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Photo-production of scalar particles in the field of a circularly polarized laser beam

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ABSTRACT

The photo-production of a pair of scalar particles in the presence of an intense, circularly polarized laser beam is investigated. Using the optical theorem within the framework of scalar quantum electrodynamics, explicit expressions are given for the pair production probability in terms of the imaginary part of the vacuum polarization tensor. Its leading asymptotic behavior is determined for various limits of interest. The influence of the absence of internal spin degrees of freedom is analyzed via a comparison with the corresponding probabilities for production of spin-1/2 particles; the lack of spin is shown to suppress the pair creation rate, as compared to the predictions from Dirac theory. Potential applications of our results for the search of minicharged particles are indicated.

1. Introduction

Understanding the nonlinear and unstable nature of the quantum vacuum in the presence of a strong electromagnetic field constitutes an important task of theoretical physics. Corresponding studies have revealed a nontrivial vacuum structure, suitable to explore the low-energy frontier of particle physics [1–3]. Moreover, perspectives of achieving ultrahigh field intensities ($I \sim 10^{26}$ W/cm^2) in short laser pulses of few femtoseconds duration [4,5] have motivated a growing interest in the phenomenology purely associated with the quantum nature of the electromagnetic interaction (see [6,7] for recent reviews). This is because the envisaged laser field strengths lie only 1–2 orders of magnitude below the critical value $E_c = 1.3 \times 10^{16}$ V/cm where QED vacuum nonlinearities become substantial and spontaneous vacuum decay into electron–positron (e^-e^+) pairs via the Schwinger mechanism is expected to occur [8–10].

In combination with an incident high-energy particle, strong laser fields can induce e^-e^+ pair production already at intensities available today. In a pioneering experiment at SLAC [11], a multi-GeV photon decayed into a pair while propagating through a moderately intense laser pulse ($I \sim 10^{18}$ W/cm^2). This process, involving the simultaneous absorption of several laser photons, represents a nonlinear version of the well-known Breit–Wheeler reaction [12–14]. The high-energy non-laser photon originated from Compton backscattering of SLAC's ultrarelativistic electron beam off the laser pulse. In the near future, corresponding studies can be conducted within all-optical setups using laser-accelerated relativistic electrons as projectiles [15]. Other pair production mechanisms may be probed in ultrarelativistic proton–laser collisions [16–20].

In view of the upcoming high-field laboratories [4,5], theoreticians are currently investigating further properties and applications of photo-induced e^-e^+ pair production in intense laser fields. For example, due to their broad frequency composition, laser pulses of ultrashort duration have been shown to modify the created particle spectra [21] and lead to characteristic enhancements in the pair production probability [22,23]. Photo-induced pair production also plays a crucial role for the development of QED cascades which may give rise to e^-e^+ plasmas of very high density [24,25]. The nonlinear Breit–Wheeler process moreover offers a promising means to measure ultrashort γ-ray pulses via e^-e^+ streaking [26]. Superimposing the field of a high-energy photon onto a strong electric field may also help catalyzing the Schwinger effect [27–29].

In the present Letter, we study the photo-induced creation of a pair of spin-0 particles in the presence of a strong monochromatic laser beam. Our motivation is twofold. First, while the probabilities for the creation of fermion pairs are known for a long time [13,14], it is relevant to establish the corresponding formulas for scalar particles because they can be useful for the ongoing search of minicharged particles which may have either fermionic or bosonic character [2,3,30]. Second, our results provide insights into the

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fundamental question as to how the spin degree of freedom affects the photo-induced pair production process. To this end, a comparison with the known results for fermion pair production will be drawn. Such an information complements previous works where spin-resolved calculations of the nonlinear Breit–Wheeler process via helicity amplitudes [31] and the internal spin polarization vector [32] were performed. We note besides that comparative studies between the behavior of bosonic and fermionic particles in strong laser fields have recently been carried out with respect to Compton, Mott and Kapitza–Dirac scattering [6,33,34], nonlinear Bethe–Heitler pair creation in proton–laser collisions [35] and the Klein paradox [36,37].

Our theoretical approach relies on the polarization tensor $\Pi_{\mu\nu}(k_1,k_2)$, of scalar Quantum Electrodynamics (in the one-loop approximation) in the presence of a strong laser field [38,39] whose imaginary part is related to the pair production probability of scalar particles via the optical theorem. While the polarization tensor for Dirac fermions has already been exploited successfully to calculate various e^-e^+ pair production processes in strong laser fields [18,19,38], to the best of our knowledge the present calculations represent the first application of the corresponding polarization tensor for the scalar case.

2. General considerations

To begin with, let us consider the field of a plane electromagnetic wave of the form[1]

$$\mathscr{A}^{\mu}(x) = a_1^{\mu}\psi_1(\varkappa x) + a_2^{\mu}\psi_2(\varkappa x), \tag{1}$$

with $a_{1,2}$ denoting the wave amplitudes and $\psi_{1,2}$ being arbitrary functions. The wave four-vector $\varkappa^{\mu} = (\varkappa^0, \boldsymbol{\varkappa})$ fulfills the relations $\varkappa^2 = 0$ and $\varkappa a_1 = \varkappa a_2 = a_1 a_2 = 0$. According to [38,39], the vacuum polarization tensor in this field,

$$\Pi^{\mu\nu}(k_1,k_2) = c_1\Lambda_1^{\mu}\Lambda_2^{\nu} + c_2\Lambda_2^{\mu}\Lambda_1^{\nu} + c_3\Lambda_1^{\mu}\Lambda_1^{\nu} + c_4\Lambda_2^{\mu}\Lambda_2^{\nu} + c_5\Lambda_3^{\mu}\Lambda_4^{\nu} \tag{2}$$

can be expanded in terms of a basis set of Lorentz covariant vectors Λ_i^{μ} which are constructed from fundamental symmetry principles. They are explicitly given by

$$\Lambda_1^{\mu}(k) = -\frac{\mathscr{F}_1^{\mu\nu}k_{\nu}}{(k\varkappa)(-a_1^2)^{1/2}}, \qquad \Lambda_2^{\mu}(k) = -\frac{\mathscr{F}_2^{\mu\nu}k_{\nu}}{(k\varkappa)(-a_2^2)^{1/2}},$$
$$\Lambda_3^{\mu}(k) = \frac{\varkappa^{\mu}k_1^2 - k_1^{\mu}(k\varkappa)}{(k\varkappa)(k_1^2)^{1/2}}, \qquad \Lambda_4^{\mu}(k) = \frac{\varkappa^{\mu}k_2^2 - k_2^{\mu}(k\varkappa)}{(k\varkappa)(k_2^2)^{1/2}}. \tag{3}$$

Here $\mathscr{F}_i^{\mu\nu} = \varkappa^{\mu}a_i^{\nu} - \varkappa^{\nu}a_i^{\mu}$ $(i=1,2)$ are the amplitudes of the external field modes whereas k_1 and k_2 denote the incoming and outgoing four-momenta of the probe photons, respectively. We note that the short-hand notation k in Eq. (3) may stand for either k_1 or k_2. It is worth mentioning at this point that, for $k=k_1$, the vectors $\Lambda_1(k_1)$, $\Lambda_2(k_1)$ and $\Lambda_3(k_1)$ are orthogonal to each other, $\Lambda_i^{\mu}(k_1)\Lambda_{j\mu}(k_1) = -\delta_{ij}$, and fulfill the completeness relation $g^{\mu\nu} - \frac{k_1^{\mu}k_1^{\nu}}{k_1^2} = -\sum_{i=1}^{3}\Lambda_i^{\mu}(k_1)\Lambda_i^{\nu}(k_1)$ with $g_{\mu\nu} = \mathrm{diag}(+1,-1,-1,-1)$ denoting the metric tensor. A similar statement applies if the set of vectors $\Lambda_1(k_2)$, $\Lambda_2(k_2)$ and $\Lambda_4(k_2)$ are considered. We emphasize that Eq. (2) does not depend on which choice of k is taken since the difference between k_1 and k_2 is proportional to $\varkappa$ [see Eq. (5) below].

[1] From now on "natural" and Gaussian units $c = \hbar = 4\pi\epsilon_0 = 1$ are used.

The form factors c_i in Eq. (2) are distribution-valued functions which depend on the field shape via the functions ψ_i. They have been evaluated thoroughly for the case of spin-$\frac{1}{2}$ particles in [38,39]. Also for the case when the virtual charge carriers in the Feynman loop are spin-0 particles general expressions for the c_i were provided in these references; but these formulas were not further evaluated.

Using the general expressions from [38,39] and assuming that the laser field is an elliptically polarized wave with

$$\psi_1 = \cos(\varkappa x) \quad \text{and} \quad \psi_2 = \sin(\varkappa x), \tag{4}$$

we find that the form factors in Eq. (2) for the scalar case are given by

$$c_i = -i\frac{\alpha}{\pi}m^2\int_{-1}^{1}dv\int_0^{\infty}\frac{d\rho}{\rho}e^{-\frac{2i\rho}{|\lambda|(1-v^2)}\left(1+A(\xi_1^2+\xi_2^2)-\frac{k_1^2(1-v^2)}{4m^2}\right)} \times (2\pi)^4\left[\delta^4(k_1-k_2)d_i^{(0)} + \sum_{\substack{N=-\infty\\N\neq 0}}^{\infty}\delta^4(k_1-k_2-2N\varkappa)d_i^{(N)}\right]. \tag{5}$$

Here, $\alpha = e^2$ is the fine structure constant, e and m denote the particle charge and mass, respectively, and

$$\lambda = \frac{\varkappa k}{2m^2}, \qquad \xi_i^2 = -\frac{e^2a_i^2}{m^2} \quad (i=1,2). \tag{6}$$

As Eq. (5) shows, the polarization tensor decomposes into elastic $(k_1=k_2)$ and inelastic $(k_1\neq k_2)$ parts. Those terms which contain the Dirac deltas $\delta^4(k_1-k_2+2N\varkappa)$ with $N\neq 0$ are responsible for the inelastic scattering of a photon in the field of the wave. For our purposes, however, only the elastic part is relevant. The corresponding functions $d_i^{(0)}$, $i=1,2,3,4,5$, contained in Eq. (5) are given by

$$d_1^{(0)} = -d_2^{(0)} = \xi_1\xi_2\rho A_0 J_0(z)\,\mathrm{sign}[\lambda], \tag{7}$$

$$d_3^{(0)} = -\frac{1}{2}\xi_1^2A_1\left(J_0(z) - iJ_0'(z)\right) + \frac{\xi_1^2}{2}\sin^2(\rho)J_0(z) + \frac{i|\lambda|(1-v^2)}{8\rho}\left(J_0(z) - e^{iy}\right), \tag{8}$$

$$d_4^{(0)} = d_3^{(0)}\left(\xi_1^2 \leftrightarrow \xi_2^2\right), \qquad d_5^{(0)} = -\frac{k_1^2}{8m^2}v^2\left(J_0(z) - e^{iy}\right) \tag{9}$$

where $J_0(z)$ is the Bessel function of zero order and $J_0'(z)$ its derivative. The remaining parameters are

$$A = \frac{1}{2}\left(1 - \frac{\sin^2(\rho)}{\rho^2}\right), \qquad A_0 = \frac{1}{2}\left(\frac{\sin^2(\rho)}{\rho^2} - \frac{\sin(2\rho)}{2\rho}\right),$$
$$z = \frac{2\rho A_0}{1-v^2}\frac{\xi_1^2-\xi_2^2}{|\lambda|}, \qquad y = \frac{2\rho A}{1-v^2}\frac{\xi_1^2+\xi_2^2}{|\lambda|}, \tag{10}$$

and $A_1 = A + 2A_0$.

A substantial simplification is achieved when the external field is taken as a circularly polarized wave ($\xi_1 = \xi_2 = \xi$). In this case, we find it convenient to express the elastic contribution as

$$\Pi_{\mu\nu}^{(\mathrm{elast})}(k_1,k_2) = i(2\pi)^4\delta^4(k_1-k_2)\Pi_{\mu\nu}(k_2) \tag{11}$$

with

$$\Pi^{\mu\nu}(k_2) = \left(\Lambda_1^{\mu}\Lambda_2^{\nu} - \Lambda_2^{\mu}\Lambda_1^{\nu}\right)\pi_1^{(0)} + \left(\Lambda_1^{\mu}\Lambda_1^{\nu} + \Lambda_2^{\mu}\Lambda_2^{\nu}\right)\pi_3^{(0)} + \Lambda_3^{\mu}\Lambda_3^{\nu}\pi_5^{(0)}. \tag{12}$$