Radiation from a polarized vacuum in a laser-particle collision

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The probability of photon emission of a charged particle traversing a strong field becomes modified if vacuum polarization is considered. This feature is important for fundamental quantum electrodynamics processes present in extreme astrophysical environments and can be studied in a collision of a charged particle with a strong laser field. We show that for today's available 700 GeV (6.5 TeV) protons and the field provided by the next generation of lasers, the emission spectra peak is enhanced due to the vacuum polarization effect by 30% (suppressed by 65%) in comparison to the traditionally considered Compton process. This striking phenomenon offers an alternative path to the laboratory-based manifestation of vacuum polarization.

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I. INTRODUCTION

According to quantum theory, under the action of a strong field, the vacuum behaves as a medium with a modified refractive index for a probe electromagnetic (EM) radiation [1–3]. As the next generation of lasers will provide very strong EM fields, they will enable the study of the fundamental physics of photon-photon interactions [4–8]. Several setups for measuring the signatures of vacuum polarization on probe photons using the intense lasers have already been suggested (see Refs. [9,10] and references therein). The basic idea is to detect probe photons whose kinematics or polarization properties differ from the original ones [11]. However, the signature of vacuum polarization on probe photons still has not been observed [12]. The possibility of experimentally studying the quantum vacuum is also attractive for modeling the processes that are important for astrophysics as the intense lasers will allow us to reach the strong-field regime of quantum electrodynamics (QED), similar to that thought to exist in extreme environments where the vacuum polarization effects are expected [13]. For example, the emission of magnetars is thought to be polarized due to the vacuum polarization effect [14,15] (see, also, the remarks in [16]). Here we demonstrate the following feature: the radiation from a charged particle traversing a strong-field region provided, e.g., by magnetars or by the ultra-intense laser can be drastically enhanced or reduced once vacuum polarization is taken into account.

When a strong EM field is applied to a vacuum, the virtual electron-positron pairs become polarized. For a photon traversing a strong-field region, the situation can be conveniently described by assigning an index of refraction to the vacuum that allows for various nonlinear phenomena such as birefringence and vacuum Cherenkov radiation [17]. Thus, as a consequence of vacuum polarization, the probability of photon emission by a charged particle interacting with a strong EM field is governed by the synergy between synchrotronlike and Cherenkov mechanisms [18–20]. It is important to note that due to the synergism, it is an indecomposable physical process when the resulting radiation exhibits features which cannot be recovered by superposing the constituents [18,19]. As the most intense EM fields are nowadays produced only by lasers [21], there is growing interest in observing the synergic Compton-Cherenkov (SCC) radiation [22] of charged particles in a strong EM field. The counterplay between the Cherenkov radiation and nonlinear Compton scattering has been studied in [22–24] and the possibility of Cherenkov radiation in the cosmic microwave background radiation has been discussed in [25]. In comparison to these papers, our study considers the whole energy spectrum of emitted photons, not only the limit of low-energy photons, i.e., Cherenkov ones. Here we show that the SCC process exhibits so-far unexplored features in photon-emission distribution. This striking phenomenon offers an alternative path to the laboratory-based manifestation of vacuum polarization.

In this paper, we explore the features in photon spectra emerging due to the emission of a charged particle in a polarized vacuum and demonstrate differences from the traditionally considered tree-level, i.e., lowest-order, process of nonlinear Compton scattering in which vacuum polarization effects are not considered [26]. We show that sizable differences in photon-emission spectra are expected for protons of today's available energy and the EM field provided by the next generation of lasers. Namely, the amount of emitted energy can be either significantly enhanced or reduced as a consequence of the synergic nature of photon emission. These features of photon emission could serve as a signature of vacuum polarization. While it is well known that in nonlinear Compton scattering the recoil effect reduces the emitted energy [27], we show that recoil also controls the amount of energy (not) emitted due to the effects of vacuum polarization implemented in the SCC process. In particular, a recoil experienced by the emitting particle, which has been considered to be negligible in the case of quantum Cherenkov radiation [28], plays an important role in the above-mentioned more general processes as it suppresses the effect of vacuum polarization on photon emission.

II. POWER SPECTRUM FOR PHOTON EMISSION IN A POLARIZED VACUUM

The interaction of the charged particle and photon with a strong EM field is characterized by two Lorentz invariant parameters, $\chi = e\hbar\sqrt{-(F^{\mu\nu}p_\nu)^2}/m^3c^4$ and $\chi_\gamma = \hbar\sqrt{-(F^{\mu\nu}k_\nu)^2}/m_e c E_S$, respectively, where e is the elementary charge, $\hbar$ is the reduced Planck constant, m is the particle mass, $F_{\mu\nu} = \partial_\mu A_\nu - \partial_\nu A_\mu$ is the EM field tensor, A_ν is the four-potential, p_ν is the four-momentum of the emitting particle, k_ν is the photon four-wave vector, m_e denotes the electron mass, and c is the speed of light in a vacuum [26]. The Schwinger-Sauter field $E_S = m_e^2c^3/e\hbar$ is approximately 1.33×10^{16} V/cm [29]. The strength of the EM wave is characterized by the Lorentz invariant parameter $a_0 = eE_0/m\omega_0 c$, where E_0 is the amplitude of the electric field, $\omega_0 = 2\pi c/\lambda$ is the angular frequency, and λ is the wavelength.

Since we are interested in photon emission in the quantum vacuum characterized by an index of refraction $n = 1 + \Delta n$, where $|\Delta n| \ll 1$ is the change of refraction index due to vacuum polarization, we need to derive the formula for photon emission that includes the effect of the index of refraction. Therefore, we use the semiclassical approach, where the motion of the energetic particle is treated classically while photon emission is treated quantum mechanically, as it presents a powerful method for numerical calculation of photon emission [30]. For $a_0 \gg 1$ (i.e., constant crossed field approximation) and under the weak-field approximation, i.e., $\chi^2 \gg \mathcal{F}, \mathcal{G}$ and $\mathcal{F}, \mathcal{G} \ll 1$ where $\mathcal{F} = |\mathbf{E}^2 - \mathbf{B}^2|e^2\hbar^2/m^4c^6$ and $\mathcal{G} = |\mathbf{E}\cdot\mathbf{B}|e^2\hbar^2/m^4c^6$ are the normalized Poincaré invariants of EM field, the following approach is applicable to any particle-field configuration characterized by the same value of a parameter χ [2,26].

The energy $\mathcal{W}$ radiated by a charged particle per unit frequency interval $d\omega$ and per unit solid angle $d\Omega$ is [31]

$$\frac{d^2\mathcal{W}}{d\omega\, d\Omega} = \frac{e^2}{4\pi^2 c}\left(\frac{\mathcal{E}'^2 + \mathcal{E}^2}{2\mathcal{E}^2}|\boldsymbol{I}|^2 + \frac{\hbar^2\omega^2 m_e^2 c^4}{2\mathcal{E}^4}|J|^2\right), \quad (1)$$

where

$$\boldsymbol{I} = \int_{-\infty}^{\infty} \frac{\boldsymbol{\nu}\times[(\boldsymbol{\nu}-\boldsymbol{\beta})\times\dot{\boldsymbol{\beta}}]}{(1-\boldsymbol{\nu}\cdot\boldsymbol{\beta})^2}\exp[i\omega' t(1-\boldsymbol{\nu}\cdot\boldsymbol{\beta})]\,dt \quad (2)$$

and

$$J = \int_{-\infty}^{\infty} \frac{\boldsymbol{\nu}\cdot\dot{\boldsymbol{\beta}}}{(1-\boldsymbol{\nu}\cdot\boldsymbol{\beta})^2}\exp[i\omega' t(1-\boldsymbol{\nu}\cdot\boldsymbol{\beta})]\,dt. \quad (3)$$

In Eqs. (1)–(3), $\mathcal{E}$ is the particle energy, $\hbar\omega$ is the photon energy, $\mathcal{E}' = \mathcal{E} - \hbar\omega$, $\omega' = \omega\mathcal{E}/\mathcal{E}'$, and $\boldsymbol{\nu}$ is the direction of emission [31]. Vector $\boldsymbol{\beta} = \boldsymbol{v}/c$, $\boldsymbol{v}$ is the particle velocity and $\dot{\boldsymbol{\beta}}$ represents the differentiation of $\boldsymbol{\beta}$ with respect to time t.

The derivation procedure of Eq. (1) is outlined in Ref. [30]. It considers the following relation:

$$\hat{\mathcal{E}}(\hat{p} - \hbar\hat{k}) = \sqrt{\mathcal{E}'^2 \underbrace{+2\hbar\omega\mathcal{E}(1-\boldsymbol{\nu}\cdot\boldsymbol{\beta})}_{A}\underbrace{-\hbar^2(\omega^2 - \boldsymbol{k}^2c^2)}_{B}}, \quad (4)$$

where $\hat{\mathcal{E}}$, $\hat{p}$, $\hat{k}$ are the operators of particle energy, particle momentum, and photon momentum, respectively (see Eq. (2.20) in [30]). Vector $\boldsymbol{k}$ is the wave vector of the emitted photon, in a medium $|\boldsymbol{k}| = (1+\Delta n)\omega/c$ [32]. In our case, the role of a medium is played by polarized vacuum [33]. However, the semiclassical approach can only be considered as an approximative solution in the regime where $|\Delta n| \ll 1$. For an exact solution, a full one-loop calculation of the Compton scattering process in a strong EM field needs to be performed as it would automatically include the effect of vacuum polarization. We note that the properties of photons characterized by $k_\nu^2 \neq 0$ have been studied in Refs. [34–36].

At first, we consider the case of $\Delta n = 0$, which corresponds to the situation in which vacuum polarization effects are neglected. Therefore, the term B vanishes in Eq. (4). As shown in Ref. [30], the factor $1 - \boldsymbol{\nu}\cdot\boldsymbol{\beta}$ appearing in part A of Eq. (4) directly passes to the argument of the exponentials in Eqs. (2) and (3). When an ultrarelativistic particle is assumed ($\gamma \gg 1$, where γ is the particle Lorentz factor), then $1 - \boldsymbol{\nu}\cdot\boldsymbol{\beta} \simeq 1/(2\gamma^2) + \theta^2/2 + O(1/\gamma^4)$ is considered in Eq. (1), where $\theta \approx 1/\gamma$ is an angle between $\boldsymbol{\nu}$ and $\boldsymbol{\beta}$ [30]. In such a case, the result of the integration of Eq. (1) over Ω can be recast in the form that is used for numerical calculation of the photon-emission power spectrum in nonlinear Compton scattering,

$$\frac{d\mathcal{P}}{d\omega} = \frac{\alpha\omega}{\sqrt{3}\pi\gamma^2}\left[\frac{\zeta^2 - 2\zeta + 1}{1-\zeta}K_{2/3}(\eta) - \int_\eta^\infty K_{1/3}(y)dy\right], \quad (5)$$

where $\alpha = e^2/\hbar c$ is the fine-structure constant, $\zeta = \hbar\omega/\mathcal{E}$, $\eta = 2\zeta/[3\chi(1-\zeta)]$, and $K_l(m)$ are the modified Bessel functions of the second kind [27,30]. Thus, $\omega^2 - \boldsymbol{k}^2c^2 = 0$ in Eq. (4) refers to a radiation process in an external field [30].

On the other hand, the term $\omega^2 - \boldsymbol{k}^2c^2 \neq 0$ is essential when radiation in a medium is considered [32]. In the case of a polarized vacuum, we assume that $|\boldsymbol{k}| = (1+\Delta n)\omega/c$ holds within the interaction region and thus we need to consider $\Delta n \neq 0$ in Eq. (4) [32]. Neglecting the small terms $\propto \Delta n/\gamma^4$ and $\propto \Delta n^2$, we obtain

$$1 - \boldsymbol{\nu}\cdot\boldsymbol{\beta} \approx \frac{1}{2\gamma^2} + \frac{\theta^2}{2} - \Delta n \quad (6)$$

and

$$B \approx 2\hbar^2\omega^2\Delta n. \quad (7)$$

Since the term A in Eq. (4) is approximately equal to

$$A \approx 2\hbar\omega\mathcal{E}\left(\frac{1}{2\gamma^2} + \frac{\theta^2}{2} - \Delta n\right), \quad (8)$$

we obtain

$$A + B = 2\hbar\omega\mathcal{E}\left[\frac{1}{2\gamma^2} + \frac{\theta^2}{2} - \Delta n\left(1 - \frac{\hbar\omega}{\mathcal{E}}\right)\right]. \quad (9)$$

The last term in the square brackets of Eq. (9) is a correction accounting for the effect of the index of refraction. This effect weakens as photon energy $\hbar\omega$ increases. Including this correction in Eq. (1) and following the procedure outlined in Ref. [31], we obtain the differential power spectrum for photon emission in a quantum vacuum characterized by $n = 1 + \Delta n$,

$$\frac{d^2\mathcal{P}}{d\omega\, d\theta} = e^2\omega^2\rho\left\{\frac{\mathcal{E}'^2+\mathcal{E}^2}{2\mathcal{E}'^2}F\left[\xi^2\mathrm{Ai}'^2(\xi)+\frac{\theta^2}{F}\mathrm{Ai}^2(\xi)\right]+\frac{\omega^2m^2}{2\mathcal{E}'^2\mathcal{E}^2}F\,\mathrm{Ai}^2(\xi)\right\}, \tag{10}$$

where

$$F = \frac{1}{\gamma^2}+\theta^2-2\Delta n\left(1-\frac{\hbar\omega}{\mathcal{E}}\right), \tag{11}$$

$$\xi = \left(\frac{\omega'\rho}{2}\right)^{2/3}F, \tag{12}$$

$\mathrm{Ai}(\xi)$ is the Airy function, and ρ is the radius of curvature of the particle trajectory.

III. SYNERGIC CHERENKOV-COMPTON RADIATION IN A VACUUM POLARIZED BY A STRONG ELECTROMAGNETIC FIELD

Let us now consider a photon characterized by χ_γ counterpropagating with respect to the constant crossed EM field ($\boldsymbol{E}\perp\boldsymbol{B}$, $|\boldsymbol{E}| = |\boldsymbol{B}|$). Due to the possibility of creating an electron-positron pair by a photon in a field, the field can be considered as a homogeneous anisotropic medium with dispersion and absorption and thus can be characterized by assigning the complex index of refraction $\tilde{n}(\chi_\gamma) = 1 + \Delta\tilde{n}(\chi_\gamma)$. In the limits $\chi_\gamma \ll 1$ and $\chi_\gamma \gg 1$, the complex $\Delta\tilde{n}(\chi_\gamma)$ can be expressed as

$$\Delta\tilde{n}(\chi_\gamma) = \frac{\alpha m_e^2c^4}{2(\hbar\omega)^2}\times\begin{cases}\left[\frac{11\mp3}{90\pi}\chi_\gamma^2 - i\sqrt{\frac{3}{2}}\frac{3\mp1}{16}\chi_\gamma\exp(-8/3\chi_\gamma)\right], & \chi_\gamma \ll 1\\ -\frac{5\mp1}{28\pi^2}\sqrt{3}\Gamma^4\left(\frac{2}{3}\right)(1-i\sqrt{3})(3\chi_\gamma)^{\frac{2}{3}}, & \chi_\gamma \gg 1,\end{cases} \tag{13}$$

where the minus (plus) sign corresponds to the parallel (perpendicular) polarization of the photon with respect to the electric field, and $\Gamma(x)$ is the Gamma function [33,37]. For a given χ_γ, the corresponding value of $\Delta\tilde{n}(\chi_\gamma)$ can be obtained from Eq. (2.7) in Ref. [33]. The real part of the index of refraction determines the angle and intensity of SCC radiation, while its imaginary part describes photon absorption by the EM field as it is connected with a probability of electron-positron pair production by the emitted photon that becomes important for large values of χ_γ [26,33]. For $\chi_\gamma \gg 1$, the ratio of imaginary and real parts of $\tilde{n}(\chi_\gamma)$ in the studied parameter range is of the same order as if $\mathrm{Re}[\Delta\tilde{n}(\chi_\gamma)] = 0$ was considered. As shown in Refs. [2,30,32], for $\mathrm{Re}[\Delta\tilde{n}(\chi_\gamma)] = 0$, the imaginary part of the refraction index is neglected in the calculation of the photon-emission process even in the case of large χ_γ. Similarly, in the following, we consider the effect of the real part of the index of refraction $\Delta n(\chi_\gamma) = \mathrm{Re}[\Delta\tilde{n}(\chi_\gamma)]$ in the calculation of photon-emission distribution as we are interested in the SCC photon spectrum and its difference from the Compton one. For further considerations related to electron-positron pair production in a polarized vacuum, we note that backreactions from pair creation, damping of an incident photon due to the decay into the electron-positron pair, and the radiative energy loss of a charged particle should be self-consistently taken into account [38,39]. Thus, the exact treatment of the interaction dynamics requires the development of a corresponding numerical model that is beyond the scope of this paper.

Since $\Delta n(\chi_\gamma) \neq 0$, the phase velocity of light is either higher $[\Delta n(\chi_\gamma) > 0]$ or lower $[\Delta n(\chi_\gamma) < 0]$ than the free-space value. While Δn is positive for $\chi_\gamma \ll 1$ and has a maximum at $\chi_\gamma \approx 0.75$, for $\chi_\gamma \gtrsim 15$ the value of Δn becomes negative, and thus the Cherenkov radiation vanishes [33]. Due to the difference of a nonlinear vacuum refraction index from unity, the probability of photon emission by a relativistic charged particle propagating in a strong EM field is modified. In general, the emission of a photon in a medium is governed by the SCC process that depends on both positive and negative values of $\Delta n(\chi_\gamma)$. Near $n\beta \approx 1$, where $\beta = |\boldsymbol{\beta}|$, there is a transition region that can be divided into two branches: the Compton branch for $n\beta < 1$ and Cherenkov branch for $n\beta > 1$. The pure Compton and Cherenkov radiation processes represent the two limit cases of SCC emission. For $\Delta n = 0$, the Compton branch of SCC radiation reduces to the nonlinear Compton scattering given by Eq. (5). To exceed the threshold for the pure Cherenkov radiation in the case when quantum theory is considered, the required Lorentz factor of the particle needs to be

$$\gamma_{\mathrm{Ch}} > \frac{1}{\sqrt{2\Delta n\left(1-\frac{\hbar\omega}{\mathcal{E}}\right)}}, \tag{14}$$

provided $\Delta n > 0$ [28]. This condition corresponds to $F < 0$ for $\theta = 0$ in Eq. (11). In such a case, the Cherenkov branch becomes significant in the SCC process and the Cherenkov photons are emitted along the particle propagation direction within the Cherenkov angle $\cos\theta_{\mathrm{Ch}} = 1/(n\beta)$ [28].

However, even if the threshold for pure Cherenkov radiation is not met, the SCC radiation can still significantly differ from the traditionally considered Compton one [20,40,41], which might be of interest for future particle-field interaction experiments studying the fundamental properties of a quantum vacuum. In the following, we consider the interaction of charged particles (electron, proton) with a strong EM field that will be provided by the next generation of lasers. Using Eq. (10), we have calculated the emission spectra of the SCC

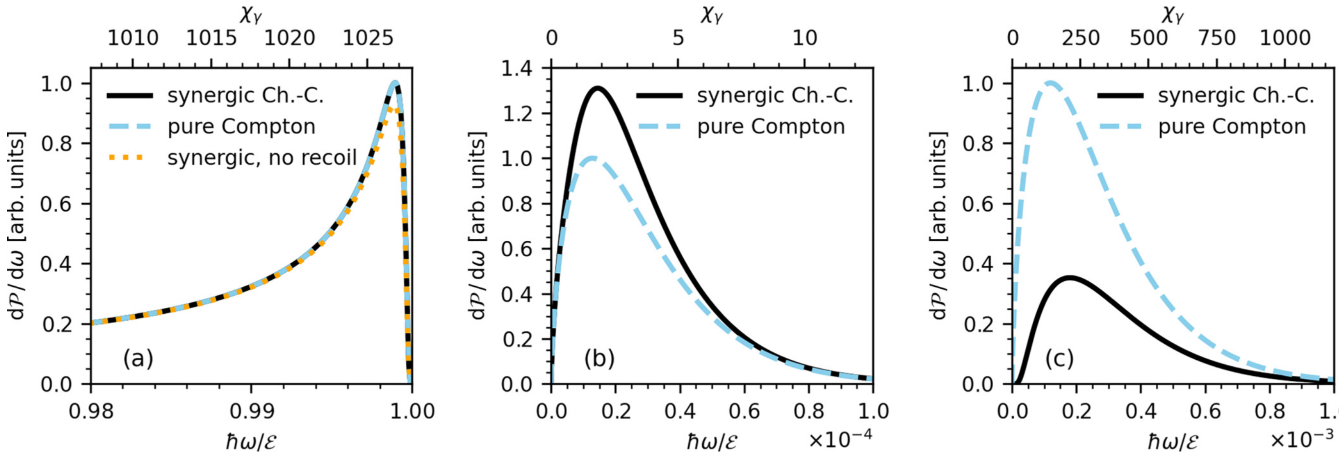

FIG. 1. Power of radiation $d\mathcal{P}/d\omega$, given by Eq. (10), emitted by the (a) 40 GeV electron, (b) 700 GeV, and (c) 6.5 TeV protons in the forward direction during a head-on collision with a strong electromagnetic field of intensity (a) 10^{25} W/cm^2 and (b),(c) 5×10^{26} W/cm^2 in the Cherenkov-Compton process that accounts for the effect of vacuum polarization, i.e., $\Delta n \neq 0$. Solid (dotted) line corresponds to the case when a recoil is (is not) considered in Eq. (11), respectively. The blue dashed curve shows the power emitted in the pure Compton process in which vacuum polarization is neglected, i.e., $\Delta n = 0$.

process for the respective cases, considering for clearness $\Delta n(\chi_\gamma)$ averaged over both polarizations. Under a weak-field approximation, the presented approach can also describe the synchrotron-Cherenkov radiation in a strong constant magnetic field [3,19]. Such an interaction setup is relevant for studying the photon emission in the strong field of magnetars [42] or near black holes [43].

A. Emission by an electron

At first, we consider the interaction of a 40 GeV electron beam with a laser field of intensity 10^{25} W/cm^2, and thus the electrons reach $\chi \approx 10^3$. For achieving such a high laser intensity with multi-petawatt lasers [44–48], one could consider the concept of relativistic flying mirrors [49–52]. For $\chi \gg 1$, a sharp peak close to the initial electron energy appears in the Compton emission spectrum, which means that photons with $\hbar\omega \approx \mathcal{E}$ are predominantly radiated [53,54]. Since for $\chi_\gamma \gg 1$ we get $\Delta n < 0$ from Eq. (13), the photon-emission spectra should be suppressed near the spectral tip. In Fig. 1(a), we show the emission spectra for the following processes: pure Compton (dashed line), SCC (solid line), and SCC with neglected recoil term $1 - \hbar\omega/\mathcal{E}$ in Eq. (11) (dotted line). As shown, there is no visible difference between the emission spectra in the case of pure Compton and SCC radiation. While the recoil term $1 - \hbar\omega/\mathcal{E}$ can be considered insignificant for quantum Cherenkov radiation [28], it is shown that in the more general case of Cherenkov-Compton radiation, this term becomes crucial as it controls the amount of energy emitted due to the effects of vacuum polarization. For this particular case, neglecting this term results in obtaining erroneous results that underestimate the emitted power by approximately 10% near the spectral tip in the case of SCC. We see that the recoil effect thus plays an important role in SCC radiation as it considerably reduces the effect of vacuum polarization on the photon-emission spectra near $\hbar\omega \approx \mathcal{E}$; see Eq. (11). As a result, in the case of the interaction with an ultrarelativistic electron characterized by $\chi \gg 1$, the evidence of photon spectra modification caused by vacuum polarization in a strong EM field is dramatically suppressed just due to the recoil.

For direct observation of the Cherenkov effect, it is reasonable to consider photons characterized by $\chi_\gamma \approx 0.75$, as in this case Δn reaches the maximum value while electron-positron pair production is considerably diminished. Since Δn acquires a positive constant value for $\chi_\gamma \ll 1$ [see Eq. (13)], it is natural to study the Cherenkov branch of SCC radiation of an electron in this limit [22–24,55,56]. We note that to exceed the threshold for pure Cherenkov radiation given by Eq. (14), the Lorentz factor $\gamma \gg 1$ is needed. In such a case, photons with $\chi_\gamma \gg 1$ are predominantly emitted while Cherenkov photons are only emitted at the low-energy end of the spectra ($\chi_\gamma \ll 1$). The detailed analysis of photon emission in this case presented in Ref. [24] shows that the Cherenkov cone is not distinctly evident for the studied range of parameters.

B. Emission by an ion

In contrast to the case of emitting electrons, the Cherenkov branch tends to dominate for heavier particles [18]. As an example, we consider the interaction of a 700 GeV proton, whose energy is an order of magnitude below the current record [57], with the EM field of intensity 5×10^{26} W/cm^2. As shown in Fig. 1(b), the SCC process (solid line) enhances the peak of the photon-emission spectra by 30% in comparison to pure Compton radiation (dashed line) in which vacuum polarization is neglected. When vacuum polarization is taken into account, more energy, by 20%, is emitted in the form of photons with $1 \lesssim \chi_\gamma \lesssim 8$ that, in turn, enables more efficient pair production. We note that in the collision of a high-energy proton beam and a strong laser field, the merging of laser photons can occur due to the polarization of vacuum [58].

The Compton-Cherenkov process can also significantly suppress the emission of photons with $\chi_\gamma \gg 1$. We note that in galactic centers, protons can be accelerated up to PeV range [59]; however, to demonstrate this phenomenon in a terrestrial laboratory, we consider a proton of today's available energy of 6.5 TeV [57]. In such case, photons with $\chi_\gamma \gg 1$ are emitted. Since $\Delta n(\chi_\gamma \gg 1) < 0$, the SCC process causes the quenching of photon emission. In Fig. 1(c), we show that in such case, the peak of the photon-emission spectra is reduced by 65% and the total emitted energy is lower by more than 50% compared to pure Compton scattering. We note that for the presented parameters, the protons are characterized by $\chi \approx 10^{-5}$ and 10^{-4}, respectively, and the formation length for photon emission is of the order of $\lambda/10$, where wavelength $\lambda = 1$ μm is assumed.

IV. CONCLUSION

We calculated the photon-energy distribution in the Compton-Cherenkov process describing the emission of a charged particle in a vacuum polarized by a very strong electromagnetic field. This phenomenon is important for the fundamental QED processes present in strong fields near magnetars and black holes in space and can be studied in collisions of accelerated particles with a strong field produced by the new generation of lasers. While for electrons characterized by $\chi \gg 1$ the effect of vacuum polarization on the emission spectra peak is not expected just due to the recoil, for today's available 6.5 TeV (700 GeV) protons accelerated with the Large Hadron Collider accelerator [57], the synergic nature of photon emission leads to suppression (enhancement) of emitted energy by tens of percent compared to a pure Compton process (i.e., when vacuum polarization effects are neglected). The features of enhancement and reduction of high-energy photon emission in the Compton-Cherenkov process can be exploited as a signature of vacuum texture in laser-particle experiments.

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