Generation of arbitrarily polarized GeV lepton beams via nonlinear Breit-Wheeler process

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A B S T R A C T
Generation of arbitrarily spin-polarized electron and positron beams has been investigated in the single-shot interaction of high-energy polarized γ-photons with an ultraintense asymmetric laser pulse via nonlinear Breit-Wheeler pair production. We develop a fully spin-resolved semi-classical Monte Carlo method to describe the pair creation and polarization. In the considered general setup, there are two sources of the polarization of created pairs: the spin angular momentum transfer from the polarized parent γ-photons, as well as the asymmetry and polarization of the driving laser field. This allows to develop a highly sensitive tool to control the polarization of created electrons and positrons. Thus, dense GeV lepton beams with average polarization degree up to about 80%, adjustable continuously between the transverse and longitudinal components, can be obtained by our all-optical method with currently achievable laser facilities, which could find an application as injectors of the polarized e+e- collider to search for new physics beyond the Standard Model.

1. Introduction

Ultrarelativistic spin-polarized electron and positron beams have fundamental applications in particle and high-energy physics [1,2]. Experiments with polarized leptons are envisaged in future e+e- colliders, such as International Linear Collider (ILC) [3], Compact Linear Collider (CLIC) [4] and Circular Electron Positron Collider (CEPC) [5], where the longitudinal polarization of leptons provides high sensitivity, in particular, suppressing background from W+W- boson and single Z boson production via WW fusion [3], enhancing different triple gauge couplings in WW pair production [3,6], and improving top vector coupling in top quark production [7]. Meanwhile, the transverse polarization can cause asymmetric azimuthal distribution of final-state particles [8], and facilitate the search of new physics beyond the Standard Model [9,10] via, e.g., measuring relative phases among helicity amplitudes in W+W- pair production [11], probing the mixture of scalar-electron states [12] and searching for graviton in extra dimensions [13]. Commonly, longitudinal and transverse polarizations are studied separately since these effects are independent of each other [3,8]. However, it has also been recognized that the arbitrarily spin-polarized (ASP) lepton beams (having both longitudinal and transverse polarization components) have special importance, in particular, introducing a spin-frame of reference (required for the full reconstruction of the density matrix of a spin-1/2 particle) and modify effective beyond the Standard Model vertices [14]. Thus, they play an unique role in future beyond the Standard Model experiments in e+e- colliders, e.g., rendering special spin structure functions observable in vector and axial-vector type beyond the Standard Model interactions [10], diagnosing spin and chirality structures of new particles in antler-topology processes [15], and producing polarized top quark pairs as a probe of new physics [16].

In conventional methods, the transversely polarized lepton beams can be directly obtained in a storage ring via Sokolov-Ternov effect [17], which demands a long polarization time since the corresponding static magnetic fields are relatively weak (~Tesla). The longitudinally polarized leptons can be created in a Bethe-Heitler pair production process [18] via high-energy circularly polarized (CP) γ-photons interacting with high-Z target [19]. In the latter the low luminosity of γ-photons beam is compensated by a requirement of its high repetition rate to yield a dense positron beam for applications [20]. Generally, the transverse and longitudinal polarizations can be transformed to

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Fig. 1. Scenario of generation of ASP lepton beams via nonlinear BW process. A LP asymmetric laser pulse (propagating along \(-z\) direction and polarizing along \(\hat{z}\) axis) head-on collides with polarized \(\gamma\)-photons \((\gamma_p)\) to create ASP electron and positron beams. (a)-(c) indicate LP, CP and EP \(\gamma\)-photons, respectively, and (d)-(f) show average polarizations of created positrons \(\vec{S}_+\) corresponding to (a)-(c), respectively. The red-solid arrow and ellipse indicate the direction and amplitude of \(\vec{S}_+\) respectively. The red-dashed arrows in (d) and (e) indicate particular cases of neglecting the polarization of \(\gamma\)-photons and employing symmetric laser field, respectively.

Each other by a spin-rotator which, however, has drawbacks demanding quasi-monoenergetic beams and is accompanied by the beam intensity reduction \([21]\).

Modern ultrashort ultraintense laser pulses \([22,23]\) enable alternative efficient methods to generate dense polarized lepton beams in femtoseconds via nonlinear quantum electrodynamics (QED) processes \([24-26]\), such as nonlinear Compton scattering and Breit-Wheeler (BW) \(e^+e^-\) pair production \([27-30]\). As reported, the leptons can be efficiently transversely polarized in a standing laser wave \([31]\), an elliptically polarized (EP) \([32]\) or a bichromatic laser field \([33]\), but not in a symmetric monochromatic laser field \([34]\). Longitudinally polarized positrons can be produced by circularly polarized (CP) \(\gamma\)-photons through the helicity transfer in the nonlinear BW process \([35]\). Two-step schemes are possible: firstly low-energy linearly polarized \(\gamma\)-photons are generated from pulsed photocathodes \([36]\), polarized atoms \([37]\) or molecular photodissociation \([38]\), and then accelerated to ultrarelativistic energies, for instance, via conventional accelerators, or via the laser- \([39,40]\) or beam-driven \([41]\) plasma wakefield. Although there are ways to generate transversely or longitudinally polarized laser beams, the engineering of ASP beams of desired polarization still remains a challenge.

Fig. 2. The polarization mechanism of ASP positron beams. (a) and (b): Spin of a sample positron \(\vec{S}_+\) created by LP (a) and EP (b) \(\gamma\)-photons, respectively. \(\vec{P}\) indicates the average polarization vector (i.e. the instantaneous SQA) with two transverse components \(D_{\text{mag}}\) and \(D_{\text{LP}}\) and one longitudinal component \(D_{\text{CP}}\) (see the analytical formula in Eq. 8). For LP \(\gamma\)-photons in (a), \(D_{\text{LP}}=0\) and the positron polarization is engineered in the transverse plane, while in the case of EP \(\gamma\)-photon in (b), 3-dimensional control of \(\vec{P}\) is possible. Compared with the case of a LP photon, the polarization \(\vec{P}\) of an EP photon is less intuitive and thus not shown in (b). In our interaction scheme (Fig. 1), \(\vec{P}_1 = \vec{E}_{-} + \xi_{\gamma}\) and \(\vec{P}_2 = -\vec{E}_{+} + \xi_{\gamma}\), \(\xi_{\gamma}\) indicates the polarization angle to \(\vec{P}_1\); \(\theta_1\) and \(\theta_2\) are the angles of \(\vec{P}\) to \(D_{\text{LP}}\) and \(D_{\text{mag}}\), respectively. (c) and (d): Averaged-over-energy polarization of the positrons \(\vec{S}_+\) created by LP (c) and EP (d) \(\gamma\)-photons, respectively. The red arrow and circle (ellipsoid) indicate the direction and amplitude of \(\vec{S}_+\) (see the analytical formula in Eq. 12), respectively.

In this article, the generation of ASP GeV lepton beams has been investigated in the interaction of polarized \(\gamma\)-photons with a counterpropagating ultraintense linearly polarized (LP) asymmetric tailored (e.g., bichromatic or frequency-chirped) laser pulse (see the interaction scenario in Fig. 1). With a fully spin-resolved semiclassical Monte Carlo algorithm, we describe photon-polarization-dependent pair production in the nonlinear BW process. We find that the polarization of created pairs is controlled by the polarization of parent \(\gamma\)-photons, and by the polarization and asymmetry of the strong laser field, via the spin angular momentum transfer and the asymmetric spin-dependent pair production, respectively (Fig. 2 and Eq. 12). We discuss how the ASP positron beam can be built up by controlling the polarization of \(\gamma\)-photons and the asymmetry parameters of the laser field. Our simulations show that dense GeV lepton beams with continuously adjustable polarization direction and degree up to about 80% can be obtained with currently achievable laser facilities to the benefit of many unique applications (Fig. 4).

2. Methods

To have a rich pair yield in the nonlinear BW process, one assumes that the nonlinear QED parameter is large: \(\chi_{\gamma} \equiv |e| \sqrt{-(\text{FWH}^2)/m^3} \geq 1\) \([24]\). For the radiative polarization of created pairs the photon emissions are important, which are enhanced at a large value of the parameter \(\chi_{\gamma} \equiv |e| \sqrt{-(\text{FWH}^2)/m^3} \geq 1\) \([24]\). Here, \(e\) and \(m\) are the electron charge and mass, respectively, \(k_\gamma\) and \(p\) the 4-momenta of \(\gamma\)-photon and positron (electron), respectively, and \(F_{\text{bw}}\) the field tensor. Relativistic units with \(c = \hbar = 1\) are used throughout. The photon polarization can be characterized by the unit vector \(\vec{P} = \cos(\theta_\gamma)\hat{\vec{E}}_+ + \sin(\theta_\gamma)\vec{E}_-\vec{e}_\phi\), and the corresponding Stokes parameters are \((\xi_1, \xi_2, \xi_3) = (\sin(2\theta_\gamma)\cos(\phi_\gamma), \sin(2\theta_\gamma)\sin(\phi_\gamma), \cos(2\theta_\gamma))\) \([42]\). Here, \(\vec{P}_1\) and \(\vec{P}_2\) are two orthogonal basis vectors, \(\theta_\gamma\) the polarization angle, \(\phi_\gamma\) the absolute phase, \(\xi_1\) and \(\xi_2\) describe the linear polarization, and \(\xi_3\) the circular polarization.

Let us first summarize our methods of numerical simulation and analytical estimation. Note that in the nonlinear BW process the polarization of the electrons is similar with that of the positrons, thus for simplicity we only discuss the case of the positrons below.

2.1. Numerical simulation method

In our Monte Carlo algorithm we employ the fully spin-resolved pair production probability \(W_{\text{pair}}\), calculated via the QED operator method of Baier-Katkov \([43]\) in the local constant field approximation \([24,44]\):

\[
d^2W_{\text{pair}} \, \frac{d\vec{S}_+}{d\vec{S}_+} = \frac{1}{2} (G_0 + \xi_1 G_1 + \xi_2 G_2 + \xi_3 G_3)
\]

Equation (1) is valid when the invariant field parameter \(a_0 \equiv |e| E_0 m_{\text{e}} \gg 1\), and the \(\gamma\)-photon energy is limited by \(E_{\gamma}\), where \(E_0\) and \(a_0\) are the laser field amplitude and frequency, respectively. The variables introduced in Eq. 1 read:

\[
G_0 = \frac{W_0}{2} \left( \text{IntK}_\gamma(\rho) + \frac{\epsilon_1^2 + \epsilon_2^2}{\epsilon_3} \text{K}_\gamma(\rho) + \left[ \text{IntK}_\gamma(\rho) - 2\text{K}_\gamma(\rho) \right] (S_+ \cdot S_-) \right)
\]

540
\begin{equation}
+K_2(\rho) \left[ \frac{\epsilon_f}{\epsilon} \left( S_{\alpha} \cdot \mathbf{b}_{\gamma} \right) + \frac{\epsilon_f}{\epsilon} \left( S_{\alpha} \cdot \mathbf{b}_{\alpha} \right) \right] + \frac{\epsilon_f + \epsilon_f^2}{\epsilon} \text{Int} K_2(\rho) \\
- \left[ \frac{\epsilon_f \epsilon_f}{\epsilon} \right] K_2(\rho) \left( S_{\alpha} \cdot \mathbf{v}_{\alpha} \right) \left( S_{\gamma} \cdot \mathbf{v}_{\gamma} \right) \right] \right]
\end{equation}

(2)

\begin{equation}
G_1 = \frac{W_0}{2} \left[ K_1(\rho) \left[ \frac{\epsilon_f}{\epsilon} \left( S_{\alpha} \cdot \mathbf{a}_{\alpha} \right) + \frac{\epsilon_f}{\epsilon} \left( S_{\alpha} \cdot \mathbf{a}_{\gamma} \right) \right] + \frac{\epsilon_f - \epsilon_f}{\epsilon} K_1(\rho) \left( S_{\alpha} \times S_{\alpha} \right) \cdot \mathbf{v}_{\alpha} \right]
\end{equation}

(3)

\begin{equation}
G_2 = \frac{W_0}{2} \left[ K_1(\rho) \left[ \frac{\epsilon_f}{\epsilon} \left( S_{\alpha} \times S_{\alpha} \right) \cdot \mathbf{a}_{\alpha} \right] + \frac{\epsilon_f - \epsilon_f}{\epsilon} K_1(\rho) \left( S_{\alpha} \times S_{\alpha} \right) \cdot \mathbf{v}_{\alpha} \right]
\end{equation}

(4)

\begin{equation}
G_3 = \frac{W_0}{2} \left[ K_1(\rho) + \frac{\epsilon_f + \epsilon_f}{2 \epsilon} K_1(\rho) \left( S_{\alpha} \cdot \mathbf{b}_{\alpha} \right) \left( S_{\alpha} \cdot \mathbf{v}_{\alpha} \right) \right]
\end{equation}

(5)

where \( W_0 = \frac{2}{\epsilon} \sum_{\alpha=0}^{\infty} \int dK_1(z) \), \( \rho = 2 \epsilon_f \left| \left( \epsilon_f / \epsilon \right) + \epsilon_f \right| \) is the fine structure constant, \( \epsilon_f \) and \( \epsilon_f \) are the energies of the parent photon, created electron and positron, \( \mathbf{b}_{\alpha} = \mathbf{v}_{\alpha} \times \mathbf{v}_{\alpha} \approx -\mathbf{k}_{\alpha} \times \mathbf{E} = -\mathbf{B} \), \( \mathbf{E} \) is the electric field, \( \mathbf{E} \), \( \mathbf{v}_{\alpha} \), \( \mathbf{v}_{\gamma} \), \( \mathbf{a}_{\alpha} \), \( \mathbf{a}_{\gamma} \), \( \mathbf{b}_{\alpha} \) and \( \mathbf{b}_{\alpha} \) are the unit vectors along the electric and magnetic fields, the positron velocity and acceleration, respectively, and \( S_{\alpha} \) denotes the positron (electron) spin vector. Note that the Stokes parameters must be transformed from the photon initial frame (\( \mathbf{P}_i \), \( \mathbf{P}_e \), \( \mathbf{n} \)) to the pair production frame (\( \mathbf{P}_1 \), \( \mathbf{P}_2 \), \( \mathbf{n} \)) [45], where \( \mathbf{P}_1 = [\mathbf{E} - \mathbf{n} \cdot \mathbf{E} - \mathbf{n} \cdot \mathbf{n}] \) and \( \mathbf{P}_2 = \mathbf{n} \times \mathbf{P}_1 \), with the photon propagation vector \( \mathbf{n} \).

To study the positron polarization \( S_{\alpha} \), we first sum over the electron polarization \( S_{\gamma} \) in \( W_{pair}^{+} \), and the probability relying on \( S_{\alpha} \) is simplified as:

\begin{equation}
\frac{d^2 W_{pair}}{d\epsilon d\epsilon} = \frac{W_0}{2} \left( C + S_{\alpha} \cdot D \right)
\end{equation}

(6)

where \( C = \text{Int} K_1(\rho) + \frac{\epsilon_f + \epsilon_f^2}{\epsilon} K_2(\rho) - \epsilon_f \), \( D = D_{mag} + D_{LP} + D_{CP} \),

\begin{align*}
D_{mag} &= D_{mag} + D_{LP} + D_{CP} + \text{Int} K_2(\rho) \\
D_{LP} &= D_{LP} + D_{LP} + D_{CP} \\
D_{CP} &= D_{CP} + D_{CP}
\end{align*}

According to Eq. 8, the average polarization of the positron with an energy \( \epsilon_f \) created at a certain moment is:

\begin{equation}
\langle \mathbf{S}_{\alpha} \rangle = \frac{\mathbf{P}_1 \left[ D_{mag} + D_{LP} + D_{CP} \right]}{\langle \mathbf{D} \rangle}
\end{equation}

(9)

Although \( \mathbf{D} \) and \( C \) in Eq. 10 change significantly with \( \epsilon_f \), \( \mathbf{P}_1 = \mathbf{D} / \langle \mathbf{D} \rangle \) is almost constant when \( \epsilon_f \) is small (Fig. 3), resulting in that \( \langle \mathbf{S}_{\alpha} \rangle \) does not change much with the increase of the laser field strength. However, the average polarization of the positrons with energy \( \epsilon_f \) in positive half laser cycles \( \langle \mathbf{S}_{\alpha} \rangle \) is different from that in negative half laser cycles \( \langle \mathbf{S}_{\alpha} \rangle \). \( \langle \mathbf{D}_{mag} \rangle \) and \( \langle \mathbf{D}_{LP} \rangle \) are anti-parallel in the positive and negative half laser cycles; while \( \langle \mathbf{D}_{CP} \rangle \) roughly doesn’t change. Thus, the average polarization of the positrons with energy \( \epsilon_f \) created in the laser field can be estimated as:

\begin{equation}
\langle \mathbf{S}_{\alpha} \rangle = \frac{N_{\alpha}(\mathbf{n}) + N_{\alpha}(\mathbf{n})}{N_{\alpha}(\mathbf{n}) + N_{\alpha}(\mathbf{n})}
\end{equation}

(10)
plays a more important role on the determination of the final positron polarization (Fig. 6a).

For LP γ−photons, \( \mathbf{D}_{\text{mag}} = |\mathbf{D}_{\text{mag}}^{(LP)}| \) with \( \theta_1 = \theta_2 \) results in that \( \mathbf{S}_\gamma \) is parallel to \(-\mathbf{P}\) (Fig. 2c). In the case of EP γ−photons, as \( \theta_2 \) increases, \( \mathbf{S}_\gamma \) will rotate anti-clockwise by an azimuth angle φ, which can be calculated by Eq. 12 (Fig. 2d). Thus, our intuitive expectations expressed in Eq. 12 are the following. The transverse polarization governs by the laser field asymmetry \( A_{\text{field}} = \mathbf{S}_\gamma / |\mathbf{S}_\gamma| \), with \( P_\gamma = \mathbf{D}_{\text{mag}} + \mathbf{D}_{\text{pol}}^{(CP)}/|\mathbf{C}| \), while the average longitudinal polarization \( \mathbf{S}_\gamma \) is solely determined by \( \mathbf{D}_{\text{pol}}^{(CP)}/|\mathbf{C}| \propto \xi_2 \). In a symmetric laser field \( A_{\text{field}} = 0 \), the transverse polarization \( \mathbf{S}_\gamma \) is vanishing, and only the longitudinal polarization \( \mathbf{S}_\gamma \) can be obtained by employing CP γ−photons (\( \xi_2 \neq 0 \)) [35] (see the red-dashed arrows in Fig. 1e).

3. Results

Our Monte Carlo simulations provide the following results for generation of ASP positron beams (Fig. 4). The employed laser and γ−photon parameters are as follows. A tightly focused LP bi-chromatic Gaussian laser pulse [33] propagates along \(-z\) direction and is polarized along the \( \hat{z} \) axis (Fig. 1). The phase difference between two-color fields is \( \Delta \phi = \pi/2 \) to obtain the maximal field asymmetry. The invariant field parameters of two laser fields are \( \alpha_1 = 60 \) and \( \alpha_2 = 15 \), and the peak intensity of the total beam \( I_0 = 1.1 \times 10^{12} \) W/cm². Other parameters are wavelength \( \lambda_2 = 2 \lambda_1 = 1 \mu \text{m} \), pulse durations \( r_1 = r_2 = 15 \) ps, with periods \( T_1 = 2 T_2 \), and focal radii \( w_1 = w_2 = 5 \mu \text{m} \). This kind of laser pulse is currently feasible in petawatt laser facilities [22,23]. The concrete form of the asymmetric laser pulse is of minor importance, e.g., a frequency-chirped laser pulse [46] can work similarly. The γ−photon beam is of a cylindrical form and propagates along \( z \) direction, with an initial energy \( \epsilon_1 = 1.8 \text{GeV} \), energy spread \( \Delta \epsilon_1/\epsilon_1 = 0.06 \), angular divergence \( \Delta \theta_1 = 0.3 \text{ mrad} \), beam radius \( w_1 = 1 \mu \text{m} \), beam length \( L_z = 5 \mu \text{m} \), photon number \( N_\gamma = 10^8 \) and density \( n_\gamma \approx 6.37 \times 10^{15} \text{cm}^{-3} \). It has a transversely Gaussian and longitudinally uniform distribution. Such a γ−photon beam may be generated via synchrotron radiation [47], bremsstrahlung [19], linear [48] or nonlinear Compton scattering [49]. At these parameters, \( \mathbf{T}_\gamma \approx 0.96 \) and the pair production is significant.

According to Eq. 12, the transverse and longitudinal polarizations of created positron beam (\( \mathbf{S}_\gamma \) and \( \mathbf{S}_\gamma^L \)) can be controlled by adjusting the polarization of parent γ−photons (\( \theta_1 \) and \( \theta_2 \), as shown in Fig. 4a and b), which indicate the spin angular momentum transfer from parent γ−photons to created pairs. The transverse polarization \( \mathbf{S}_\gamma \propto \sqrt{\xi_1} \) and the pair production is significant.

Furthermore, the angle-resolved transverse polarization is asymmetric, because it relies on the asymmetry \( A_{\text{field}} \) of the laser field. In contrast, the longitudinal polarization is transferred from the circular polarization component of the γ−photon \( \mathbf{S}_\gamma^L \propto \xi_2 \), and therefore, its symmetric angular distribution is not disturbed by the laser field asym-
Fig. 4. Results of analytical estimations and numerical simulations for ASP positron beams. (a) and (b): $\bar{S}_x$ and $\bar{S}_y$ of the positrons with respect to $\theta_x$ and $\theta_y$, respectively, analytically estimated by Eq. 12 with $A_{\text{rad},\psi}$ = 0.8378. The black points correspond to ($\theta_y = 50^\circ$, $\theta_x = 70^\circ$). (c)-(f) are our numerical simulation results with ($\theta_y = 50^\circ$, $\theta_x = 70^\circ$) after the interaction, including the quantum radiative depolarization effect of the positrons propagating through the laser field. (c)-(e): $\bar{S}_x$, $\bar{S}_y$ and the angle-resolved positron density $\log_{10}(d^2N/d\theta_xd\theta_y)$ (rad$^{-2}$) with respect to the deflection angles $\theta_x = \arctan(p_{x,\gamma}/p_{\gamma})$ and $\theta_y = \arctan(p_{y,\gamma}/p_{\gamma})$, respectively. The black curves indicate $\bar{S}_x$, $\bar{S}_y$ and $\log_{10}(dN/d\theta_xd\theta_y)$ (rad$^{-2}$) summing over $\theta_x$ vs $\theta_y$, respectively. (f) $\bar{S}_x$ (red), $\bar{S}_y$ (blue) and the energy-resolved positron density $\log_{10}(dN_e/d\epsilon_e)$ (GeV$^{-1}$) (green) vs $\epsilon_e$. Other parameters are given in the text.

Fig. 5. Simulation results for the case excluding the quantum radiative depolarization effect. (a) and (b): $\bar{S}_x$ and $\bar{S}_y$ with respect to $\theta_x$ and $\theta_y$, respectively. Other parameters are the same as those in Fig. 4.

metry. In the given ultrastark laser field, the yield rate of the positrons $N_p/N_e \approx 0.1$ ($N_e$ can be calculated via Fig. 4e, and $N_e$ is known) is much higher than that of the common method employing Bethe-Heitler pair production ($\sim 10^{-3} - 10^{-4}$) [19,47,48], since $N_e \sim a_{\text{tox}}$ is rather large in the ultraintense laser field [24,43]. The flux of the positron beam is approximately $3 \times 10^{15}/\text{s}$, and the pulse duration is determined by that of the $\gamma$-photon beam. The emittance of the positron beam is about $\epsilon \approx \epsilon_{\text{div}} \times 10^{-4}$ mm-mrad, which fulfills the experimental requirements of the beam injectors [50], with radius $w_y = l_y = 1 \mu$m and angular divergence $\theta_{\text{div}} \approx 30$ mrad (see FWTH in Fig. 4e). Because of the stochastic effects of the pair production and further radiation, the energy of the positron beam spreads, with a peak of the energy distribution at $\epsilon_e \approx 0.3$ GeV (Fig. 4f).

The transverse and longitudinal polarizations show opposite behaviour with respect to the positron energy. With the increase of $\epsilon_e$, $\bar{S}_x$ ($\bar{S}_y$) monotonically decreases (increases) from above 90% (0%) to about 50% (above 85%), because the pair polarization is mainly determined by the energy of the laser (parent $\gamma$-photons) at low (high) $\epsilon_e$.

4. Discussion

For the experimental feasibility, we have investigated the impact of the laser and $\gamma$-photon parameters on the quality of the positron beam polarization, as shown in Fig. 7. The transverse and longitudinal polarizations $\bar{S}_x$ and $\bar{S}_y$ are inversely proportional to $\chi_x \propto \dot{\alpha}e_\epsilon \propto \dot{\alpha}e_\gamma$ (Fig. 6a, b), while the pair production probability is proportional to $\chi_y$. Thus, as $\alpha_1$ and $\epsilon_\gamma$ increase, $\bar{S}_x$ and $\bar{S}_y$ decrease, while $N_p$ increases (Fig. 7a, b). When the laser pulse duration $\tau_\epsilon$ increases, more pairs can be generated, and $\bar{S}_x$ and $\bar{S}_y$ decrease due to the radiative depolarization effect being enhanced (Fig. 7c). Furthermore, the yield and polarization of the positrons is quite stable with respect to the variations of the colliding angle $\theta_\gamma$, the energy spread $\Delta\epsilon_e/\epsilon_e$, and angular divergence $\Delta\theta_\gamma$ of the $\gamma$-photon beam in reasonable limits (Fig. 7d-f). Thus, the quality of
5. Conclusion

A method is put forward for the creation and 3-dimensional polarization engineering of high-energy positron beams. We show that employing polarized γ-rays and an asymmetric laser field, the positron beams with desired polarizations can be produced. The direction and magnitude of the average polarization vectors are sensitively controlled by the laser field asymmetry and polarization as well as the γ-photon polarization. With this method dense arbitrarily spin-polarized GeV lepton beams with polarization degree up to about 80% can be obtained with currently achievable laser facilities [22,23], which have unique applications for high-energy and particle physics.

Declaration of Competing Interest

The authors declare that they have no conflicts of interest in this work.

Acknowledgments

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