

All-optical ultrafast spin rotation for relativistic charged particle beams

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An all-optical method of ultrafast spin rotation is put forward to precisely manipulate the polarization of relativistic charged particle beams of leptons or ions. In particular, laser-driven dense ultrashort beams are manipulated via single-shot interaction with a co-propagating moderate temporally asymmetric (frequency-chirped or subcycle THz) laser pulse. Using semi-classical numerical simulations, we find that in a temporally asymmetrical laser field, the spin rotation of a particle can be determined from the flexibly controllable phase retardation between its spin precession and momentum oscillation. An initial polarization of a proton beam can be rotated to any desired orientation (e.g., from the common transverse to the more useful longitudinal polarization) with extraordinary precision (better than 1%) in tens of femtoseconds using a feasible frequency-chirped laser pulse. Moreover, the beam qualities, in terms of energy and angular divergence, can be significantly improved in the rotation process. This method has potential applications in various areas involving ultrafast spin manipulation, like laser-plasma, laser-nuclear and high-energy particle physics.

Relativistic beams of spin-polarized charged particles, such as leptons and ions, play a crucial role in exploring the fundamental structures [1–3] and interactions [4–7], with the advantage of an increased number of measurable observables and unprecedented test precision [8, 9]. They can provide direct access to study new physics beyond the standard model [10], characterize the quantum numbers and chiral couplings of new particles [8, 11], investigate spin-dependent nuclear interactions [12–14], and increase the cross-section and control the angular distribution of the reaction products in nuclear fusion reactions [15, 16]. In the famous proton-spin puzzle [17, 18], energetic polarized proton beams are applied to precisely measure the internal spin-flavor structure of the quarks and the gluon spin distribution of the target proton in proton-proton collisions [19–23]. In these experiments, the longitudinal spin polarization is generally sensitive to the helicity distributions of the quarks and gluons [19, 23], while the transverse one is sensitive to the transversity distributions of quarks and parton orbital angular momenta [3, 20]. Production of longitudinally spin-polarized (LSP) and transversely spin-polarized (TSP) particle beams with a high degree of polarization ($\gtrsim 70\%$ [3, 8]) is, therefore, quite desirable.

In general, production of a relativistic polarized proton beam is accomplished in two steps. First, a low-energy beam is produced from an optically-pumped polarized H^- ion source [24] or a polarized atomic beam source [25], and then injected into a high-energy accelerator, where it is commonly rotated to the TSP state in order to suppress the depolarization effect. A relativistic TSP electron beam, by comparison, can be obtained from a low-energy one produced by photoemission from a GaAs-

based cathode [26] and then accelerated by conventional means, or directly employing the Sokolov-Ternov effect in a storage ring [27]. Relativistic LSP proton and electron beams are obtained via spin-polarization manipulation in conventional spin rotators [28, 29] and Siberian Snakes [30, 31] (the latter are employed to avoid the polarization loss due to imperfections and intrinsic depolarizing spin resonances [30, 32]), both of which consisting of a sequence of vertical and horizontal arc dipoles and superconducting solenoids (or specially arranged magnets). The conventional spin rotators manipulate the polarization adiabatically, and the manipulation precision is sensitive to the beam energy spread and the perturbing fields around the particle trajectories. They generally take tens of minutes or even hours to switch between TSP and LSP states, and typically run few times per day in order to cancel the remaining potentially undesirable systematic effects [3, 29].

Rapid advances in ultraintense ultrashort laser technique, with peak intensities reaching the scale of 10^{23} W/cm², pulse durations of tens of femtoseconds and energy fluctuations of about 1% [33–35], offer new avenues for polarized beam generation. Relatively low-cost and efficient laser-driven plasma accelerators with a gradient exceeding 0.1 TeV/m are capable of providing dense tens-of-MeV proton [36–38] and multi-GeV electron beams [39]. In all-optical setups, leptons can be transversely polarized in a standing-wave [40–42], elliptically polarized [43, 44], or bichromatic laser pulses [45–48] due to the quantum radiative spin-flip effect, since the transverse laser field dominates the spin polarization process; while the more useful LSP lepton beams can only be produced indirectly via the helicity transfer from cir-

cularly polarized γ photons in linear [49] or nonlinear Breit-Wheeler pair production processes [50, 51]. The latter is generally pre-produced via Compton scattering [52] or bremsstrahlung [53]. High-energy polarized electron [54–56] and proton beams [57–59] can also be produced through laser-driven wakefield acceleration of pre-polarized low-energy ones, generated via a photodissociated hydrogen halide gas target [60–62]. Unfortunately, ultrafast spin manipulation (in particular, spin rotation from TSP to LSP state) of these beams is still a great challenge.

In this Letter, an all-optical ultrafast spin rotation method is put forward to precisely manipulate the polarization of relativistic (laser-driven) ultrashort particle beams through interaction with a co-propagating moderate temporally asymmetrical (e.g., frequency-chirped or sub-cycle THz) laser pulse (see Fig. 1). Focus will be on a TSP ion or lepton beam to be rotated to the LSP state (or any desired polarization orientation) by a frequency-chirped laser pulse. This will be achieved with extraordinary precision (better than 1% in tens of femtoseconds) while significantly improving beam qualities in terms of the energy and angular divergence (see Fig. 2). We find that in a strong laser field the rotation of a particle's spin is determined from the phase retardation between its spin precession and momentum oscillation, which can be quite pronounced and flexibly controllable for temporally asymmetrical laser fields (see Fig. 3). This method is demonstrated to be robust with respect to the parameters of the laser and particle beams (see Fig. 4 and [63]), is realizable with currently achievable laser facilities, and thus has significant applications in broad areas involving ultrafast spin manipulation.

In our simulations the particle's spin and momentum dynamics will be treated semiclassically. For an electron the quantum radiation effects are characterized by the invariant parameter $\chi \equiv |e|\hbar\sqrt{(F_{\mu\nu}p^\nu)^2}/m_e^3c^4 \approx a_0\gamma(\hbar\omega_0/m_e c^2)(1 - |\beta|\cos\theta)$ [64, 65], where \hbar is the Planck constant and c the speed of light; p^ν , β , γ , e and m_e are the 4-momentum, velocity normalized by c , Lorentz factor, charge and mass of the electron, respectively. Also, $F_{\mu\nu}$, ω_0 , $a_0 \equiv |e|E_0/(m_e\omega_0c)$, and E_0 are the field tensor, frequency, invariant intensity parameter, and amplitude of the laser pulse, respectively. Note that χ for a much heavier lepton or ion is much smaller than that of an electron in the same field. With $\chi \ll 1$, the quantum radiation reaction (RR) effects on the particle's momentum and spin (Sokolov-Ternov effect) are negligible [54–59, 66–70]. Furthermore, the Stern-Gerlach force is much smaller than the Lorentz force in our simulations (see Fig. 2) and can be dropped [71, 72]. Thus the particle motion can be adequately described by Newton-Lorentz equations, and its spin precession is governed by Thomas-Bargmann-Michel-Telegdi (T-BMT) equation, $d\mathbf{S}/d(\omega_0 t) = \mathbf{\Omega}_s \times \mathbf{S}$, with $\mathbf{\Omega}_s = -\frac{g}{m}[(a + \frac{1}{\gamma})\mathbf{B} - \frac{a\gamma}{\gamma+1}(\boldsymbol{\beta} \cdot \mathbf{B})\boldsymbol{\beta} - (a + \frac{1}{\gamma+1})(\boldsymbol{\beta} \times \mathbf{E})]$

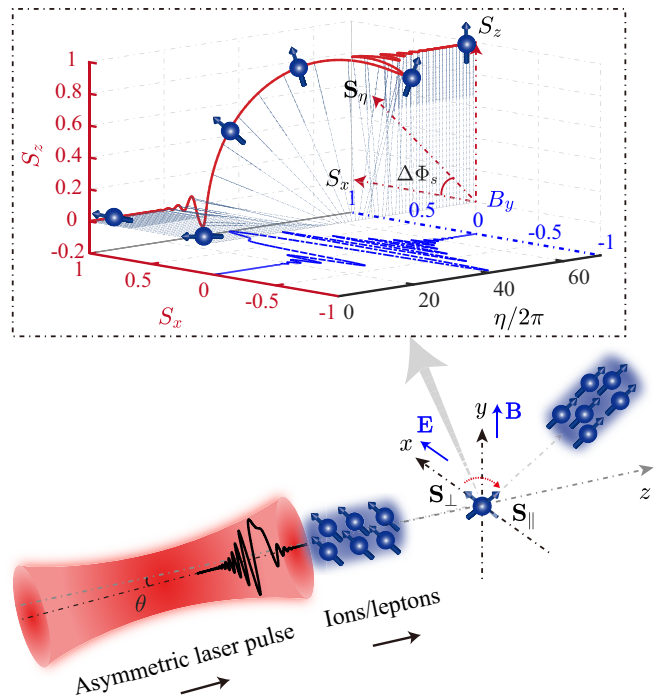


FIG. 1. Interaction scenario. The TSP particle beam (“ \mathbf{S}_\perp ”, along $+x$ direction and perpendicular to its initial momentum direction) is rotated to the LSP state (“ \mathbf{S}_\parallel ”, parallel to its final momentum direction). \mathbf{E} and \mathbf{B} indicate the electric and magnetic components of the laser field, respectively. (Inset) Particle spin evolution. θ is the laser incidence angle with respect to the particle beam, $\Delta\Phi_s$ the spin rotation angle and η the laser phase.

[73–75]. Note that we use dimensionless units throughout. q , m and $\mathbf{\Omega}_s$ are the particle charge, mass and precession frequency normalized by e , m_e and ω_0 , respectively, \mathbf{E} and \mathbf{B} are normalized by $|e|/(m_e\omega_0c)$, \mathbf{S} is the particle's spin vector in its rest frame, and $a = (g - 2)/2$ the anomalous magnetic moment of the particle with the gyromagnetic factor g . For the proton and electron, $a \approx 1.793$ and 1.16×10^{-3} , respectively. Quantum RR effects as photon recoils and stochastic spin flips must be included for $\chi \gtrsim 1$ [40–48, 70], while for $10^{-3} \lesssim \chi \ll 1$, RR effects can be treated classically [68, 70, 76].

Sample results of spin polarization rotation of proton and electron beams in moderately frequency-chirped laser pulses are illustrated in Fig. 2. An initially TSP particle beam co-propagating with the laser pulse has spin $(\tilde{S}_x, \tilde{S}_y, \tilde{S}_z) = (1, 0, 0)$. The particles are initially distributed in a cylinder of radius $R = 1.5\lambda_0$ and length $L = 0.5\lambda_0$, and have a number density $n_0 \approx 2.26 \times 10^{16} \text{ cm}^{-3}$, and transversely Gaussian and longitudinally uniform profiles. The initial kinetic energy of the proton (electron) beam is $\varepsilon_p = 1 \text{ MeV}$ ($\varepsilon_e = 0.1 \text{ MeV}$) with 1% energy spread and angular divergence (FWHM) of about 11.8° . Such pre-polarized beams can be obtained via photodissociation of the aligned hydrogen halide molecules with

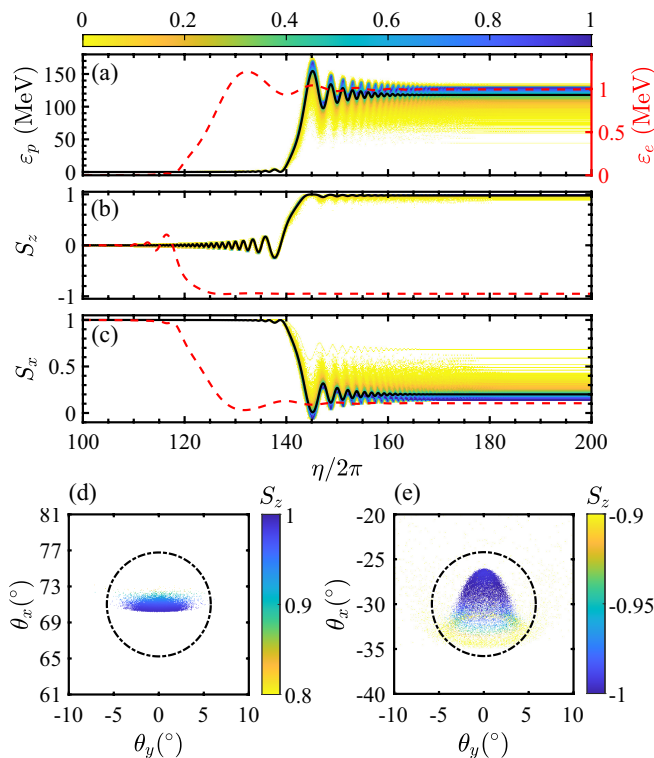


FIG. 2. (a)-(c): Variations of $dN/d\epsilon_p$, dN/dS_z and dN/dS_x of the proton beam (color) with respect to η , respectively, and the black curves show the corresponding average values. N is the particle number. The red-dashed curves show the average values for the case of employing an electron beam. (d) and (e): Angle-resolved distributions of S_z vs the transverse deflection angles $\theta_x = \arctan(p_x/p_z)$ and $\theta_y = \arctan(p_y/p_z)$ for the proton and electron beams, respectively. The black-dashed circles denote the corresponding beam initial profile sizes (FWHM). Simulation parameters are given in the text.

circularly polarized ultraviolet light [60–62]. We employ a focused, linearly-polarized, and linearly-chirped Gaussian laser pulse, polarized along $+x$ and propagating along $+z$ ($\theta = 5^\circ$ in Fig. 1), with an instantaneous frequency $\omega(\eta) = \omega_0(1 + b\eta)$ [77–79], where b is the dimensionless chirp parameter, $\eta = \omega_0(t - z/c)$, and ω_0 the initial frequency at $z = 0$. For the proton (electron) beam, the laser peak intensity is $I_0 \approx 3.46 \times 10^{21}$ W/cm² (5.54×10^{16} W/cm²) with $a_0 \approx 8.5 \times 10^{-10} \lambda_0 [\mu\text{m}] \sqrt{I_0 [\text{W}/\text{cm}^2]} \approx 50$ (0.2), wavelength $\lambda_0 = 1 \mu\text{m}$, pulse duration $\tau_0 = 30T_0$ ($10T_0$), period T_0 , focal radius $w_0 = 5 \mu\text{m}$ and $b = -0.00539$ (-0.0167). Such a chirped laser pulse could in principle be produced in experiments by using dispersion filters [80, 81] (in particular, the linear chirp is adjusted by the group delay dispersion [82]) or by the reflection from a relativistic ionization front (interface gas plasma) [83]; see also [84–86]. For these values, the quantum parameters are $\chi_p \sim 10^{-8}$ ($\chi_e \sim 10^{-7}$) and, hence, the RR effects are negligible in both cases.

Consider what happens as the front of the laser pulse

catches up and interacts with the protons [in the range of $\eta/2\pi \lesssim 137$ in Figs. 2(a)-(c)]. Since the front part of the pulse is almost symmetrical [see Fig. 3(a)], the proton spins oscillate initially only slightly, and the initial TSP state is well preserved. This is followed by interaction with the highly asymmetrical part of the laser pulse, resulting in precise manipulation and rotation of the proton spin from the TSP state (\mathbf{S}_\perp , denoted by S_x in the interaction geometry of Fig. 1, in which $S_y \approx 0$) to the LSP state (\mathbf{S}_\parallel , denoted by S_z) during the range of $137 \lesssim \eta/2\pi \lesssim 147$. Due to the negative chirp effect, the proton spin polarization can be tuned in a nearly single elongated cycle [see Fig. 3(a)]. Finally, the proton interacts with the approximately symmetrical rear part of the laser pulse in the range of $\eta/2\pi \gtrsim 147$, and its spin oscillations become weaker and weaker and eventually die down. The final average longitudinal polarization is $\bar{S}_z \approx 98\%$, symmetrically distributed with respect to θ_y [see Fig. 2(d)], and the proton beam is evidently compressed by the laser field (transverse ponderomotive force), particularly in θ_x . Moreover, the proton beam is significantly accelerated with an average energy gain of about 120 MeV (that further remarkably decreases the angular divergence $\sim \Delta p_{x,y}/p_z$), due to the phase synchronization between the laser and proton beams [77, 87], with an energy spread $< 10\%$ [63]. The energy gain may reach the GeV scale with the appropriate presently available laser parameters [63, 78]. These results may be optimized further by moderately increasing the laser focal radius, allowing the protons to experience a quasi-uniform laser field. Such a highly-LSP low-divergence proton beam with a hundred-MeV energy can be employed as the polarized injection source for high-energy accelerators [71], or to search for new physics by investigating CP or T -symmetry [88] (e.g., null high-precision tests of T -violation by polarized proton-deuteron scattering at energy of 135 MeV [89]).

For heavier charged particles like a deuteron or a triton, the spin rotations and energy gains are moderately reduced compared with those of a proton [63], because the spin manipulation efficiency mainly depends on the charge-to-mass ratio and g -factor of the particle [90]. While, for much lighter leptons such as electrons and positrons (with smaller g -factors, too), it's vice versa, and consequently a much weaker laser pulse can satisfy the requirements. For instance, interacting with a tens-of-GW laser of $I_0 \approx 5.54 \times 10^{16}$ W/cm², the electron beam can achieve a final polarization of $\bar{S}_z \approx -0.95$, with a bell-shaped angular distribution [see Fig. 2(e)].

The spin manipulation mechanism is analyzed in Fig. 3. In the interaction geometry of Fig. 1, $\mathbf{E} = \hat{i}E_x$ and $\mathbf{B} = \hat{j}B_y$. For protons with $\gamma \approx 1.128$ (see Fig. 2), the spin precession frequency $|\boldsymbol{\Omega}_s| \approx -\frac{q}{\gamma m}(\gamma a + 1)|\mathbf{B}| \propto B_y$, yielding $\Omega_{s,x} \approx \Omega_{s,z} \approx 0$ and $\Omega_{s,y} \approx 0.064$, according to the T-BMT equation above. Using a plane-wave laser pulse [63], the

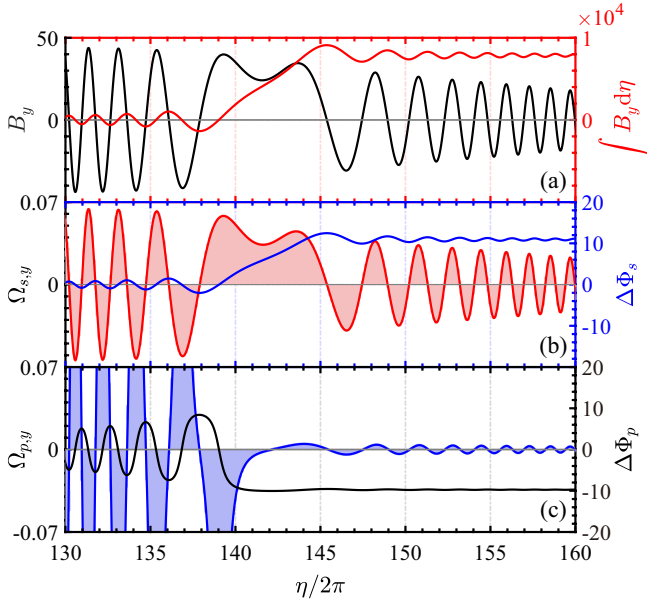


FIG. 3. (a) Evolutions of the normalized magnetic field B_y sensed by the proton (black) and its time-integrated profile (red). (b) Spin precession frequency $\Omega_{s,y}$ (red) and $\Delta\Phi_s$ (blue) vs η . (c) Momentum angular frequency $\Omega_{p,y}$ (blue) and $\Delta\Phi_p$ (black) vs η . Other simulation parameters are the same as those in Fig. 2.

spin rotation angle can also be analytically estimated from $\Delta\Phi_s = \int_0^\eta |\mathbf{\Omega}_s| d\eta' \approx a_0 C \sqrt{\frac{\pi}{2b}} [\cos(\frac{1}{4b}) \text{F}_c(\frac{1+2b\eta}{\sqrt{2\pi b}}) + \sin(\frac{1}{4b}) \text{F}_s(\frac{1+2b\eta}{\sqrt{2\pi b}})] + C_0$, where $C = (a + 1/\gamma)q/m$, C_0 is a constant of integration, and $\text{F}_c(\cdot) = \text{FresnelC}(\cdot)$ and $\text{F}_s(\cdot) = \text{FresnelS}(\cdot)$ are Fresnel cosine and sine integrals. Thus, the proton spin can be expressed as $S_z(\eta) \approx \sin[\theta_0 + \Delta\Phi_s(\eta)]$ [63] with the initial polarization angle θ_0 (for a TSP proton beam $\theta_0 = 0$). $\Delta\Phi_s$ depends sensitively on the time-integrated $\mathbf{\Omega}_s$ (subsequently on \mathbf{B}) and further on a_0 , τ_0 , b and g -factor. With the proton momentum direction defined as $\hat{\mathbf{n}} \equiv \boldsymbol{\beta}/\beta$, its angular frequency is $\mathbf{\Omega}_p = (\hat{\mathbf{n}} \times \mathbf{E}/\beta - \mathbf{B})q/\gamma m = [\gamma^2 \boldsymbol{\beta} \times \mathbf{E}/(\gamma^2 - 1) - \mathbf{B}]q/\gamma m$ [63], which yields $\Omega_{p,x} \simeq \Omega_{p,z} \approx 0$ and $\Omega_{p,y} \approx 0.5$. Thus, during interaction with the temporally symmetrical laser field, $|\mathbf{\Omega}_p| \gg |\mathbf{\Omega}_s|$ with opposite precession directions and the phase of $\Omega_{p,y}$ is ahead of that of $\Omega_{s,y}$ by π [see Figs. 3(b) and (c) in $\eta/2\pi \lesssim 137$ and $\eta/2\pi \gtrsim 147$]. This clearly demonstrates that the proton spin direction with respect to its momentum direction, i.e., the spin polarization, is unchanged during interaction with the symmetrical part of the laser pulse, where the spin precession does not result in a net change in $\Delta\Phi_s$ and $\Delta\Phi_p$ simultaneously. Tuning the laser field by an optimal linear negative chirp, its temporal symmetry can be destroyed. This results in an intense quasi-static positive part covering the range $137 \lesssim \eta/2\pi \lesssim 147$. This part exhibits low local instantaneous frequency and small phase variations [see Fig. 3(a)]. The time-integrated nor-

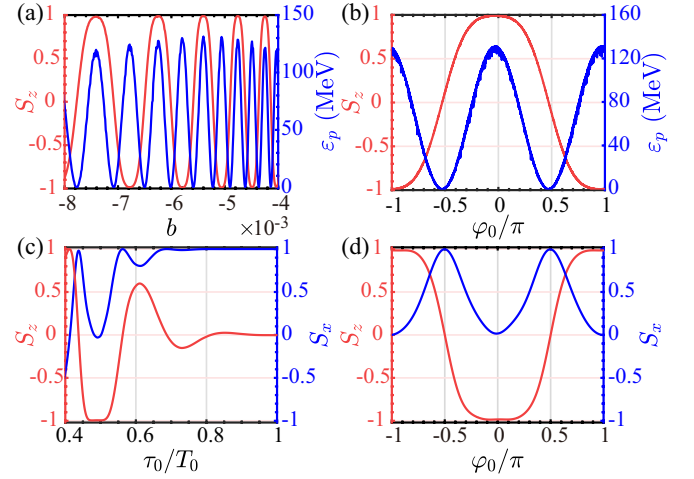


FIG. 4. (a) and (b): Final S_z and ε_p of the proton beam vs b and the carrier-envelope phase (CEP) φ_0 of the chirped laser pulse, respectively. (c) and (d): Final S_z and S_x of the electron beam vs τ_0 and φ_0 of the THz pulse. In (c) and (d), $a_0 = 2.5$ and $\lambda_0 = 300 \mu\text{m}$. Other parameters are the same as those in Fig. 2.

malized magnetic field is quite large ($\sim 8 \times 10^3$). During interaction with the quasi-static part of the pulse, the positive and negative half-cycles affect the proton momentum and spin dynamics substantially asymmetrically. Meanwhile, the phase retardation between $\Delta\Phi_s$ and $\Delta\Phi_p$ decreases rapidly and then varies slightly over a relatively long time. Thus a comparable match of integrated spin precession frequency $|\Delta\Phi_s| \simeq 6.83\pi/2$ and momentum angular frequency $|\Delta\Phi_p| \simeq 6.19\pi/2$ of opposite signs can be acquired with the employed parameters, giving rise to a net phase retardation of about $\pi/2$ [see Figs. 3(b) and 3(c)]. Subsequently, sufficient spin polarization rotation (from TSP to LSP state) can be achieved. In addition, our numerical simulations demonstrate that spin manipulation of the protons from TSP to LSP state, or vice versa, should satisfy the condition $|\Delta\Phi_s - \Delta\Phi_p| = (2j - 1)\pi/2$ [or $\gamma \simeq 1 + (2j - 1) \times 0.13$, for given parameters], where j is a positive integer.

For experimental feasibility, impact of other laser and particle beam parameters on the efficiency of the spin rotation is investigated in Fig. 4 and [63]. Viewed as functions of the chirp parameter b the final S_z and kinetic energy gain ε_p of the proton beam exhibit local fluctuations of the order of 10^{-3} , with the optimal rotation occurring, for an optimal ε_p , for $b \approx -0.00539$ [see Fig. 4(a)] and $\varphi_0 = 0, \pm\pi$ [see Fig. 4(b)]. The adjustable precision, $\delta S_z/S_z \sim 0.27$, is estimated using feasible b and a_0 parameter values which may result in fluctuations not exceeding 10^{-5} [80, 81] and 0.01 [33–35], respectively. Note that in experiments the linear chirp parameter b can achieve the required precision via mainly adjusting the group delay dispersion, by employing the devices, such as acousto-optic programmable dispersive filters (Daz-

zler) [80, 81], and, the fluctuations in Fig. 4(a) can be further optimized (slowed down) by simultaneously controlling φ_0 and moderately enlarging τ_0 , which would remarkably reduce the required precision of b . Also, \bar{S}_z increases exponentially with a_0 and decreases with τ_0 , while \bar{S}_x follows an opposite trend. The final average energy, $\bar{\varepsilon}_p$, increases almost linearly with a_0 and τ_0 , while the energy spread remains less than 10% [63].

Besides, the spin rotation process of a realistic sub-cycle THz pulse [91–93] interacting with a pre-polarized electron beam is investigated here, too, with the fields $B_y \simeq E_x = a_0 \cos(\varphi_0 + \eta) \exp[-(\eta - 4\tau_0/T_0)^2/(\tau_0/T_0)^2]$. In Figs. 4(c) and (d), the optimal LSP state is obtained for $\tau_0 = T_0/2$, which corresponds to $\varphi_0 = 0, \pm\pi$. The results stay stable within the adjustable range of $\tau_0 \pm 5\%$ ($\varphi_0 \pm 9.3\%$) for the parameters employed. Finally, our method is widely effective for temporally asymmetrical laser pulses, also including the case of laser pulses accompanied by additional intense low-frequency magnetic fields [94, 95], however, it is not applicable for scenarios that are spatially asymmetrical (e.g., bichromatic) [63].

In conclusion, an all-optical ultrafast spin rotation method has been put forward to precisely manipulate the polarization of relativistic charged particle beams using a moderate temporally asymmetrical laser pulse. This method is feasible with currently available laser facilities and would provide a new technique for carrying out polarized beam experiments and high-precision spin-dependent measurements in laser-plasma physics, laser-nuclear physics, and high-energy particle physics.

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