Polarized electron-beam acceleration driven by vortex laser pulses

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1. Introduction

Spin is an intrinsic form of angular momentum carried by elementary particles [1]. Numerous studies in particle physics and material science have been carried out using spin-polarized electron beams [2–6]. Generally, generating polarized electrons requires conventional accelerators (Storage ring or Linac) that are typically very large in scale and budget [7, 8]. In some cases, it also needs sufficient long time to attain high polarization (couple of hours for storage rings [9]). Thanks to the rapid development of laser technology, the focal light intensities are now well beyond \(10^{20} \text{ W cm}^{-2}\) [10–12], which paves a new path to obtain high energy electron beams based on the concept of laser-driven wakefield acceleration (LWFA) [13–15]. The latter, due to the extremely high acceleration gradient, promises a more compact and cost-efficient approach for electron acceleration. However, for the realization of a laser-driven accelerator for polarized electron beams several challenges need to be addressed: (i) since a significant build-up of electron polarization from an initially unpolarized target during laser acceleration does not happen [16, 17], it requires the use of a gas target where the electron spins are already aligned before laser irradiation. (ii) Polarization losses during the injection of as many as possible electrons into a bubble structure and (iii) subsequent acceleration in the wake field must be kept under control and minimized.

With regard to polarized electron targets suitable for LWFA, we have witnessed promising developments in recent years. Electron spin polarization in strong-field ionization of atoms has been widely studied both theoretically [18–23] and experimentally [18–20]. Gas targets with spin polarization \(\sim 40\%\) have been achieved...
As stated above, preparation a pre-polarized electron target is crucial to our scheme. Known from previous literatures, the photodissociation of hydrogen halides with circularly polarized UV light yields highly spin-polarized hydrogen and halogen atoms \([16, 24, 25]\), reaching gas densities of approx. \(10^{19} \text{ cm}^{-3}\) at high polarization of the electrons in the hydrogen and halogen atoms \([26–28]\) (see detail in section 2).

The perspective of all-optical laser-driven polarized electron acceleration therefore relies on addressing the key issues (ii) and (iii), i.e. electron depolarization during LWFA. As shown below, spin depolarization mainly happens in the injection phase, which is in line with previous studies \([29]\). Recently, a method is proposed to mitigate this issue by fine tuning the focal position of the Gaussian laser beam to weaken the electron injection \([30]\). Here we propose that compared to Gaussian lasers, vortex laser beams are capable of creating a unique topology in LWFA that significantly suppresses the beam depolarization without sacrifice of the injected electron charge. Based on the proposed all-optical experiment set-up available for present techniques, we demonstrate in full three-dimensional particle-in-cell simulations the generation of high polarization electrons at very high beam charge. Such laser beams are readily accessible, for example, with the Laguerre-Gaussian (LG) mode \([31, 32]\).

The paper is organized as follows: in section 2 we give a detailed description of the method to generate fully polarized electron target. In section 3 we introduce the spin dynamics and the parameters chosen for our PIC simulations. In section 4 we present the results of simulations and analysis. Finally, in section 5, we give a brief summary.

2. Description of the all-optical set-up

As stated above, preparation a pre-polarized electron target is crucial to our scheme. Known from previous literatures, the photodissociation of hydrogen halides with circularly polarized UV light yields highly spin-polarized hydrogen and halogen atoms \([26, 27]\), at gas densities of at least \(10^{19} \text{ cm}^{-3}\) \([28]\). The polarization can approach 100% for specific photodissociation wavelengths (for HCl this is near 213 nm), and if the molecular bonds are aligned prior to photodissociation \([25]\); otherwise, if the bonds are isotropic, the polarization is reduced to 40%. However, the electronic polarization of the halogen atoms is very low; specifically, the core electrons are completely unpolarized. Therefore, to maximize the polarization of all accelerated electrons, the halogen atoms can be ionized with a resonant laser (e.g. via \(2 + 1\) resonance-enhanced multiphoton ionization, for Cl atoms at 234.62 nm \([26]\)) and removed from the target volume before irradiation by the accelerating laser pulse.

With this in mind, the all-optical scheme is sketched in figure 1. A spin-polarized hydrogen-atom target is prepared \([16]\) from the photodissociation of HCl gas emitted from a nozzle. The bolds of HCl molecule are aligned by an 1064 nm infrared (IR) light (laser) first. Then polarized electrons are obtained via photodissociation with a circularly polarized UV light (213 nm, right-handed, perpendicular to IR) so that all the electron spins are oriented along the UV light propagation direction. After that Cl atoms are ionized by another UV light (235 nm, not shown here) and removed by either external electric field or thermal expansion. This is followed by the driving vortex LG pulse with a well-controlled time synchronization. The latter, propagating coaxially with the IR light, stimulates a plasma wakefield in the gas target and accelerates pre-polarized electrons to high energies.

Figure 1. Sketch of the all-optical laser-driven polarized electron acceleration scheme with (a) the proposed experimental configuration and (b) the procedure to generate fully polarized electron target. The 1064 nm IR laser propagates along the x axis to align the bonds of the HCl molecules, then a UV light propagates along the z axis with a wavelength of 213 nm is used to photodissociate the HCl molecules. A 234.62 nm UV light is used to ionize the Cl atoms. Thermal expansion of the electrons creates large Coulomb field that expels the Cl ions. A fully polarized electron target is therefore produced for sequential acceleration by the LG laser pulse propagating along the x axis. A density ramp of length scale \(L\) is employed to ensure efficient electron injection.
The all-optical set-up requires good time synchronization between the multiple laser pulses. While aligning and photodissociation can be done simultaneously, the key issue is to control the time delay of the driving laser with respect to the previous light sources. As a matter of fact, the electron spins of hydrogen-atoms are aligned after photodissociation. For electron spin of $s = +1/2$, $n = +1/2 >$ and $|s| = +1/2, n = -1/2 >$ in the atoms ($n$ is the nuclei spin). The first hyperfine state is an eigenstate and remains unchanged but the latter oscillates between $|s| = +1/2, n = -1/2 >$ and $|s| = -1/2, n = +1/2 >$ at a period of 350 ps [16, 26, 27]. To maximize the polarization of electrons for LWFA, the time delay between the driving pulse and the photodissociation UV light should be controlled, for example, at a precision of one-tenths of the oscillation period. State-of-art technique offers synchronization precision at ps level [33, 34], therefore is sufficient for our purpose. It should be mentioned that the driving laser should also be delayed in such a way that the Cl ions are completely removed.

3. Spin dynamics and simulation parameters

Three relevant mechanisms may have influence on beam polarization, i.e. spin precession in electric and magnetic fields according to the Thomas–Bargmann–Michel–Telegdi (T-BMT) equation [35], the Sokolov–Ternov (S–T) effect (spin flip) and the Stern–Gerlach force (gradient forces). In reality, for LWFA schemes, spin precession according to T-BMT is the main influence while other effects are negligibly small [16, 17, 36]. Simple estimations are given here to explain why other effects can be omitted here. Typically the polarization time for electrons via S–T effect is about $T_{\text{pol,S–T}} = 8m_e^3c^6/5\sqrt{3}\hbar^2F^2\gamma_e^2$ [2, 17]. The S–T like effect could be significant in the strong-field QED regime and lead to certain degree of beam polarization [37, 38], relying on the collision of multi-GeV electrons with a laser pulse of focal intensity $> 5 \times 10^{22}$ W cm$^{-2}$. However, in LWFA the typical relativistic factor of electrons and field strength are $\gamma_e \sim 10^6$ and $F \sim 10^{16}$ V m$^{-1}$, respectively. One finds $T_{\text{pol,S–T}} \sim 1 \mu$s much larger than the acceleration duration (~ns scale), thus the S–T effect can be neglected in plasma wakefield. Comparing the Stern–Gerlach force to the Lorentz force $|F_{\text{SG}}|/|F_L| \sim |\nabla (S \cdot B)/\gamma_e^2c^2Bm_e|/\hbar\lambda c\gamma_e^2 \ll 1$ [2, 17] also suggests its minor effect in LWFA. Electron spin dynamics is therefore integrated into the PIC simulation code VLPL [39] according to the T-BMT equation [40]:

$$\frac{ds}{dr} = \Omega \times s$$

with

$$\Omega = \frac{e}{m}\left(\frac{B}{\gamma} - \frac{1}{\gamma + 1/c^2} \times E\right) +\alpha_e \frac{e}{m}\left(\frac{B}{\gamma} - \frac{\gamma}{\gamma + 1/c^2}(v \cdot B) - \frac{v}{c^2} \times E\right).$$

Here $e$ is the fundamental charge; $m$ is the electron mass; $v$ is the electron velocity; $\gamma = 1/(1 - v^2/c^2)^{-1/2}$ is the relativistic factor; $\alpha_e = (g - 2)/2 \approx 1.16 \times 10^{-3}$ ($g$ is the gyromagnetic factor); and the vector $s$ is the electron spin in its rest frame [2], respectively. For simplicity, the spin vector is normalized to $|s| = 1$ instead of $\hbar/2$ in our simulations. We have adopted the rotation matrix method in our moving particle module of the PIC code to minimize the numerical error in solving the T-BMT equation.

The LG laser propagates along the $x$ axis from left side of simulations window of 48 $\mu$m(x) $\times$ 48 $\mu$m ($y$) $\times$ 48 $\mu$m(z) in size and 1200 $\times$ 600 $\times$ 600 in cell numbers, and passes through the fully ionized cold plasma where the electron spins are initially aligned to $+z$ axis. For the sake of tracking the bubble wake during whole simulations, the window moves with the velocity of $v$ before the laser front reaches the right boundary. We employ the LG$_{01}$ mode laser [41–43]:

$$E(r, t) = E_0\sqrt{2}r_{\text{w}(x)} \exp\left[-\frac{r^2}{w^2(x)}\right]\sin^2\left(\frac{\pi t}{2\tau}\right)\exp\left\{-i\left[2\tan^{-1}\frac{x}{x_R} - \phi - \frac{k_0\rho^2x}{2(x^2 + x_R^2)}\right]\right\}$$

with $r^2 = y^2 + z^2$, $\phi = \arctan(z/y)$, wavelength $\lambda = 800$ nm, $k_0 = 2\pi/\lambda$, $w(x) = w_0\left[1-(x-x_0^2+x_R^2)/x_R^2\right]^{1/2}$, $x_R = \pi w_0^2/\lambda$ and the focusing position $x_R$ respectively. The laser pulse is linearly polarized along the $y$ axis. We also give results for a Gaussian laser for comparison. We use the density ramp injection method [44–46] with following profile:

$$n(\xi) = \{(\kappa - \Theta(\xi - L) - (\kappa - 1)\Theta(\xi - 2L))\sin(\pi\xi/2L) + \Theta(\xi - L)\}n_0$$

similar to the one used in [30], for better comparison. Here $\Theta(x)$ is the step function, $\xi = x - x_0, x_0 = 20$ $\mu$m, $L = 16$ $\mu$m, ratio between the peak density of the ramp and the background density $\kappa = n_p/n_0 = 4$, $n_p$ is the peak density while $n_0 = 10^{18}$ cm$^{-3}$ is the background density, respectively.
4. Results and theoretical analysis

The acceleration processes for both laser modes are illustrated in figure 2, at the same peak intensity of $8.6 \times 10^{18}$ W cm$^{-2}$ and pulse duration of 21.4 fs. The ordinary Gaussian beam drives a bubble wakefield that traps local electrons due to the decreasing phase velocity in the plasma density bump [47]. A cylindrical electron bunch is formed in the bubble center, as depicted in figures 2(c) and (d). The spins of the injected electrons are well aligned in the same direction at 75$T_e$, However, at a later stage of 150$T_e$, the spin orientations are strongly diverged at different positions. Averaging the spin projections onto the $z$ axis, one finds that the overall polarization is vanishing, i.e. the electron beam is depolarized.

To better understand the spin dynamics of electrons, we count all electrons injected into the bubble and calculate the polarization with $P = \sqrt{P_x^2 + P_y^2 + P_z^2}$ where $P_i = \sum_i s_{zi}/N$ and the averaged relativistic factor $\gamma = \sum_i \gamma_i/N$ at each interaction time. As illustrated in figure 3(a), the beam polarization barely changes when interacting with the laser field (50–65$T_e$) since the initial spin directions are parallel/anti-parallel to laser magnetic field (along the $+/-z$ axis). The polarization decreases significantly with $\gamma$ varying slowly and then followed by steady acceleration of electrons where the polarization is almost constant. The former, corresponding to the injection phase of LWFA, is when the major depolarization happens. Along the simulation, one finds that the polarization is maintained at a very high level ($\sim$88%) for the LG case, while the one for the Gaussian beam is only about 30%.

To further illustrate the polarization variation depending on the driving laser geometry, the evolution of the transverse spin component $s_z$ is shown in figure 3(b). During the interaction, the spins dilute towards the $z = -1$ end for the Gaussian pulse, while most trapped electrons accumulate at $z = 1$ for the LG laser. From the energy spectra in figure 3(c) we see that, while the cutoff and peak energies are marginally smaller for the LG case, the total number of electrons is significantly higher than that for the Gaussian beam. This generates an enormous electron-beam flux. For instance, the peak current of the LG case reaches 20 kA (polarization $\sim$88%), about 4 times larger than in the Gaussian case (5 kA, polarization $<30\%$). The boosted bunch charge or the peak flux benefits from the new geometry of the LG-laser-driven wakefield. The vortex beam produces a donut-shaped electron bunch in the vicinity of $r_0 \pm \Delta r/2$, as compared to a cylinder-like beam of radius $\Delta r$ for the Gaussian driver, corresponding to a cross-section area of $2\pi r_0\Delta r^2$ and $\pi\Delta r^2$ respectively. For a simple estimation, we use $\Delta r \sim \lambda_p/(x_0)\approx 4$ $\mu$m [30, 48–50] as the bunch radius of the trapped electrons and $r_0 = w_0/\sqrt{2} \approx 7$ $\mu$m as the center of the trapped region for the LG case [48, 49]. Accordingly, the peak-current ratio between the LG and the Gaussian case is about $2r_0/\Delta r \sim 3.5$, a factor that is well reproduced in simulations.

Figure 2. Electron density and laser field distributions (iso-surface) with $a = E_0/\max = 2$ and $x_T = 250\lambda$ at $75T_e$ and $150T_e$ for the LG laser (a) and (b) and for the Gaussian laser (c) and (d). The black arrows denote the electron spin directions. The electron density is projected onto the $x-y$ (at $z = 0$) and $y-z$ planes, while the laser electric field is projected onto the $x-z$ plane (at $y = 0$).
For the LG case separate the spin vectors $s$ within the bubble. Considering and the precession duration $s$ mainly along the $P_z$ direction. From equation $P_z$ averaged amplitudes of the averaged $Ex$ as a function of $s$. From the initial condition $\gamma = eB_0(2 + \beta_0/2)me$. To solve the equation for each particle, we separate the spin vectors $s$ into the component parallel to $\Omega$ with $s_\parallel = (s \cdot e_\phi)e_\phi$ and perpendicular to $\Omega$ with $s_\perp = (s \cdot e_\perp)e_\perp$. From the initial condition $s_0 = e_\perp$. One acquires evolution of the spin vector as $s = (e_\perp \cdot e_\phi)e_\phi + \cos(\Delta \theta_\parallel)(e_\perp \cdot e_\phi)e_\phi + \sin(\Delta \theta_\parallel)(e_\perp \cdot e_\phi)e_\perp$. Here $e_\perp$, $e_\phi$ and $e_\phi$ are normalized base vectors in cylindrical coordinates, the rotation angle $\Delta \theta_\parallel \approx (\Omega)\Delta t$ depends on the time averaged precession frequency and the precession duration $\Delta t$. Then the beam polarization in each direction follows

$$P_x = \frac{1}{N} \sum \sin(\Delta \theta_\parallel)(e_\perp \cdot e_\phi) = 0, \quad P_y = \frac{1}{N} \sum (e_\perp \cdot e_\phi)(e_\phi \cdot e_\phi) + \cos(\Delta \theta_\parallel)(e_\perp \cdot e_\phi)(e_\phi \cdot e_\phi) = 0 \quad \text{and} \quad P_z = \frac{1}{N} \sum (e_\perp \cdot e_\phi)^2 + \cos(\Delta \theta_\parallel)(e_\perp \cdot e_\phi)(e_\phi) = 0$$

The above analysis shows that the spin precession is strongly related to the azimuthal magnetic field. We averaged $B_0$ over the trapped electron region and display its values in the $y-z$ plane during injection in figure 4. For the LG case (figure 4(a)), $B_0$ is anticlockwise in the near axis region ($-5 \mu m < r < 5 \mu m$), decays to zero as the radius increases, and changes to clockwise in the outer region. Differently, the $B_0$ field for the Gaussian laser is clockwise in the whole displayed area and gradually declines (figure 4(c)). One sees that the peak magnetic field

Figure 3. Results for trapped electrons located in bubble wake. Displayed are the evolution for (a) polarization (solid lines) and electron relativistic factor $\gamma$ (averaged, dashed lines), (b) the $s_z$ distribution and (c) energy spectrum. Red and blue colors represent the LG and Gaussian laser beams.

Figure 4. Color-map of the averaged azimuthal magnetic field $B_\phi$ and current density along the $x$ direction $I_x$, during injection in the trapping region in the $y-z$ plane for the LG case (a) and (b) and for the Gaussian case (c) and (d). Cyan arrows show the directions and amplitudes of the averaged $B_\phi$. The averaged $B_\phi$ and $I_x$ as a function of the radius $r = \sqrt{y^2 + z^2}$ (e) and the longitudinal electric field $E_z$ as a function of $x-ct$ (f).

In LWFA, one usually has $B \sim B_0$, $E_x \sim -B_0$ and $\nu \sim v_x$ [47, 49], where $B_0$ is the azimuthal magnetic field within the bubble. Considering $\gamma \sim 1 \gg a$, during the injection (see figure 3(a)), the spin precession frequency from equation (3) takes the simplified form $\Omega \approx eB_0(2 + \beta_0/2)me$. To solve the equation for each particle, we separate the spin vectors $s$ into the component parallel to $\Omega$ with $s_\parallel = (s \cdot e_\phi)e_\phi$ and perpendicular to $\Omega$ with $s_\perp = (s \cdot e_\perp)e_\perp$. From the initial condition $s_0 = e_\perp$. One acquires evolution of the spin vector as $s = (e_\perp \cdot e_\phi)e_\phi + \cos(\Delta \theta_\parallel)(e_\perp \cdot e_\phi)e_\phi + \sin(\Delta \theta_\parallel)(e_\perp \cdot e_\phi)e_\perp$. Here $e_\perp$, $e_\phi$ and $e_\phi$ are normalized base vectors in cylindrical coordinates, the rotation angle $\Delta \theta_\parallel \approx (\Omega)\Delta t$ depends on the time averaged precession frequency and the precession duration $\Delta t$. Then the beam polarization in each direction follows

$$P_x = \frac{1}{N} \sum \sin(\Delta \theta_\parallel)(e_\perp \cdot e_\phi) = 0, \quad P_y = \frac{1}{N} \sum (e_\perp \cdot e_\phi)(e_\phi \cdot e_\phi) + \cos(\Delta \theta_\parallel)(e_\perp \cdot e_\phi)(e_\phi \cdot e_\phi) = 0 \quad \text{and} \quad P_z = \frac{1}{N} \sum (e_\perp \cdot e_\phi)^2 + \cos(\Delta \theta_\parallel)(e_\perp \cdot e_\phi)(e_\phi) = 0$$

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in the former is less than half of that in the latter. As a matter of fact, the azimuthal magnetic field $B_\phi$ in the cavity satisfies [45, 46]:

$$B_\phi(r) = \frac{1}{r} \int_0^r \left( \frac{j_x}{\epsilon_0 c^2} - \frac{\partial E_x}{\partial x} \right) r'dr'.$$

We assume $\partial E_x/\partial r \sim 0$ since in the blowout regime of LWFA the longitudinal electric field is slowly varying in the trapping area. Therefore, the field strength depends on two quantities: the longitudinal current density ($j_x$) and the electric field ($E_x$) along the $x$ axis. Since it is known for the blow-out regime that $|j_x/\epsilon_0 c^2|/|\partial E_x/\partial x| \sim n_0/\gamma_0 \sim 4$ [49–51], the difference is mainly due to the self-generated magnetic field of the beam current.

For a Gaussian laser, the light intensity peaks on-axis such that trapped electrons are concentrated at the center of the bubble, leading to a well-directed current in the counter-propagation direction, as shown in figure 4(d). One finds the highest field strength near the symmetry axis. On the contrary, the intensity of the LG laser beam is maximized off-axis, leaving a hollow space in the propagation center. Two consequences immediately arise: first, trapped electrons are distributed in a circular ring with their density peaking at around $r = 7 \mu$m. The new topology significantly lowers down the electron areal density and the current density for a certain amount of beam charge. Second, electrons located near the symmetry axis leak through the beam center and become the source of a counter-propagating return current in the region of $r < 5 \mu$m. These effects lead to the unique current density profile for the LG driver in figure 4(b), where the peak current density in the trapping region is only one-third of that for a Gaussian beam. The already weakened magnetic field is further effectively compensated by the anti-clockwise field resulting from the counter-propagating flux in the center. Together, they strongly reduce $B_\phi$ while maintaining the total beam charge or flux, as lined out along the radial distance in figure 4(e). We also note that the longitudinal accelerating field $E_x$, linearly dependent on the phase [50–52], is notably smaller in the LG case as seen from figure 4(f). It leads to less acceleration compared to the Gaussian case (figure 3(c)). A possible explanation is that the beam loading effect is more significant as compared to the Gaussian case due to the larger amounts of injected electrons caused by donut-shaped vortex distribution, which weakens the longitudinal electric field $E_x$ formed by the background ions.

The essence of beam depolarization is that electron spins precess at difference frequencies. Eventually the unsynchronized spins are oriented in various directions such that the averaged spin, i.e. the beam polarization, vanishes. A direct observation of the depolarization process will be tracking the specific trajectories of electrons during injection and acceleration. We chose three electrons for each case, at injection radii of 0, $\Delta r/2$, $r_0 - \Delta r/2$ for the Gaussian case and $r_0, r_0 + \Delta r/2, r_0 - \Delta r/2$ for the LG case respectively. (c) Beam polarization as a function of the peak current at 300$\tau_0$ (800 fs) for case 1: varying laser amplitude $a$ with fixed values of $\gamma_0 = 250 \lambda, \kappa = 4, n_0 = 10^{19}$ cm$^{-3}$; and case 2: varying $\kappa, n_0$ and $\gamma_0$ at fixed laser amplitude $a = 5$. The pink dotted represents polarization rate of 80%.

![Figure 5](image-url)
injection time for $\Delta t \sim \pi |r_i - r_0|/(\beta c) \sim 4\pi |r_i - r_0|/c$ LG case and $\Delta t \sim 2\pi c/\lambda_e$. Assuming the injection density is homogenous among whole injection region, the whole polarization after injection is given as:

$$P_{\text{LG}} = \left[ 1 + \sum \cos (\Delta \theta_i)/N \right]/2 \approx 0.5$$

$$P_G \approx 0.5 + \int_{r_0}^{r_0 + \Delta r/2} \cos \left( \frac{5\pi e}{B_\phi(r)} > \frac{r - r_0}{\alpha mc} \right) dr / (2\beta r \Delta r) \quad (6a)$$

$$P_G \approx 0.5 + \int_{r_0}^{r_0 + \Delta r/2} \cos \left( \frac{5\pi e}{B_\phi(r)} < \frac{r - r_0}{\alpha mc} \right) dr / (2\beta r \Delta r) \quad (6b)$$

where $P_{\text{LG}}$ and $P_G$ is the whole polarization of the beam for LG and Gaussian case respectively. We calculate a polarization of about 0.91 for the LG case and 0.31 for the Gaussian case from the $B_\phi$ profiles in figure 4(e), in good agreement with the simulation results from figure 3(a).

After electrons gaining sufficiently high energies, the spin directions remain almost unchanged (see figure 3(a)). In fact, the spin precession angle $\Delta \theta_i$ can be roughly estimated during this phase. In steady acceleration, one has $E_i \sim cB_\phi \gamma \sim \beta / c$ [45], and the precession frequency can be written as $\Omega \sim \gamma B_\phi/m \gamma (\gamma + 1)$, for $\beta \approx 1$ and $\gamma \gg 1$, suggesting the electron spin precession is slowed down due to $\gamma \approx 1$. Substituting the equation of motion for electrons $\gamma \Delta t \sim \gamma dt/dt \sim E_\phi/mc$ into the precession frequency, we finally obtain $\Delta \theta_i \sim \gamma B_\phi/E_\phi (\gamma + 1) \sim 1/\gamma \ll 1$, i.e. the spin change is negligible during this phase.

Spin depolarization imposes strong restrictions on the charge or the current of the electron beam from LWFA. One can find out the criteria for the Gaussian laser beam by taking $B_\phi \sim B_\phi (\gamma + 1)$ and $\gamma \gg 1$ for the Gaussian case in figure 3(a). In all cases the beam current is limited to $I_{\text{peak}} < 2.2$ kA for the Gaussian laser to preserve polarizations over 80% (pink dotted line in figure 5(c)), consistent with the prediction in equation (7). However, the limitation on the beam flux is released because of the vortex beam structure. The peak current for the LG case $I_{\text{peak}}$ reaches $\sim 20$ kA where the polarization is $\sim 90\%$, an order of magnitudes higher.

We find interesting oscillations of the beam polarization for the Gaussian laser in figure 5(c). This can be seen from equation (7), where the second term in the numerator periodically changes with the peak current, describing the oscillation very well. It is known that only the component perpendicular to the magnetic field in the rest frame of the electron $s_\perp$ processes. For larger peak current, the precession frequencies increase such that the transverse beam polarization, averaged over accelerated electrons, vanishes for $\Delta \theta_{\text{max}} \gg 2\pi$. In other words, $P_\perp = \sum \cos (\Delta \theta_i)/N \approx 0$, leaving the non-changing parallel to the magnetic field in the rest frame of the electron. Therefore the oscillation converges to $P = P_{\perp} = 1/2$ for growing peak current.

Our simulations are carried out for initial electron polarization along the laser magnetic field. Nevertheless, our scenario is valid for all starting spin directions. For initial polarizations $s_0 = e_\parallel$ (parallel to propagation direction) and $s_0 = e_\perp$ (parallel to electric direction) the results are illustrated in figure 6(a). The polarizations of both cases are 81.2% for $s_0 = e_\parallel$ and 88.6% for $s_0 = e_\perp$, validating the novel effect of the vortex laser geometry. For the longitudinal pre-polarized case ($s_0 = e_\parallel$), the polarization oscillation still happens as shown in figure 6(b). Unlike in the transverse polarization case, one has $P = P_\perp = \sum \cos (\Delta \theta_i)/N$ for $s_0 = e_\parallel$, converging to 0 polarization at sufficiently large currents. The restriction on peak current is also stronger, e.g. to preserve polarization $>80\%$, $I_{\text{peak}} < 1.5$ kA.

5. Conclusions

In conclusion, we proposed a promising scheme to generate polarized electron beams via LWFA driven by a vortex LG Laser. According to our 3D-PIC simulations involving particle spin dynamics, electron beam is of polarization over 80% is achieved with high beam charge and peak flux. Compared to the Gaussian laser driven acceleration, the restriction on the electron beam current density is released, thanks to the novel topology of the
vortex LG laser. By releasing the limitation on the beam loading, the averaged electron energy is naturally lower for the LG case. This could be compensated, for example, by enlarging the laser pulse energy as compared to the Gaussian case (at no loss of beam polarization). The scheme relies on an all optical set-up that is accessible at state-of-the-art facilities.

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