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Pair production in short laser pulses near threshold

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ABSTRACT

The e^+e^- pair production by a probe photon traversing a linearly polarized laser pulse is treated as generalized nonlinear Breit–Wheeler process. For short laser pulses with very few oscillations of the electromagnetic field we find below the perturbative weak-field threshold $\sqrt{s} = 2m$ a similar enhancement of the pair production rate as for circular polarization. The strong enhancement below the weak-field threshold is traced back to the finite bandwidth of the laser pulse. A folding model is developed which accounts for the interplay of the frequency spectrum and the intensity distribution in the course of the pulse.

1. Introduction

Strong-field QED processes, in particular pair production processes, in short intense laser pulses are currently of great interest in view of upcoming high-power laser facilities such as HiPER [1] or ELI [2]. Strong electromagnetic fields produced by lasers polarize the quantum vacuum. This vacuum polarization can be tested by probe photons leading to pair production, provided the probe photon has a sufficiently high energy to separate the virtual quantum fluctuations lifting the particles to their mass shell.

Pair production in photon–photon collisions is described via the Breit–Wheeler process [3], $\gamma' + \gamma \to e^+ + e^-$, as t channel process which is sometimes termed "conversion of light into matter". Considering it as 2-to-2 process in a particle physics language, the Breit–Wheeler pair creation is a threshold process, i.e. the available energy, expressed by the Mandelstam variable $s = (k' + k)^2 = 2\omega'\omega(1 - \cos\theta_{\vec{k}'\vec{k}})$, must fulfill $s > s_{\text{thr}} \equiv 4m^2$, where m is the electron mass. Supposed, γ describes optical (laser) photons with energy $\omega \sim 1$ eV, the energy of the counter propagating photon γ' must obey $\omega' > 250$ GeV. While such high energy photons might be generated by various processes at high energy accelerators/colliders, e.g. in the time reversed Breit–Wheeler process (pair annihilation) or Compton backscattering of a laser beam off a very energetic electron beam, or in astrophysical environments, the availability of laboratory based 250 GeV photon sources is quite scarce. Nevertheless, the Breit–Wheeler process has been identified experimentally. The SLAC experiment E-144 [4] was a special set-up of one of the above mentioned options.

In a sufficiently strong laser field, however, multi-photon processes are enabled [5–8], schematically $\gamma' + n\gamma \to e^+ + e^-$, often termed nonlinear Breit–Wheeler process. Indeed, the interpretation [9] of the SLAC experiment E-144 has proved this possibility. In fact, at least $n = 5$ laser photons were needed to produce a pair [10]. For an all-optical set-up with $\omega' \sim \omega$, $n > 10^{11}$ laser photons would be required to provide the threshold energy. Strong laser fields are necessary for such nonlinear effects. The electron and positron in a strong laser field acquire an effective mass, $m_* = m\sqrt{1 + a_0^2/2}$, where a_0 is the dimensionless laser-strength parameter related to the laser intensity I and wavelength λ as $a_0^2 = 7.309 \times 10^{-19} I[\text{W/cm}^2]\lambda^2[\mu\text{m}]$ (cf. [11–13] for a recent analysis of the mass dressing in strong pulsed laser fields). That is, the above threshold estimate would read $s > s^*_{\text{thr}} \equiv 4m_*^2$ in disfavor of achieving the threshold for strong optical laser fields. The up-shift of the threshold, however, is compensated to some extent by the higher harmonics, i.e. multi-photon channels. The n-photon channels are open at $s > s_n = 4m_*^2/n$. That means, in principle, the higher harmonics (labeled by n) lead to pair production for arbitrarily small values of $s < s_{\text{thr}}$, termed subthreshold production. For a given value of s, the minimum number of laser photons n_0 is the smallest integer such that $s_{n_0} \leqslant s$; channels with $n < n_0$ photons are closed. In the perturbative regime, where $a_0 \ll 1$, the harmonics are suppressed by factors of a_0^{2n}, and a noticeable

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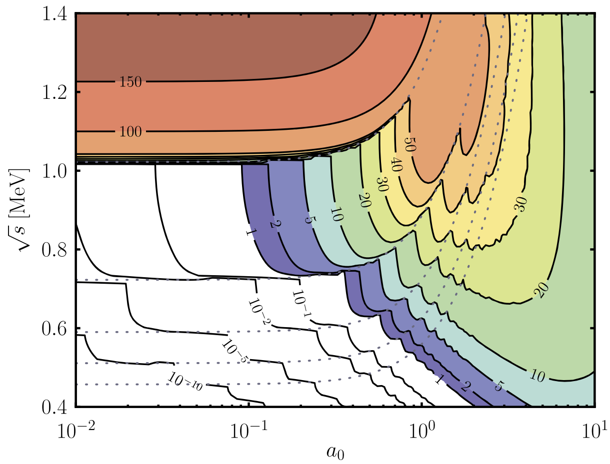

Fig. 1. Contour plot of the cross section (1) over the $\sqrt{s}$ vs. a_0 plane. The numbers label curves of constant cross sections in mb. Dotted curves depict the loci of closing the n-photon channels at $s = s_n$ for $n = 1 \dots 5$ (from top to bottom).

pair production therefore ceases shortly below the 2-to-2 threshold of $s_{\rm thr}$. Denoting by $\kappa = s a_0/(2m^2)$ the Ritus parameter for pair production [7], in the deep subthreshold region $\kappa \ll 1$ and at large values of $a_0 \gg 1$, a large number of laser photons participate in the formation of pairs with cross section $\sigma \propto e^{-8/(3\kappa)}$. It has a similar non-analytic dependence on the field strength parameter encoded in κ as the Schwinger effect which depends on the electric field amplitude (cf. [14]). At fixed a_0, there is a steep decline of the pair creation probability for $s < 2m^2/a_0$, i.e. large values of a_0 shift the region of noticeable pair production to small values of s.

The cross section for pair production in the nonlinear Breit–Wheeler process by an unpolarized probe photon γ' in a plane wave with linear polarization reads in such a case [7] $\sigma(s) = \sum_{n \geqslant n_0}^{\infty} \sigma_n(s)$ with

$$\sigma_n(s) = \frac{4\alpha^2}{s a_0^2} \int_0^{2\pi} d\varphi \int_1^{u_n^*} \frac{du}{u\sqrt{u(u-1)}} \times \left\{ A_0^2 + a_0^2(2u-1)\left(A_1^2 - A_0 A_2\right) \right\}, \tag{1}$$

where $A_j \equiv A_j(n, a, b) = \frac{1}{2\pi}\int_{-\pi}^{\pi} dt \cos^j t \exp\{i(a \sin t - b \sin 2t - nt)\}$ are generalized Bessel functions with arguments n, $a = \frac{8m^2}{s} a_0 \sqrt{1 + a_0^2/2}\sqrt{u(u_n^* - u)} \cos\varphi$, and $b = a_0^2 u m/s$; u_n^* is defined by $u_n^* = n\frac{s}{4m_*^2}$. Note that the relevant energy variable is still $s = (k' + k)^2$. In Fig. 1, a survey of the cross section is exhibited as a contour plot over the $\sqrt{s}$ vs. a_0 plane. The apparent structures visible in the region $0.8 < \sqrt{s}\,[\mathrm{MeV}] < 1$ for $a_0 < 1$ are related to the turn over from the first to the second harmonic. The features propagate to increasing values of $\sqrt{s}$ for $a_0 > 1$. The imprints of the onset of the second and third harmonics are still visible. In the weak-field region, $a_0 < 1$, the steep decline of the cross section as a function of $\sqrt{s}$ is evident below $\sqrt{s} = 2m$.

In the case of circular polarization, Eq. (1) becomes modified [7,8]. The overall pattern remains, but with less pronounced structures due to azimuthal symmetry. In the strong-field region, $a_0 \gg 1$, the subthreshold production with a large number of photons is suppressed due to the large angular momentum of the pair [7].

The hitherto discussed examples are for infinitely extended plane waves. In reality, strong laser fields are generated presently by the chirped pulse amplification technique, i.e. pulse compression, and the asymptotic final state refers to a free electron–positron pair, where m_* does not matter. Therefore, a modified threshold behavior is to be expected. Moreover, a pulse of finite duration is not longer monochromatic, instead the power spectrum has a support of finite width. Also this effect will have imprints on the pair production, as does the variation of the intensity in the course of the pulse. Considering the (nonlinear) Breit–Wheeler process as cross channel of the (nonlinear) Compton scattering one can expect similar strong effects of the temporal pulse shape, as has been found for the latter one [15–18]. In fact, the differential e^+e^- spectra in the (general) Breit–Wheeler process depend sensitively on the laser pulse shape [11]. Even for weak laser intensities $a_0 < 1$, a significant enhancement of the total pair production rate just below the threshold $s_{\rm thr}$ has been found recently [19] for short circularly polarized pulses.

Given this motivation we study here the pair production off a probe photon in a linearly polarized laser pulse. In Section 2, we recapitulate the basic formulas within the Furry picture, where the interaction with the probe photon γ' is treated perturbatively, while the interaction of the electron and positron with the laser pulse is accounted for by Volkov states. The numerical results are discussed in Section 3 for the total pair production probability. A folding model is introduced in Section 4 to access to the effect of enhanced subthreshold pair production, thus identifying the relevant physical mechanisms. The summary is given in Section 5.

2. Pair production in pulsed laser fields

The pair creation process is a cross channel of the Compton scattering. Accordingly, we evaluate the cross section of the process as decay of a probe photon γ' with four-momentum k' and polarization four-vector ϵ' into a pair $e^+(\gamma) + e^-(\gamma)$, where $e^\pm(\gamma)$ are laser dressed Volkov states (cf. [20]) which encode the interaction with the laser field. For weak laser fields, $a_0 \ll 1$, this process may be resolved perturbatively into a series of diagrams, where, in addition to the probe photon γ', n laser photons γ interact with the outgoing $e^\pm$. The diagram with $n = 1$ at $a_0 \ll 1$ reproduces the Breit–Wheeler process.

The linearly polarized laser pulse is described here by the real four-vector potential

$$A^\mu(\phi) = \frac{m a_0}{e} \epsilon^\mu g(\phi) \cos\phi \tag{2}$$

with transverse polarization four-vector ϵ^μ obeying $k \cdot \epsilon = 0$; e denotes the elementary charge. The envelope function $g(\phi)$, with $\phi = k \cdot x$ as invariant phase, encodes the temporal shape of the pulse with wave four-vector k and normalization $g(0) = 1$. We chose in this Letter, following [11],

$$g(\phi) = \cos^2\left(\frac{\phi}{2N}\right), \tag{3}$$

for $|\phi| \leqslant \pi N$ and $g(\phi) = 0$ for $|\phi| > \pi N$, where N is number of cycles in the pulse. In line with the notation in [19] we denote the considered process as finite pulse approximation (FPA), in contrast to the infinite pulse approximation (IPA) where $g \to 1$. Both, IPA and FPA, use plane waves, i.e. $\epsilon \cdot k = 0$.

The S matrix for the generalized nonlinear Breit–Wheeler process is

$$S = -ie \int d^4x\, \bar{\Psi}_{p'}(x) \not{\epsilon}' \frac{e^{-ik' \cdot x}}{\sqrt{2\omega'}} \Psi_{-p}(x), \tag{4}$$

where $\bar{\Psi}_{p'}$ and Ψ_{-p} denote the Volkov wave functions (to be taken with the potential (2)) of the outgoing electron and positron with four-momenta p' and p, respectively. For instance, the positron wave function is given by