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SELF-CONSISTENT DESCRIPTION OF $e^-e^+\gamma$-PLASMA CREATED FROM THE VACUUM IN A STRONG ELECTRIC LASER FIELD

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In the present work a closed system of kinetic equations is obtained for the description of the vacuum creation of an electron–positron plasma and secondary photons due to a strong laser field. An estimate for the photon energy distribution is obtained. In the Markovian approximation the photon distribution has a $1/k$ spectrum (flicker noise).

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INTRODUCTION

Up to now, the Schwinger effect [1] has resisted an experimental verification. This is basically due to the huge critical field strength E_c which could not yet be reached in the laboratory. Recently, main attention was devoted to theoretical studies of pair creation by time-varying electric fields [2–4] where sufficiently strong electric fields can be achieved at modern high-intensity laser facilities. It has been shown [2–5] that pair creation by a single laser pulse with $E \ll E_c$ could hardly be observed. More optimistic results have been obtained for X-ray free electron lasers [6–8] and for counter-propagating beams of optical lasers [9–11]. It is obvious, that for subcritical fields $E \ll E_c$, the electron–positron excitations have quasiparticle character, and S-matrix methods cannot be applied [12], and existing estimates [8] are not reliable. An adequate method is the kinetic theory.

We will construct here the system of kinetic equations for a self-consistent description of the electron–positron–photon system generated from the vacuum

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by a time-dependent electric field. As a first step, we will consider the one-photon annihilation process only, which in the presence of a strong field is not forbidden [13].

Below we will assume an external electric field $A^\mu(t) = (0, \mathbf{A}_{\rm ext}(t))$ that is spatially homogeneous. In the case of a strong external electric field $\mathbf{A}_{\rm ext}(t)$ also some internal field $\mathbf{A}_{\rm int}(t)$ will be generated. The total acting field will be equal to $\mathbf{A}(t) = \mathbf{A}_{\rm int}(t) + \mathbf{A}_{\rm ext}(t)$ and this field is quasi-classical. Some fluctuations of the electromagnetic field can arise against this background. They can be interpreted as photon excitations. These photons, in principle, can be registered far from the active zone.

1. ELECTRON–POSITRON SECTOR

In general, the complete system of equations for a self-consistent description of the electron–positron–photon plasma consists of: 1) the KEs for the electron and positron quasiparticle components with distribution functions $f_{e,p}(\mathbf{p}, t)$ in the presence of a total electric field $\mathbf{E} = -\dot{\mathbf{A}}(t)$, 2) the KE for the photon component, and 3) the Maxwell equation for the internal field $\mathbf{A}_{\rm int}(t)$. We assume the electroneutrality condition $f_e(\mathbf{p}, t) = f^c(-\mathbf{p}, t) = f(\mathbf{p}, t)$ to be fulfilled.

We start from the standard QED Lagrangian $\mathcal{L} = \mathcal{L}_0 + \mathcal{L}'$ taking into account in $\mathcal{L}_0$ the interaction with a quasi-classical ($A_\mu(t)$) and in $\mathcal{L}'$ with a quantized ($\hat{A}_\mu(t)$) electromagnetic field,

$$\mathcal{L}_0 = \frac{i}{2}\{\overline{\psi}\gamma^\mu D_\mu \Psi - (D^*_\mu \overline{\psi})\gamma^\mu \Psi\} - m\overline{\psi}\psi, \quad \mathcal{L}' = -e\overline{\psi}\gamma^\mu \hat{A}_\mu \Psi, \tag{1}$$

where $D_\mu = \partial_\mu + ieA_\mu(t)$. It is assumed that the intensity of the quantized field is rather weak so that there is no backreaction influence on the state of the system. In other words, the electron–positron system plays the role of a photon source only. The in-vacuum $|{\rm in}\rangle = |\rangle$ is defined such that $\langle \hat{A}_\mu(x)\rangle = 0$. Below we will not consider the backreaction problem because for subcritical fields $E \ll E_c$ the internal field is negligible.

The kinetics of electron–positron vacuum pair creation due to a linearly polarized electric field was studied in a large number of works. See, e.g., [14, 15] and the works quoted therein. The corresponding generalization to the case of an arbitrarily polarized time-dependent electric field was obtained in [16–18]. The oscillator representation in its different realizations [14, 15] leads to the nonstationary orthonormalized spinor basis

$$\begin{aligned} u_1^+(\mathbf{p}, t) &= A(\mathbf{p})\left[\omega_+, 0, p^3, p_-\right], \qquad & u_2^+(\mathbf{p}, t) &= A(\mathbf{p})\left[0, \omega_+, p_+, -p^3\right], \\ v_1^+(-\mathbf{p}, t) &= A(\mathbf{p})\left[-p^3, -p_-, \omega_+, 0\right], \quad & v_2^+(-\mathbf{p}, t) &= A(\mathbf{p})\left[-p_+, p^3, 0, \omega_+\right], \end{aligned} \tag{2}$$

where $\omega(\mathbf{p},t)=\sqrt{m^2+\mathbf{P}^2}=\omega$, $\mathbf{P}=\mathbf{p}-e\mathbf{A}$, $p_\pm=p^1\pm ip^2$, $\omega_+=\omega+m$ and $A(\mathbf{p})=[2\omega\omega_+]^{-1/2}$.

Finally, the Dirac equation in the presence of an external quasiparticle field $\mathbf{A}_{\rm ext}(t)$ generates the Heisenberg-like equations of motion for the construction operators

$$\begin{aligned}\dot a(\mathbf{p},t)&=-U_{(1)}(\mathbf{p},t)a(\mathbf{p},t)-U_{(2)}(\mathbf{p},t)b^+(-\mathbf{p},t)-i\omega(\mathbf{p},t)a(\mathbf{p},t),\\ \dot b(-\mathbf{p},t)&=b(-\mathbf{p},t)U_{(1)}(\mathbf{p},t)+a^+(\mathbf{p},t)U_{(2)}(\mathbf{p},t)-i\omega(\mathbf{p},t)b(-\mathbf{p},t),\end{aligned}\tag{3}$$

where in the representation (2) the matrices are $U_{(1)}(\mathbf{p},t)=i\omega a[\mathbf{pE}]\boldsymbol{\sigma}=iU_k\sigma_k$ and $U_{(2)}(\mathbf{p},t)=a[\mathbf{P}(\mathbf{PE})-\mathbf{E}\omega\omega_+]\boldsymbol{\sigma}$ with $a=e/(2\omega^2\omega_+)$.

KEs for the electron–positron component of the plasma follow from the equations of motion (3) and the definitions of the electron and positron distribution functions in the instantaneous representation

$$f_{\alpha\beta}(\mathbf{p},t)=\langle a^+_\beta(\mathbf{p},t)a_\alpha(\mathbf{p},t)\rangle,\quad f^c_{\alpha\beta}(\mathbf{p},t)=\langle b_\beta(-\mathbf{p},t)b^+_\alpha(-\mathbf{p},t)\rangle\tag{4}$$

and also the two additional functions

$$f^{(+)}_{\alpha\beta}(\mathbf{p},t)=\langle a^+_\beta(\mathbf{p},t)b^+_\alpha(-\mathbf{p},t)\rangle,\quad f^{(-)}_{\alpha\beta}(\mathbf{p},t)=\langle b_\beta(-\mathbf{p},t)a_\alpha(\mathbf{p},t)\rangle\tag{5}$$

describing vacuum polarization. The system of KEs is then [16–18]

$$\begin{aligned}\dot f&=[f,U_{(1)}]-\big(U_{(2)}f^{(+)}+f^{(-)}U_{(2)}\big),\\ \dot f^c&=[f^c,U_{(1)}]+\big(f^{(+)}U_{(2)}+U_{(2)}f^{(-)}\big),\\ \dot f^{(+)}&=[f^{(+)},U_{(1)}]+\big(U_{(2)}f-f^cU_{(2)}\big)+2i\omega f^{(+)},\\ \dot f^{(-)}&=[f^{(-)},U_{(1)}]+\big(fU_{(2)}-U_{(2)}f^c\big)-2i\omega f^{(-)}.\end{aligned}\tag{6}$$

If the standard decomposition in the basis of Pauli matrices is used $f=f_0+f_k\sigma_k$, where $f_0=\mathrm{Tr}\,\{f\}/2$, and $f_k=\mathrm{Tr}\,\{f\sigma_k\}/2$, the system KEs (6) can be rewritten in the spin representation. Some of the simplest applications of the corresponding system of KEs can be found in [16–18].

2. PHOTON SECTOR

In order to construct the photon kinetics it is necessary to derive the corresponding generalization of the quasiparticle formalism developed in Sec. 1, using it as a nonperturbative basis. The interaction with the quantized electromagnetic field $\hat A_\mu(x)$ in the fermion sector of the theory is introduced by means of the substitution $H_0\to H_0+H'$ in the Heisenberg-like equation of motion (3),

$$\dot a(\mathbf{p},t)+U_{(1)}(\mathbf{p},t)a(\mathbf{p},t)+U_{(2)}(\mathbf{p},t)b^+(-\mathbf{p},t)=-i[a(\mathbf{p},t),H_0+H'],\tag{7}$$

where H_0 is the Hamiltonian of the fermion field in the quasiparticle representation and H' is the usual Hamiltonian of interaction with the quantized field. The time dependence of $H'(t)$ makes manifest the nonstationarity of the system that is also reflected in the decomposition of the field operators Ψ, Ψ^+ in the nonstationary basis (2). The same source (external field) induces the nonstationarity of the quantized electromagnetic field. However, this does not alter the mass shell of the photon field, $k^2 = 0$ (in contrast to electron–positron field, $\omega(\mathbf{p}, t)$), so that the standard decomposition is valid,

$$\hat{A}_\mu(x) = (2\pi)^{-3/2} \int \frac{d^3k}{\sqrt{2k}} A_\mu(\mathbf{k}, t)\, \mathrm{e}^{-i\mathbf{kx}}, \tag{8}$$

where $A_\mu(\mathbf{k}, t) = A_\mu^{(+)}(\mathbf{k}, t) + A_\mu^{(-)}(-\mathbf{k}, t)$.

The system of Heisenberg-like equations of motion taking into account the photon subsystem can be written now in an explicit form. For example,

$$\begin{aligned} i\dot{A}_r^{(\pm)}(\mathbf{k}, t) = \mp k A_r^{(\pm)}(\mathbf{k}, t) \mp e(2\pi)^{-3/2} \frac{1}{\sqrt{2k}} \int d^3p_1 d^3p_2\, \delta(\mathbf{p}_1 - \mathbf{p}_2 \mp \mathbf{k}) \times \\ \times \big\{ a^+(\mathbf{p}_1, t)[\bar{u}u]^r(\mathbf{p}_1, \mathbf{p}_2, t) a(\mathbf{p}_2, t) + a^+(\mathbf{p}_1, t)[\bar{u}v]^r(\mathbf{p}_1, \mathbf{p}_2, t) b^+(-\mathbf{p}_2, t) + \\ + b(-\mathbf{p}_1, t)[\bar{v}u]^r(\mathbf{p}_1, \mathbf{p}_2, t) a(\mathbf{p}_2, t) + b(-\mathbf{p}_1, t)[\bar{v}v]^r(\mathbf{p}_1, \mathbf{p}_2, t) b^+(-\mathbf{p}_2, t) \big\}. \end{aligned} \tag{9}$$

Here and below the vectors $\mathbf{p}_1, \mathbf{p}_2, \ldots$ denote the canonical momenta of fermions and $\mathbf{k}_1, \mathbf{k}_2, \ldots$ correspond to the momenta of photons; $[\bar{\xi}\eta]^r_{\beta\alpha}(\mathbf{p}_1, \mathbf{p}_2; t) = \bar{\xi}_\alpha(\mathbf{p}_1, t)\gamma^\mu \eta_\beta(\mathbf{p}_2, t) e^r_\mu$, and $\mathbf{e}^1$, $\mathbf{e}^2$ are the polarization unit vectors, $\mathbf{e}^3 = \mathbf{k}/k$ and $e^0_\mu = \delta^0_\mu$. The photon correlation function is defined as

$$F_{rr'}(\mathbf{k}, \mathbf{k}', t) = \langle A_r^+(\mathbf{k}, t) A_{r'}^-(\mathbf{k}', t) \rangle. \tag{10}$$

With the help of Eq. (9) we can then obtain the first equation of the BBGKY hierarchy

$$\begin{aligned} \dot{F}_{rr'}(\mathbf{k}, \mathbf{k}', t) = ie(2\pi)^{-3/2} \sum_{\alpha\beta} \int d^3p_1 d^3p_2 \Big\{ -\frac{1}{\sqrt{2k}} \delta(\mathbf{p}_1 - \mathbf{p}_2 - \mathbf{k}) \times \\ \times \big[[\bar{u}u]^r_{\beta\alpha}(\mathbf{p}_1, \mathbf{p}_2, t) \langle a^+_\alpha(\mathbf{p}_1, t) a_\beta(\mathbf{p}_2, t) A_r^{(-)}(\mathbf{k}, t) \rangle + \\ + [\bar{u}v]^r_{\beta\alpha}(\mathbf{p}_1, \mathbf{p}_2, t) \langle a^+_\alpha(\mathbf{p}_1, t) b^+_\beta(-\mathbf{p}_2, t) A_r^{(-)}(\mathbf{k}, t) \rangle + \\ + [\bar{v}u]^r_{\beta\alpha}(\mathbf{p}_1, \mathbf{p}_2, t) \langle b_\alpha(-\mathbf{p}_1, t) a_\beta(\mathbf{p}_2, t) A_r^{(-)}(\mathbf{k}, t) \rangle + \\ + [\bar{v}v]^r_{\beta\alpha}(\mathbf{p}_1, \mathbf{p}_2, t) \langle b_\alpha(-\mathbf{p}_1, t) b^+_\beta(-\mathbf{p}_2, t) A_r^{(-)}(\mathbf{k}, t) \rangle \big] + \\ + \frac{1}{\sqrt{2k}} \delta(\mathbf{p}_1 - \mathbf{p}_2 + \mathbf{k}') \big[[\bar{u}u]^{r'}_{\beta\alpha}(\mathbf{p}_1, \mathbf{p}_2, t) \langle a^+_\alpha(\mathbf{p}_1, t) a_\beta(\mathbf{p}_2, t) A_{r'}^{(+)}(\mathbf{k}', t) \rangle + \end{aligned}$$

$$
\begin{aligned}
&+ [\bar{u}v]^{r'}_{\beta\alpha}(\mathbf{p}_1, \mathbf{p}_2, t)\langle a^+_\alpha(\mathbf{p}_1, t) b^+_\beta(-\mathbf{p}_2, t) A^{(+)}_{r'}(\mathbf{k}', t)\rangle + \\
&+ [\bar{v}u]^{r'}_{\beta\alpha}(\mathbf{p}_1, \mathbf{p}_2, t)\langle b_\alpha(-\mathbf{p}_1, t) a_\beta(\mathbf{p}_2, t) A^{(+)}_{r'}(\mathbf{k}', t)\rangle + \\
&\quad + [\bar{v}v]^{r'}_{\beta\alpha}(\mathbf{p}_1, \mathbf{p}_2, t)\langle b_\alpha(-\mathbf{p}_1, t) b^+_\beta(-\mathbf{p}_2, t) A^{(+)}_{r'}(\mathbf{k}', t)\rangle\Big] \Bigg\}.
\end{aligned} \tag{11}
$$

3. TRUNCATION PROCEDURE

The simplest decoupling of the hierarchy,

$$\langle a^+_\alpha(\mathbf{p}_1, t) a_\beta(\mathbf{p}_2, t) A^{(\pm)}_r(\mathbf{k}, t)\rangle \simeq \langle a^+_\alpha(\mathbf{p}_1, t) a_\beta(\mathbf{p}_2, t)\rangle\langle A^{(\pm)}_r(\mathbf{k}, t)\rangle = 0, \tag{12}$$

is not effective. Therefore, we will consider the equation at the second order for the annihilation process (the next-to-last line in Eq. (11)) and the reverse one (the fourth line). We will not write the equations of the second level for these correlators completely in view of their awkwardness and will discuss only the simplest truncation scheme resulting in correlators of the type:

$$
\begin{aligned}
&\langle a^+(\mathbf{p}_1, t) a(\mathbf{p}_2, t) A^{(\pm)}_r(\mathbf{k}, t) A^{(\pm)}_{r'}(\mathbf{k}', t)\rangle \simeq \\
&\qquad \simeq \langle a^+(\mathbf{p}_1, t) a(\mathbf{p}_2, t)\rangle\langle A^{(\pm)}_r(\mathbf{k}, t) A^{(\pm)}_{r'}(\mathbf{k}', t)\rangle,
\end{aligned} \tag{13}
$$

which occur in the RPA approximation as well. We ignore other processes as, e.g., the vacuum polarization effects contained in correlators of the type $\langle a(\mathbf{p}_1, t) b(-\mathbf{p}_2, t) A^{(\pm)}_r(\mathbf{k}, t) A^{(\pm)}_{r'}(\mathbf{k}', t)\rangle$ leading to the polarization functions (5). Ignoring spin effects (see Sec. 2), the approximation (13) in combination with the diagonalization of the photon and fermion correlation functions

$$
\begin{aligned}
\langle A^{(+)}_r(\mathbf{k}, t) A^{(-)}_{r'}(\mathbf{k}', t)\rangle &= \delta_{rr'}\delta(\mathbf{k} - \mathbf{k}') F_r(\mathbf{k}, t), \\
\langle a^{(+)}_\alpha(\mathbf{p}, t) a_\beta(\mathbf{p}', t)\rangle &= \delta_{\alpha\beta}\delta(\mathbf{p} - \mathbf{p}') f(\mathbf{p}, t)
\end{aligned} \tag{14}
$$

leads to the following photon KE for zero initial condition:

$$
\begin{aligned}
\dot{F}(\mathbf{k}, t) = &-\frac{e^2}{2(2\pi)^3 k}\int d^3p \int_{t_0}^{t} dt' K(\mathbf{p}, \mathbf{p} - \mathbf{k}; t, t')\big[1 + F(\mathbf{k}, t')\big] \times \\
&\times \big[f(\mathbf{p}, t') + f(\mathbf{p} - \mathbf{k}, t') - 1\big] \cos\left\{\int_{t'}^{t} d\tau\big[\omega(\mathbf{p}, \tau) + \omega(\mathbf{p} - \mathbf{k}, \tau) - k\big]\right\},
\end{aligned} \tag{15}
$$

where the nonlocal kernel is $K(\mathbf{p}, \mathbf{p}'; t, t') = [\bar{v}u]^r_{\beta\alpha}(\mathbf{p}, \mathbf{p}'; t)[\bar{u}v]^r_{\alpha\beta}(\mathbf{p}', \mathbf{p}; t')$. The next step is based on the Markovian approximation that allows one to ignore the memory effect in the photon distribution, i.e., in the r.h.s. of Eq. (15) we replace $F(\mathbf{k},t') \to F(\mathbf{k},t)$. That brings us to the following quadrature formula

$$F(\mathbf{k},t) = \exp\left[\Phi(\mathbf{k},t)\right] - 1 \simeq \Phi(\mathbf{k},t), \tag{16}$$

where the last approximation is valid for the subcritical fields and

$$\Phi(\mathbf{k},t) = -\frac{e^2}{2(2\pi)^3 k} \int\limits_{t_0}^{t} dt' \int\limits_{t_0}^{t'} dt'' \int d^3p K(\mathbf{p},\mathbf{p}-\mathbf{k};t',t'')\times$$
$$\times \left[f(\mathbf{p},t'') + f(\mathbf{p}-\mathbf{k},t'') - 1\right] \cos\left\{ \int\limits_{t''}^{t'} d\tau \left[\omega(\mathbf{p},\tau) + \omega(\mathbf{p}-\mathbf{k},\tau) - k\right] \right\}.$$

The kernel $K(\mathbf{p},\mathbf{p}';t,t')$ is a slowly varying function of the momentum arguments $\mathbf{p},\mathbf{p}'$ at fixed t and t'. There are also some complicated fast temporal oscillations on this background. For a rough estimate of the effect let us substitute the kernel K by its average value $K(\mathbf{p},\mathbf{p}-\mathbf{k};t,t') \to K_0 = -5$. In addition, we neglect the non-Markovian effect in the fermion distribution functions. Since the main part of fermions is created from vacuum with small momenta, they can be neglected due to the momentum and field ($E \ll E_c$) dependence in the high frequency factor in the r.h.s. of Eq. (17). We obtain the result

$$F(\mathbf{k},t) = \frac{5e^2 n(t)}{2k\delta^2}, \tag{17}$$

where $\delta = 2m - k$ is the frequency mismatch and $n(t) = 2\int d^3p f(\mathbf{p},t)/(2\pi)^3$ is the pair density with the factor 2 from the spin degeneracy. Thus, in the optical region $k \ll m$ the distribution (17) gives $F(k) \sim 1/k$, which is characteristic of the flicker noise (see, e.g., [20]). In the high frequency region $k \sim m$, the Markovian approximation is not justified. Here, a more detailed investigation is necessary.

SUMMARY

Our main result in Eq. (17) defines the frequency dependence of the photon distribution by the factor $1/k$. That this is a multiphoton process can be seen when one identifies the frequency mismatch with the energy Nk of the photon system which is necessary for the energy conservation in the one-photon annihilation process: we obtain $N \sim 2m/k$. For optical lasers this is a huge number and therefore such kind of events is very rare. This conclusion about the role of multi-photon processes is conformed with the analysis of the absorption coefficient of the electron–positron plasma created from the vacuum in the infrared region [19].

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