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Strong-field Breit–Wheeler pair production in two consecutive laser pulses with variable time delay

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ABSTRACT

Photoproduction of electron–positron pairs by the strong-field Breit–Wheeler process in an intense laser field is studied. The laser field is assumed to consist of two consecutive short pulses, with a variable time delay in between. By numerical calculations within the framework of scalar quantum electrodynamics, we demonstrate that the time delay exerts a strong impact on the pair-creation probability. For the case when both pulses are identical, the effect is traced back to the relative quantum phase of the interfering *S*-matrix amplitudes and explained within a simplified analytical model. Conversely, when the two laser pulses differ from each other, the pair-creation probability depends not only on the time delay but, in general, also on the temporal order of the pulses.

1. Introduction

The generation of matter–antimatter particle pairs from the electromagnetic energy of photons belongs to the most striking predictions of quantum electrodynamics (QED). It can be realized through the strong-field Breit–Wheeler (SFBW) reaction,

$$\omega_\gamma + n\omega \to e^+e^- \,, \qquad (1)$$

where a high-energy gamma ray of frequency ω_γ collides with a high-intensity laser field from which n photons of frequency ω are absorbed to overcome the pair creation threshold. Experimental evidence for SFBW pair production was found in highly relativistic electron–laser collisions at SLAC [1]. In the foreseeable future, further studies of the process are planned at high-intensity laser facilities such as the Extreme-Light Infrastructure [2], the Exawatt Center for Extreme Light Studies [3], the Diocles Petawatt Laser [4] or the HIBEF project [5]. These campaigns are going to cover large areas of the parameter space for SFBW which have not been probed yet. As an alternative experimental approach to the Breit–Wheeler process, the usage of a thermal photon target has been proposed [6,7].

Since high laser intensities are generated in short pulses, theoreticians have started a few years ago to calculate pair production in laser fields of finite extent. With respect to the SFBW process, it was found that the broad frequency spectrum of a short pulse can strongly modify the energy and angular distributions of created particles [8–16]. In particular, the carrier-envelope phase of a few-cycle pulse was shown to exert a characteristic impact [17,18]. Besides, when several laser pulses follow each other, their partial contributions add up coherently, leading to a comb-like structure of emitted positrons [19]. In certain laser parameter domains, the spectral broadness of a short pulse may also strongly affect the total pair creation probability due to subthreshold enhancement effects [20,21]. In laser-driven recollisions of a created electron and positron, even muon–antimuon pair production can result as a subsequent high-energy reaction [22,23].

The impact of finite pulses on pair production was also analyzed in other electromagnetic field configurations, such as time-dependent electric fields of finite duration or spatially localized electric and magnetic fields (see, e.g., [24–27]). In particular, multiple-slit interference phenomena in the time domain were observed in sequences of electric-field pulses [28–30]. Moreover, strong enhancement effects have been predicted when a rather weak, but fast oscillating field component is superimposed onto an intense, slowly varying field [31–39]. Systematic analyses to find the pulse shapes which optimize the pair yields were carried out [40–42]. Also the creation of multiple pairs in electromagnetic fields of finite extension was addressed [43,44].

In this paper, we study SFBW pair production in a laser field which consists of two consecutive pulses, see Fig. 1. Our focus lies on effects arising from variations of the time delay between both pulses. Two scenarios are considered: When both pulses are identical, the time delay is shown to strongly influence the energy spectrum of created particles and, remarkably, the total produc-

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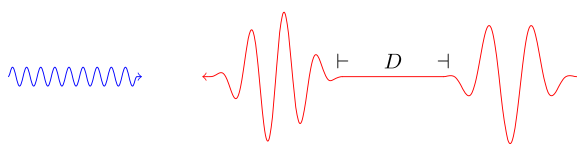

Fig. 1. Scheme of the field configuration: The gamma quantum (blue) collides with two consecutive short laser pulses (red) with a variable distance D. (For interpretation of the references to color in this figure legend, the reader is referred to the web version of this article.)

tion probability, as well. Our numerical results are corroborated by a simplified analytical model. When the pulses are different, the interesting question arises in addition whether their temporal sequence can affect the pair creation process. As we will show, in general the pulse order matters.

Gaussian units with $\hbar = c = 1$ are employed throughout. The positron charge and mass are denoted by e and m, respectively, and $\lambda_e = 1/m$ is the reduced Compton wavelength.

2. Theoretical framework

The SFBW process is induced by the decay of a high-energy photon, which is described as one mode $\{\mathbf{k}_\gamma, \lambda_\gamma\}$ of a quantized radiation field $\hat{\mathcal{A}}^\mu$. Effectively, we employ the scattering potential

$$\mathcal{A}^\mu_\gamma = \langle 0|\hat{\mathcal{A}}^\mu|\mathbf{k}_\gamma\lambda_\gamma\rangle = \sqrt{\frac{2\pi}{V\omega_\gamma}}\,e^{-ik_\gamma\cdot x}\epsilon^\mu_\gamma\,, \tag{2}$$

with the wave four-vector $k^\mu_\gamma = (\omega_\gamma, \mathbf{k}_\gamma)$ and a real polarization vector ϵ^μ_γ fulfilling $k_\gamma\cdot\epsilon_\gamma = 0$ and being referenced by a mode index λ_γ. We use similar notation and conventions as in [45].

The two consecutive laser pulses are described classically by means of their combined vector potential

$$\mathcal{A}^\mu = \mathcal{A}^\mu_1 + \mathcal{A}^\mu_2\,, \tag{3}$$

where each of the single pulses is of the form $\mathcal{A}^\mu_j = \mathcal{A}^\mu_j(\phi_j) = a_j f_j(\phi_j - \delta_j)\chi_{[0,2\pi]}(\phi_j - \delta_j)\epsilon^\mu_j$, with the amplitude parameter a_j and phase variable $\phi_j = k_j\cdot x$ for $j = 1, 2$. The actual shape is determined by the shape functions f_j and the characteristic function $\chi_{[0,2\pi]}(\phi)$ which is unity for $0 \le \phi \le 2\pi$ and zero otherwise. The wave four-vectors k_1 and k_2 fulfill $k_1\cdot k_2 = 0$ and $\epsilon_j\cdot k_j = 0$, and ϵ^μ_j is a real polarization four-vector. The phase-shift parameters $\delta_j \ge 0$ are chosen such that the pulses are strictly separated. The particle states in the combined laser field $\mathcal{A}$ can be described by Gordon–Volkov solutions $\Psi_{p_\pm}$ (see, e.g., Eq. (5) in [45]). For calculational simplicity, the high-energy photon is assumed to collide head-on with the laser pulses.

Our calculations are performed within scalar QED, disregarding the electron and positron spin. This simplification helps us to render the main physical content of our study more transparent. Note that, in general, there can be significant differences between the creation of Klein–Gordon versus Dirac pairs, in particular on the basis of fully differential production probabilities. However, in terms of total probabilities, these differences diminish and reduce to an overall factor of about 3–5 for SFBW pair production in short laser pulses of moderate intensity [45]. Also in the strong-field limit, the production rates of scalar and fermion pairs are known to coincide with each other, up to on overall prefactor [46]. In the present paper, we shall mostly consider total production probabilities in double pulses, which are set into relation with the corresponding probabilities in single pulses. The basic influence from the double-pulse structure of the laser field can, thus, be expected to hold qualitatively for Dirac particles, as well. For further recent studies of strong-field pair creation within scalar QED, we refer to [25,29,35,36].

The pair-creation amplitude is obtained from the S matrix

$$\mathcal{S}_{p_+p_-} = -i\int d^4x\,\Psi^*_{p_-}\,\mathcal{H}_{\rm int}\,\Psi_{p_+}\,, \tag{4}$$

with $\mathcal{H}_{\rm int} = -ie\left(\mathcal{A}_\gamma\cdot\overset{\rightarrow}{\partial} - \overset{\leftarrow}{\partial}\cdot\mathcal{A}_\gamma\right) - 2e^2\mathcal{A}\cdot\mathcal{A}_\gamma$ being the interaction Hamiltonian. The S matrix can be brought into the form

$$\mathcal{S}_{p_+p_-} = S_0\int d^4x\, C\, e^{-iQ\cdot x - iH}\,, \tag{5}$$

with $S_0 = iem\sqrt{\frac{\pi}{2V^3E_{p_+}E_{p_-}\omega_\gamma}}$ and the combined momentum vector $Q^\mu = k^\mu_\gamma - (p^\mu_+ + p^\mu_-)$. The reduced matrix element $C = C_0 + \sum_{j=1}^2 C_j$ contains the terms $C_0 = \frac{p_- - p_+}{m}\cdot\epsilon_\gamma$ and $C_j = \frac{2e\mathcal{A}_j(\phi_j)}{m}\cdot\epsilon_\gamma$. The auxiliary function $H = H_1 + H_2$ can be decomposed into contributions from the individual pulses

$$H_j = \int_0^{\phi_j}\sum_{l=1}^{2} h_{l,j}\, f^l_j(\phi - \delta_j)\chi_{[0,2\pi]}(\phi - \delta_j)\,d\phi\,, \tag{6}$$

with $h_{1,j} = -ea_j\left[\frac{\epsilon_j\cdot p_+}{k_j\cdot p_+} - \frac{\epsilon_j\cdot p_-}{k_j\cdot p_-}\right]$ and $h_{2,j} = -\frac{1}{2}e^2a_j^2\left[\frac{1}{k_j\cdot p_+} + \frac{1}{k_j\cdot p_-}\right]$. For $\phi_j > \delta_j + 2\pi$, the value of H_j is constant and denoted as $H^\star_j$.

Switching to light-cone coordinates with $x^- = x^0 - x^\parallel$ and $x^+ = \frac{1}{2}(x^0 + x^\parallel)$, where $x^\parallel = \mathbf{k}_j\cdot\mathbf{x}/k^0_j$, we obtain

$$\mathcal{S}_{p_+p_-} = (2\pi)^3 S_0\delta(Q^-)\delta^{(2)}(\mathbf{Q}^\perp)\int dx^-\, C\, e^{-iQ^0x^- - iH}\,. \tag{7}$$

The remaining integral requires a regularization in analogy to the treatment presented in App. B of [47]. Effectively, we have to replace C in Eq. (7) by the new matrix element $\tilde{C} = \tilde{C}_1 + \tilde{C}_2$ where each part $\tilde{C}_j = C_j - \frac{k^0_j}{Q^0}\frac{dH_j}{d\phi_j}C_0$ contains a characteristic function. The pair-creation probability for unpolarized gamma quanta is obtained as $\mathcal{P} = \frac{1}{2}\sum_{\lambda_\gamma}\int\frac{V d^3p_+}{(2\pi)^3}\int\frac{V d^3p_-}{(2\pi)^3}|\mathcal{S}_{p_+p_-}|^2$.

The pair-creation amplitude $\mathcal{S}_{p_+p_-}$ shall now be decomposed into contributions from the individual pulses, revealing the explicit dependence on the phase-shift parameters and thus the nature of the interaction. To this end, we apply the formal substitution $x^- = (\Phi_j + \delta_j)/k^0_j$ to the integrals

$$I_j = \int dx^-\tilde{C}_j e^{-iQ^0x^- - iH_j} \tag{8}$$

in order to shift the integration domains to $[0, 2\pi]$. Accordingly, we can separate the dependence on δ_j and obtain

$$I_j = F_j\, e^{-iQ^0\delta_j/k^0_j} \tag{9}$$

with $F_j = \frac{1}{k^0_j}\int_0^{2\pi} d\Phi_j\,\tilde{C}_j e^{-iQ^0\Phi_j/k^0_j - iH_j}$ being independent of δ_j, since we can rewrite H_j inside the integration domain of F_j as $H_j = \sum_{l=1}^2 h_{l,j}\int_0^{\Phi_j} f^l_j(\tilde{\Phi}_j)\,d\tilde{\Phi}_j$, with $\Phi_j = \phi_j - \delta_j$. Similarly, $\tilde{C}_j$ is a function of $f_j(\Phi_j)$. This way, the combined amplitude can be brought into the form

$$\mathcal{S}_{p_+p_-} = (2\pi)^3 S_0\,\delta(Q^-)\delta^{(2)}(\mathbf{Q}^\perp)\left(F_1 + F_2\, e^{-i\varphi}\right)\,. \tag{10}$$

Here, $\delta_1 = 0$ was chosen without loss of generality. The contributions of the individual pulses to the pair-creation amplitude are given by F_j. The dynamical phase

$$\varphi = H^\star_1 + Q^0\Delta \tag{11}$$

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