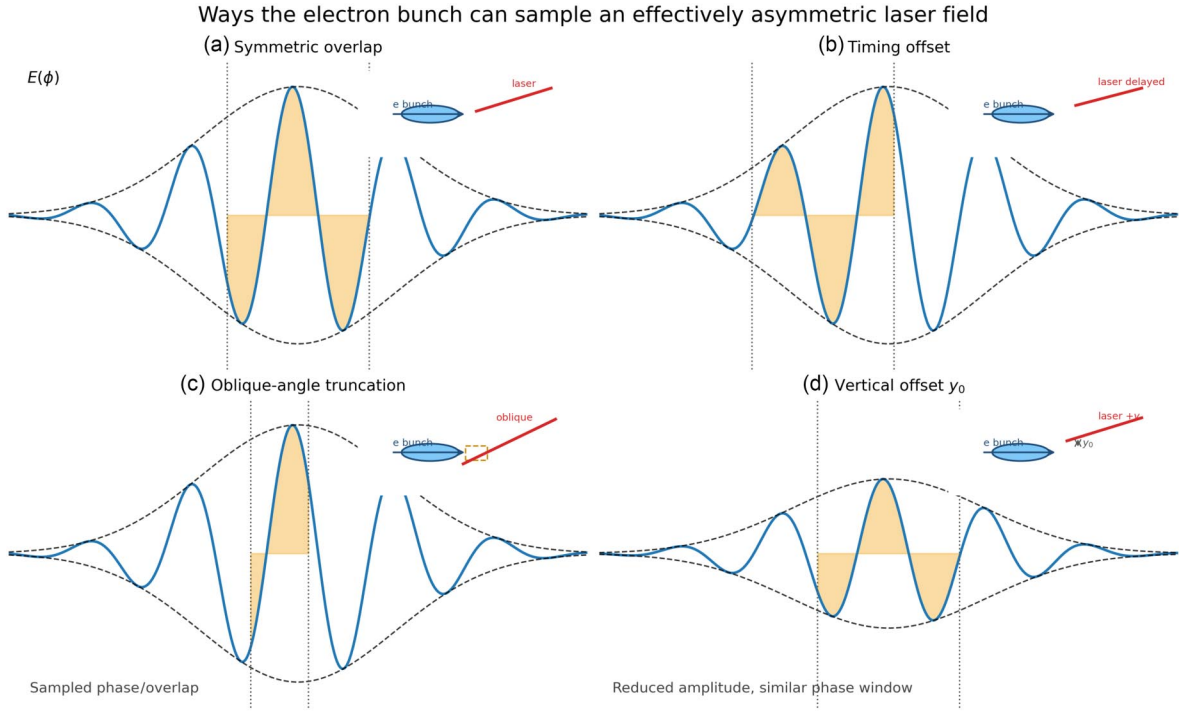


FIG. 2. Schematic examples of how the electron bunch can sample an effectively asymmetric laser field. Panel (a) shows a nominal centered overlap, for which the sampled waveform remains nearly symmetric. Panel (b) shows a timing offset, which biases the interaction toward the leading or trailing part of the laser pulse. Panel (c) shows truncation caused by oblique incidence and finite waist, which limits the phase interval sampled by the electron while also reducing χ_e . Panel (d) shows a vertical offset y_0 , which mainly reduces the sampled field amplitude and therefore weakens the interaction without strongly changing the waveform symmetry by itself. The small insets indicate the corresponding laser-bunch overlap geometries.

1.6–30 [Fig. 4(c)]. For $\chi_e = 0.7$, the polarization plateaus at about 30% after 200 shots, while at 0.2 and 0.1, it can reach about 70% and 80%, respectively.

The repeated-shot results show that lower values of the quantum efficiency parameter χ_e lead to larger asymptotic polarization, but at the cost of requiring many more shots to approach saturation. This trend is broadly consistent with the behavior discussed in Ref. [8], where larger values

of χ_e correspond to lower equilibrium polarization. In the present repeated-shot study, the response is nontrivial because the polarization after each shot becomes the initial condition for the next one. At larger χ_e , the spin-flip probability increases, but the net directional asymmetry of the spin-flip process appears to decrease. A plausible explanation is that the buildup is tied to the envelope asymmetry of the ultrashort laser pulse: polarization generated in one part of the pulse is only partially canceled by later parts, and this balance changes as χ_e increases. Once the initial polarization becomes nonzero, the terms proportional to Ξ_y in Eq. (2) also contribute and modify the subsequent evolution.

Assuming the trend observed in Fig. 4 remains valid in the machine context, lower quantum efficiency values are preferable because they are associated with weaker field strengths, higher asymptotic polarization, smaller energy kicks, and therefore reduced beam loss. Typically, weaker laser fields are also more compatible with higher repetition rates, so more shots can be delivered per second to compensate for the slower buildup.

VI. ENERGY LOSS ESTIMATES

The same LCFA kernel used to compute the spin map also provides the mean fractional energy loss per laser

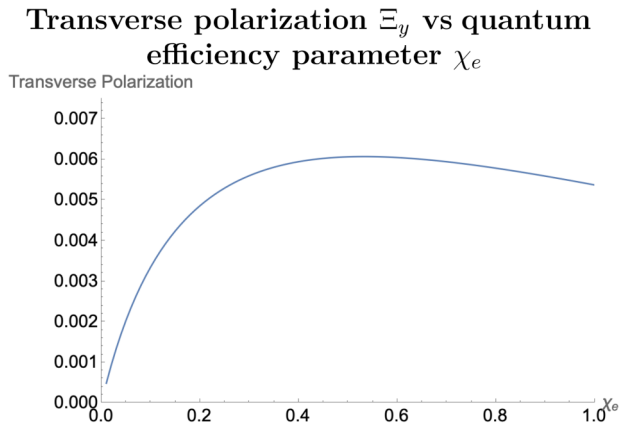


FIG. 3. Transverse polarization Ξ_y as a function of quantum efficiency parameter χ_e . The results are consistent with those presented in Ref. [10], with χ_e peaking at 0.5.

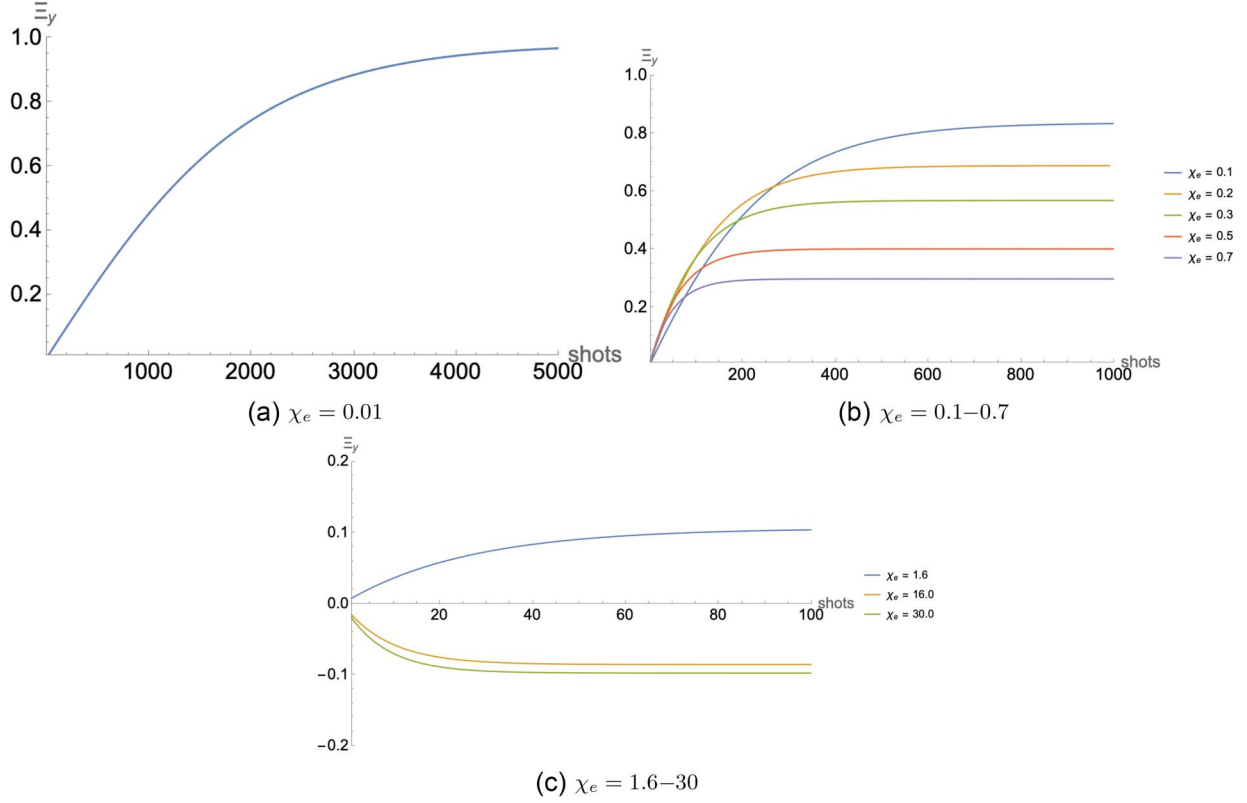
Transverse Polarization Ξ_y vs Number of Shots

FIG. 4. Number of shots vs transverse polarization for $\chi_e = 0.01$ (a), $0.1-0.7$ (b), and $1.6-30$ (c). Although polarization plateaus at greater values when χ_e is low, the number of shots, and by extension the time, required to achieve maximal polarization is longer. At much higher χ_e values, the polarization direction eventually flips and appears to settle at approximately -10% , with higher χ_e , corresponding only to fewer shots needed to reach the maximum polarization value.

interaction for a single electron. For an incoming polarization Ξ_y and quantum parameter χ_e , the expectation value of the fractional energy carried by emitted photons is

$$\left\langle \frac{\Delta E}{E} \right\rangle = \frac{\int d\tau dt t W(t, \tau; \Xi_y, \chi_e)}{\int d\tau dt W(t, \tau; \Xi_y, \chi_e)}, \quad (22)$$

where the weight W is exactly the denominator integrand of Eq. (2):

$$W = \text{Ai}_1(z) + \left(1 + \frac{t^2}{2(1-t)}\right) \frac{2\text{Ai}'(z)}{z} + \Xi_y \text{sgn}(\dot{h}(\tau)) t \frac{\text{Ai}(z)}{\sqrt{z}}, \quad (23)$$

with the same definitions of t , τ , $z(t, \tau)$, and pulse shape $h(\phi)$ as in the polarization formula. In practice, the integrals are evaluated by splitting W into the three Airy terms shown above and reusing the numerical pieces computed for the spin map. Because t is strictly positive, this LCFA expression only yields an energy

loss; a net gain would require stimulated absorption or external acceleration and is not contained in the emission kernel.

An estimate of the mean energy loss as a function of χ_e is shown in Fig. 5. The curve is shown out to 0.8% , which is at the edge of what a typical rf bucket could contain. This sets only a rough mean-loss limit on the tolerable quantum efficiency parameter of $\chi_e < 0.008$. The result also shows a nearly one-to-one relationship between χ_e and $\frac{\Delta E}{E}$ over the plotted interval.

The beam-loss issue cannot be assessed only from the mean loss. In the low- χ_e limit, the strong-field emission spectrum approaches the familiar synchrotron-radiation form rather than a monochromatic or single-edge recoil spectrum. The characteristic photon energy is

$$u_c = \frac{3}{2} \chi_e E. \quad (24)$$

Here u denotes the emitted photon energy, related to the LCFA integration variable by $u \simeq tE$ for ultrarelativistic forward emission. At the operating point $\chi_e = 0.002$,

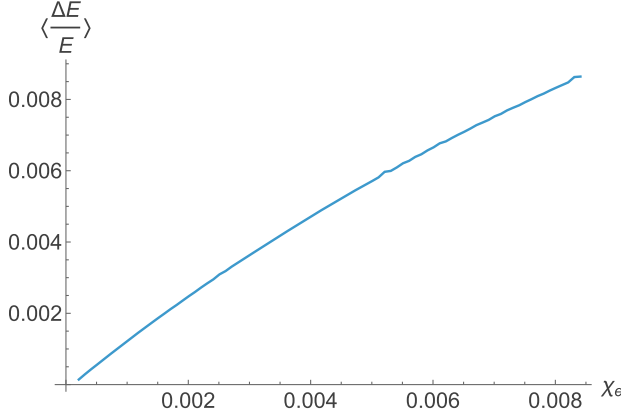


FIG. 5. Energy Loss $\langle \frac{\Delta E}{E} \rangle$ vs quantum efficiency parameter χ_e .

one has $u_c/E \simeq 0.003$. The high-energy tail is expected to fall approximately as:

$$\frac{dN_\gamma}{du} \propto \exp\left(-\frac{u}{u_c}\right), \quad (25)$$

up to algebraic prefactors, as follows from the large-argument asymptotic behavior of the Airy functions in Eq. (23). Thus rare photons with energies several times u_c are exponentially suppressed but are part of the LCFA spectrum. For example, at $\chi_e = 0.002$, a photon carrying 0.8% of the electron energy corresponds to $u/u_c \simeq 2.7$ and is in the tail of the spectrum. This tail is especially relevant for stored beams because a single large recoil can move a particle close to, or outside, the rf bucket even when the mean loss is tolerable.

The same numbers also show that the number of emitted photons per laser interaction is of order unity at the benchmark point. A simple estimate of the photon multiplicity is

$$N_\gamma \sim \frac{\langle \Delta E \rangle}{\langle u \rangle}, \quad (26)$$

where $\langle u \rangle$ is the mean photon energy in the emission spectrum. Since Fig. 5 gives $\langle \Delta E/E \rangle \simeq 0.2\%$ at $\chi_e = 0.002$, while $u_c/E \simeq 0.3\%$, the total mean radiated energy per interaction is comparable to the characteristic energy scale of a single photon. The synchrotron spectrum is broad, so $\langle u \rangle$ is an order-unity fraction of u_c rather than exactly equal to u_c ; nevertheless, the ratio above gives N_γ of order one, not a large many-photon number. This estimate is sufficient to show that photon-number fluctuations and the high-energy tail can matter. The present BMAD implementation applies the mean LCFA recoil from the precomputed lookup table rather than sampling individual photon emissions. Therefore, the results should be interpreted as a mean-kick beam-survival estimate; stochastic photon-number and photon-energy fluctuations,

which may enhance energy-spread growth and rare-particle loss, will be treated in future work using the full density-matrix framework.

VII. BMAD LASER MODULE

Based on our Python LCFA code, we implemented a module in BMAD [12] and utilized it in the long-term tracking code to simulate the interaction between a laser pulse and an electron beam in an rf bucket. This code includes the FCC-ee Z pole lattice optics, spin-orbit tracking, synchrotron motion, rf acceleration, ordinary synchrotron-radiation damping and quantum excitation, and aperture checks. Our added module simulates the timing of a laser that repeatedly scans across the length of a bunch and models the laser-overlap envelope using a user-defined pulse width and timing scan step size. It applies both the energy kick and the spin kick by interpolation from a precomputed LCFA lookup table. The present simulations therefore include stochastic synchrotron-radiation from the ring but not stochastic laser recoil. They also do not yet constitute a dedicated dynamic-aperture study with field errors and nonlinear aperture scans, and they do not apply the geometric effects of Eq. (17), which will be deferred to later work.

Thus, this simulation helps us estimate the tolerable mean-kick limits of χ_e for a stored bunch in a sample lattice, as well as the polarization growth rate. Using the FCC-ee Z pole lattice, we simulated 2214 macroparticles distributed in a 90 MV rf bucket at 45.6 GeV, with laser interactions in the range from $\chi_e = 0.006$ down to $\chi_e = 0.002$. We found that particle loss begins to occur for $\chi_e > 0.002$. Thus, even before adding stochastic laser recoil, the full 6D mean-kick simulations set a tighter bound for the FCC-ee Z pole lattice than the simpler mean-loss estimate of $\chi_e = 0.008$. At $\chi_e = 0.002$, the polarization builds up to 0.69% over 100 000 turns (see Fig. 6), with no particle loss: all 2214 macroparticles remain in the tracked bunch through the final turn. In Fig. 7, the initial and final longitudinal distributions are shown for the $\chi_e = 0.002$ case. Some emittance growth is visible, and the tracked particles remain within the rf separatrix, which is also plotted.

Since the growth is nearly linear over the tracked interval, one may form a rough scaling estimate using the observed rate of $\approx 6.9 \times 10^{-6}\%/turn$. Under the highly idealized assumption that this rate continues unchanged, 5%–10% polarization would correspond to about 220–440 s at a repetition rate matched to the FCC-ee revolution period. This estimate should be interpreted cautiously: it is only an order-of-magnitude extrapolation from the present 100 000-turn, 2214-particle, mean-kick tracking study and does not by itself demonstrate an immediately practical operating scenario. In particular, a laser system capable of delivering the required combination of repetition rate and peak field within the nominal LCFA regime is not

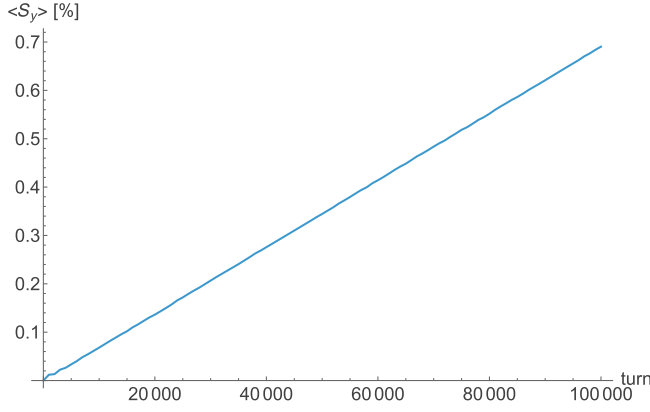


FIG. 6. Average vertical polarization per turn for an FCC-ee Z pole bunch stored at 45.6 GeV, with bunch length 18 ps and $\frac{\Delta E}{E} = 3.9 \times 10^{-4}$, for a laser sweeping over ± 20 ps, with a pulse width of 35 fs, corresponding to a quantum efficiency parameter of $\chi_e = 0.002$.

presently available. If the beam were instead struck only every third or fourth turn, the same linear scaling would suggest a substantially longer buildup time of order 1400s for 10% polarization in the FCC-ee Z pole case. Similar extrapolations for the EIC cases are even more tentative, since the relevant dynamic aperture, damping, rf acceptance, energy-compensation scheme, and polarization lifetime are machine-dependent and were not modeled here.

The average laser energy loss also raises an operational issue for any storage-ring application. At $\chi_e = 0.002$, the mean loss per laser interaction is about 0.2%, larger than the ordinary FCC-ee Z pole synchrotron-radiation loss per turn. If only selected FCC-ee pilot bunches are irradiated, this loss would have to be recovered with bunch-selective rf manipulation, an energy-recovery waiting period, or a

laser pulse pattern slow enough that synchrotron damping and rf focusing can restore the bunch before the next effective kick. For the EIC, where one would normally need all physics bunches polarized, such bunch-selective compensation is much less natural; this is one reason the EIC discussion in this paper should be viewed only as exploratory.

VIII. LASER REQUIREMENTS AND CANDIDATE PLATFORMS

To reach the $\xi \approx 10$ range, lasers with pulse energies above 0.04 J are needed. Several candidates are shown in Table I. The estimated ξ_{peak} values are computed using a conservative, physically achievable focusing benchmark, $w_0 = \lambda$, rather than an idealized diffraction-floor spot size. This keeps the Yb:YAG entries in the comparison while making clear that their relevance is primarily their high repetition rate and moderate- ξ operation, rather than straightforward access to the $\xi \gtrsim 10$ regime.

Table II summarizes the maximum values of the parameter b for representative laser wavelengths and electron-beam energies in the head-on collision geometry.

As mentioned earlier, energy-loss considerations indicate that we should restrict attention to $\chi_e \lesssim 0.002$. For the EIC cases and the FCC-ee Z pole case considered here, the laser-selection problem is therefore not simply to maximize ξ , but rather to choose a system that can simultaneously satisfy three conditions: remain close to the LCFA validity threshold $\xi \lesssim 10$, reproduce the BMAD operating point $\chi_e = 0.002$, and provide a sufficiently asymmetric sampled field history to generate a useful polarization signal.

The plane-wave result $\Delta\phi = 3.27$ remains a useful benchmark because at 800 nm it corresponds to a pulse

Longitudinal distribution before and after tracking with laser interaction at $\chi_e = 0.002$

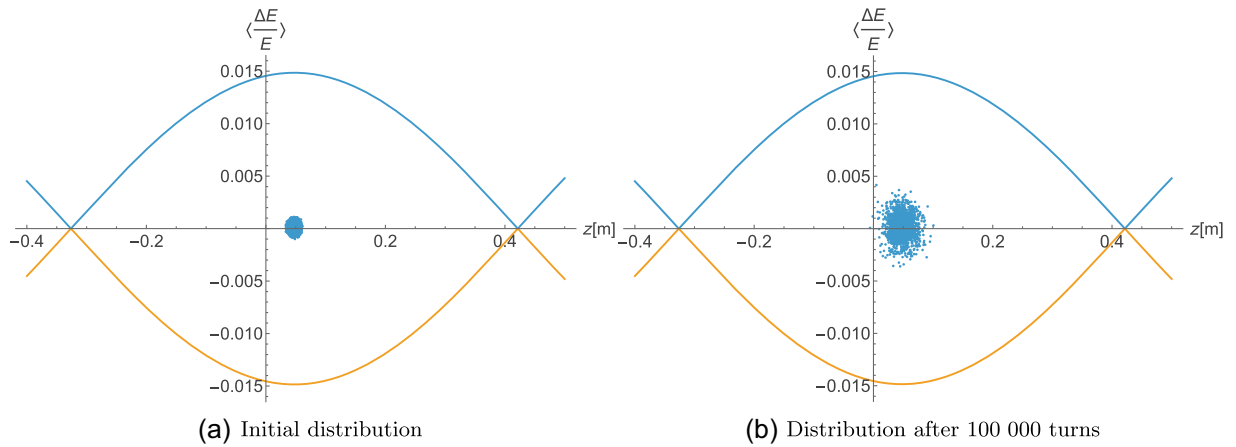


FIG. 7. Longitudinal distribution before and after tracking with laser interactions over 100 000 turns at $\chi_e = 0.002$ for an FCC-ee Z pole bunch stored at 45.6 GeV, with bunch length 18 ps and $\frac{\Delta E}{E} = 3.9 \times 10^{-4}$, for a laser sweeping over ± 20 ps, with a pulse width of 35 fs. All 2214 macroparticles were retained through the final turn. The rf separatrix for 90 MV, including radiative loss, is also shown.

TABLE I. Representative ultrafast laser platforms relevant to strong-field spin dynamics. To avoid implying idealized near-diffraction-limited focusing, the spot-size column uses a conservative free-space benchmark, $w_0 = \lambda$, and the corresponding ξ_{peak} values are estimated from the listed pulse parameters. Under this benchmark, Ti:sapphire chirped pulse amplification (CPA), few-cycle near-infrared (NIR) optical parametric chirped pulse amplification (OPCPA), and long-wavelength infrared (LWIR) CO₂ CPA reach $\xi \gtrsim 10$, while current thin-disk and cryogenic Yb:YAG systems remain in the moderate- ξ regime.

Laser platform	λ (μm)	Representative pulse energy (J)	Peak rep. rate (Hz)	Pulse duration	Assumed $w_0 = \lambda$ (μm)	Estimated ξ_{peak} at assumed focus
Ti:Sapphire (CPA) [16]	0.80	0.75	100	30 fs	0.80	33.0
Ti:Sapphire (postcompressed)	0.80	0.0063–0.0070	10^3	3.8–5 fs	0.80	7.4–9.0
Yb:YAG (thin-disk) [17]	1.03	0.72	10^3	920 fs	1.03	5.85
Yb:YAG (cryo CPA/slab) [18]	1.03	1.10	10^3	4.5 ps	1.03	3.27
OPCPA (few-cycle, NIR) [19]	0.88	0.053	10^3	9 fs	0.88	16.0
OPCPA (mid-IR) [20]	3.10	0.0011	10^4	58 fs	3.10	0.91
CO ₂ (LWIR CPA) [21]	9.20	10	0.05	2 ps	9.20	14.8

TABLE II. Dimensionless values of the head-on collision parameter $b = 2\gamma\omega_0/m$ for representative laser wavelengths and electron-beam energies relevant to the EIC and the FCC-ee Z pole.

Wavelength (μm)	EIC			FCC-ee
	5 GeV	10 GeV	18 GeV	45.6 GeV
0.80	0.0594	0.1187	0.2137	0.540
1.03	0.0461	0.0922	0.1660	0.420
9.20	0.0052	0.0103	0.0186	0.047

duration of about 2.44 fs, close to the regime where the intrinsic asymmetry is largest. However, once geometry is included, this should not be interpreted as a strict requirement that the sampled phase interval also equal 3.27.

Rather, it is a reference scale against which one can judge whether the combination of pulse duration and geometry is likely to produce a strongly asymmetric sampled field while remaining reasonably smooth relative to the benchmark pulse.

Table III presents detailed geometry calculations for the subset of laser platforms from Table I with $\xi_{\text{peak}} \geq 9$ at the benchmark operating point $\chi_e = 0.002$. The incident angle θ is evaluated using the achievable field for each platform: systems with $\xi_{\text{peak}} > 10$ are normalized to $\xi = 10$, while systems with $9 \leq \xi_{\text{peak}} < 10$ are evaluated at their listed peak field. The pulse-duration parameter $\Delta\phi$ is obtained from the listed pulse length, while σ_ϕ is evaluated using an elliptical focus with $w_y = \lambda$ and $w_x = \max[\lambda, (\xi_{\text{peak}}/10)^2\lambda]$, so that only platforms with $\xi_{\text{peak}} > 10$ are stretched horizontally to reduce the peak

TABLE III. Geometric figures of merit for high- ξ laser platforms at the benchmark operating point $\chi_e = 0.002$. Calculation details are described in the preceding paragraph.

Beam case	Laser platform	θ (mrad)	$\Delta\phi$ from pulse	σ_ϕ	R_{geometry}	$\mathcal{A}_{\text{eff}}$
FCC-ee Z	Ti:Sapphire (CPA)	38.45	40.13	1.315	0.735	−0.703
FCC-ee Z	Ti:Sapphire (postcompressed)	40.53	5.08	0.127	5.122	−0.707
FCC-ee Z	OPCPA (few-cycle, NIR)	40.32	10.95	0.324	2.058	−0.707
FCC-ee Z	CO ₂ (LWIR CPA)	130.46	232.66	0.899	0.878	−0.717
EIC 18 GeV	Ti:Sapphire (CPA)	61.20	40.13	2.094	0.690	−0.542
EIC 18 GeV	Ti:Sapphire (postcompressed)	64.51	5.08	0.203	3.248	−0.707
EIC 18 GeV	OPCPA (few-cycle, NIR)	64.19	10.95	0.516	1.352	−0.709
EIC 18 GeV	CO ₂ (LWIR CPA)	207.88	232.66	1.436	0.716	−0.689
EIC 10 GeV	Ti:Sapphire (CPA)	82.12	40.13	2.811	0.684	−0.353
EIC 10 GeV	Ti:Sapphire (postcompressed)	86.56	5.08	0.272	2.451	−0.707
EIC 10 GeV	OPCPA (few-cycle, NIR)	86.13	10.95	0.693	1.067	−0.713
EIC 10 GeV	CO ₂ (LWIR CPA)	279.30	232.66	1.935	0.690	−0.584
EIC 5 GeV	Ti:Sapphire (CPA)	116.16	40.13	3.979	0.670	−0.135
EIC 5 GeV	Ti:Sapphire (postcompressed)	122.46	5.08	0.385	1.782	−0.708
EIC 5 GeV	OPCPA (few-cycle, NIR)	121.84	10.95	0.981	0.854	−0.718
EIC 5 GeV	CO ₂ (LWIR CPA)	396.30	232.66	2.763	0.683	−0.363

field to the nominal benchmark. For the ranged postcompressed Ti:sapphire entry, the values shown correspond to the shortest pulse and highest listed ξ_{peak} . The quantities R_{geometry} and $\mathcal{A}_{\text{eff}}$ are evaluated with the Mathematica definitions described in the text.

Table III makes the trends across the candidate platforms explicit. For the FCC-ee Z case, the Ti:sapphire CPA option remains unusual in that its longer intrinsic pulse can still yield $R_{\text{geometry}} < 1$ once the horizontal waist is enlarged enough to bring ξ_{peak} down to the nominal $\xi = 10$ benchmark, while retaining a large effective asymmetry. By contrast, the shortest-pulse near-infrared platforms generally produce much smaller σ_ϕ and correspondingly larger values of R_{geometry} , indicating a sampled waveform that is more strongly modified relative to the benchmark pulse. The CO₂ case is qualitatively different: despite its very long pulse, the larger wavelength forces a substantially larger collision angle, which in turn pushes σ_ϕ back toward order unity.

The same table also clarifies how the picture changes across the beam energies considered in this paper. As the electron energy is reduced from the FCC-ee Z case to the EIC cases, the required incident angle increases for every platform, and therefore, σ_ϕ generally increases as well. This drives several of the R_{geometry} values downward, especially for the Ti:sapphire CPA and CO₂ cases. Within the reduced set of higher- ξ platforms retained here, the few-cycle NIR options remain the most direct route to large $|\mathcal{A}_{\text{eff}}|$ with comparatively modest pulse durations, while the longer-pulse platforms tend to give smaller R_{geometry} once the required angle becomes large enough.

IX. CONCLUSIONS

The use of an intense laser field to induce electron polarization has been proposed previously [7,9,10]. We first used a Python code to reproduce the local constant crossed-field approximation and determine the buildup of transverse polarization over repeated laser shots for non-zero initial polarization. Our results show that lower values of the quantum efficiency parameter, χ_e , yield higher final transverse polarization but require a larger number of shots to approach the asymptotic value. We also implemented the LCFA-based spin and mean-energy kicks in Fortran within the BMAD long-term tracking code. Simulations with the FCC-ee Z lattice show that at $\chi_e = 0.002$, with a laser scanning across the length of the bunch, the polarization grows to 0.69% in 100 000 turns, corresponding to a rate of about $6.9 \times 10^{-6}\%$ per shot, with all 2214 macroparticles retained through the final turn. The associated photon spectrum should be regarded as synchrotron-radiationlike in the low- χ_e limit, with a high-energy tail controlled by $u_c = (3/2)\chi_e E$, rather than by the linear Compton edge. The expanded platform comparison in

Table III shows that, for the EIC cases and the FCC-ee Z pole case considered here, the resulting geometric figures of merit vary substantially across wavelength, pulse duration, and beam energy. In particular, long-pulse platforms can still yield comparatively small R_{geometry} after the waist is enlarged to reduce the peak field, while few-cycle near-infrared systems provide the most direct route to large effective asymmetry. Full integration of the geometry-gated effective waveform and stochastic laser recoil, using a full density-matrix framework, remains as future work and may refine these operating-point estimates.

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