Inverse Compton effect induced by an intense circularly polarized wave

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(Submitted 15 October 1982)
Zh. Eksp. Teor. Fiz. **84**, 1605–1614 (May 1983)

Nonlinear effects related to the polarization properties of high-energy gamma quanta produced in collisions between relativistic electrons and an intense laser beam are investigated. Exact equations are obtained for the Stokes parameter with allowance for nonlinear effects, as well as the corresponding expansions in terms of the parameter $x^2/(1+x^2)$, where $x^2 = -e^2\overline{A_\mu^2}/m^2$ (A_μ is the electromagnetic potential of the laser wave and m is the electron mass), for the first and second harmonics. The polarization characteristics of the γ quanta produced in the inverse Compton effect are obtained as functions of the polarization states of the colliding particles (the laser photons are circularly polarized in this case). It is shown that circular polarization of the γ quanta produced when two laser-wave photons are absorbed exceeds the circular polarization of the γ quanta in the usual Compton scattering.

PACS numbers: 13.10. + q

INTRODUCTION

The feasibility of producing high-energy colliding γe and $\gamma\gamma$ beams (of reasonable intensity), on the basis of the new-generation linear accelerators (with colliding e^+e^- and e^-e^- beams) was theoretically proved recently.[1–4] Intense photon beams should result from collision of relativistic electrons with a powerful and focused laser beam. The high intensity of the electromagnetic wave can lead to effects that are of noticeable magnitude and are nonlinear in the wave-energy density (see Ref. 5, which contains references to the original works and an extensive bibliography). The characteristic parameters that determine the magnitude of these nonlinear effects are the combinations

$$x^2=-e^2\overline{A_\mu^2}/m^2,\quad \varkappa=xkp/m^2, \tag{1}$$

where A_μ is the 4-potential of the wave electromagnetic field, k and p are the 4-momenta of the colliding γ quantum and electron, m is the electron mass, and the superior bar in (1) denotes averaging over the time. To gain an idea of the scale of the possible nonlinear effects, let us estimate the expected values of the parameters x^2 and $\varkappa$ for the conditions of a realistic future experiment: at a neodymium-laser flash energy ($\omega = 1.17$ eV) equal to 40 J, at a flash duration $\tau = 10$ psec and a focal-spot area 10^{-5} cm^2 we have $x^2 = 0.3$, with $\varkappa = 0.24$ and 1.47 for initial electron energies 50 and 300 GeV, respectively.

The nonlinear effects in the inverse Compton effect lead to a number of qualitatively new effects. Foremost is the appearance of additional peaks in the spectrum of the produced γ quanta, due to absorption of two and more wave photons and corresponding to high energies (compared with the usual Compton effect). In other words the nonlinear effects increase substantially the energies of the scattered γ quanta. This is specially important, for the maximum energy of the γ quanta produced when one photon of the wave is absorbed is decreased by the nonlinear effects (owing to the increase of the effective mass of the electron in the field of the powerful electromagnetic wave).

We study in this article the polarization properties of the high-energy γ quanta produced in inverse Compton scattering of relativistic electrons of an intense circularly polarized electromagnetic wave. We consider collisions of unpolarized as well as polarized electrons. We shall show that it is possible to obtain in such γe collisions high-energy monochromatic beams of γ quanta having a high degree of circular polarization. Such beams will be of fundamental importance for weak-interaction physics, primarily for in the search for W and Z bosons in collisions, $\gamma + e^- \to W^- + \nu$ and $\gamma + e^- \to e^- + Z^0$ (Ref. 6) and for the subsequent study of the properties of these bosons (since the cross sections for the processes $\gamma + e^- \to W^- + \nu$ and $\gamma + e^- \to e^- + Z^0$ can be reliably predicted theoretically).

We shall pay principal attention to the analysis of nonlinear effects in polarization phenomena in the kinematic region in which two and more photons of the wave are absorbed. It must be noted the polarization phenomena in the inverse Compton effect were previously investigated only in the linear approximation[7,8] (one photon is absorbed). Nonlinear effects in polarization phenomena were considered only for collisions of laser photons with immobile electrons, and only particular (simplest) cases were investigated. A systematic analysis of polarization phenomena in the inverse Compton effect with allowance for nonlinear effects is undertaken here for the first time ever.

1. DIFFERENTIAL PROBABILITY OF PHOTON EMISSION

We obtain in this section the general dependence of the polarization characteristics of γ quanta produced in the inverse Compton effect on the polarizations of the laser electromagnetic wave and of the relativistic electrons that collide with this wave. The polarization states of the γ quanta will be characterized by the standard Stokes parameters[9]:

$$\rho_{\mu\nu}={}^1/_2(e_\mu^{(1)}e_\nu^{(1)}+e_\mu^{(2)}e_\nu^{(2)})+{}^1/_2\xi_1(e_\mu^{(1)}e_\nu^{(2)}+e_\mu^{(2)}e_\nu^{(1)})$$
$$-{}^1/_2 i\xi_2(e_\mu^{(1)}e_\nu^{(2)}-e_\mu^{(2)}e_\nu^{(1)})+{}^1/_2\xi_3(e_\mu^{(1)}e_\nu^{(1)}-e_\mu^{(2)}e_\nu^{(2)}),$$

where the relativistic covariant unit vectors $e^{(1)}$ and $e^{(2)}$ are connected with the 4-momenta of the problem in the following manner:

 0038-5646/83/050935-06$04.00

$$e_\mu^{(1)}=\varepsilon_{\mu\nu\rho\sigma}k_\nu p_\rho k_\sigma'/c, \quad e_\mu^{(2)}=[k_\mu(k'p)+k_\mu'(kp)-p_\mu(kk')]/c,$$
$$c^2=(kk')/[2(kp)(k'p)-m^2(kk')],$$

k' is the 4-momentum of the produced γ quantum.

The matrix element of emission of a γ quantum with 4-momentum k' by an electron situated in the field of a powerful circularly polarized wave is defined by the equation

$$\mathfrak{M}=\frac{e}{(2\omega')^{1/2}}\sum_{l=-\infty}^{\infty}(2\pi)^4\delta(k'+q'-k-q)\bar{u}(p')\times\left(B_l\hat{e}+\frac{\hat{r}\hat{k}\hat{e}}{2(kp')}+\frac{\hat{e}\hat{k}\hat{r}}{2(kp)}\right)u(p),$$

$$\begin{gathered}
r_\mu=e(a_{1\mu}B_{1l}+a_{2\mu}B_{2l}), \quad B_l=J_l(z)\exp(i\lambda l\varphi_0),\\
B_{1l}=[(l/z)J_l(z)\cos\varphi_0-i\lambda J_l'(z)\sin\varphi_0]\exp(i\lambda l\varphi_0),\\
B_{2l}=[(l/z)J_l(z)\sin\varphi_0-i\lambda J_l'(z)\cos\varphi_0]\exp(i\lambda l\varphi_0),\\
z^2=\alpha_1^2+\alpha_2^2=\frac{4x^2l^2}{1+x^2}\frac{u}{u_l}\left(1-\frac{u}{u_l}\right),\\
\cos\varphi_0=\frac{\alpha_1}{z}, \quad \sin\varphi_0=\frac{\alpha_2}{z},\\
\alpha_1=e\left[\frac{(a_1p)}{(kp)}-\frac{(a_1p')}{(kp')}\right], \quad \alpha_2=e\left[\frac{(a_2p)}{(kp)}-\frac{(a_2p')}{(kp')}\right],\\
u=\frac{(kk')}{(kp')}, \quad u_l=2l\frac{kp}{(1+x^2)m^2},\\
q_\mu=p_\mu-k_\mu\frac{e^2a^2}{2(kp)}, \quad q_\mu'=p_\mu'-k_\mu\frac{e^2a^2}{2(kp)},
\end{gathered} \tag{2}$$

$J_l(z)$ is a Bessel function.

The helicity $\lambda=\pm1$ corresponding to a wave with right and left polarizations defines a vector potential A_μ:

$$A_\mu=a_{1\mu}\cos kx+\lambda a_{2\mu}\sin kx,$$

where k is the wave 4-vector, $a_{1\mu}$ and $a_{2\mu}$ are the amplitudes of the 4-potential, with $a_1^2=a_2^2=a^2$, $(a_1a_2)=(a_1k)=(a_2k)=0$. The differential probability of emitting a polarized photon by a relativistic polarized electron colliding with a laser beam, per unit volume and per unit time, will have then the following general structure (summation has been carried out over the polarizations of the final electron):

$$\frac{d^2W}{d\varphi du}=\frac{e^2m^2x^2n_e}{16\pi q_0(1+u)^2}\sum_{l=1}^{\infty}[F_0^{(l)}+G_0^{(l)}(s)+\xi_1(F_1^{(l)}+G_1^{(l)}(s))+\xi_2(F_2^{(l)}+G_2^{(l)}(s))+\xi_3(F_3^{(l)}+G_3^{(l)}(s))], \tag{3}$$

where s the 4-vector of the initial-electron polarization and n_e is the density of the initial electrons. The nonzero quantities $F_i^{(l)}$ and $G_i^{(l)}(s)$ are defined as follows:

$$\begin{aligned}
F_0^{(l)}=&\left(2+\frac{u^2}{1+u}\right)\left[\frac{2l^2}{z^2}J_l^2(z)+2J_l'^2(z)\right]\\
&-2\frac{u}{u_l}\left(1-\frac{u}{u_l}\right)\left(2+\frac{x^2}{1+x^2}\right)\frac{4l^2}{z^2}J_l^2(z),\\
F_2^{(l)}=&\lambda\left(2+\frac{u^2}{1+u}\right)\left(1-2\frac{u}{u_l}\right)\frac{4l}{z}J_l(z)J_l'(z),\\
F_3^{(l)}=&16\frac{u}{u_l}\left(1-\frac{u}{u_l}\right)\frac{l^2}{z^2}J_l^2(z)-4\frac{l^2}{z^2}J_l^2(z)+4J_l'^2(z),\\
(1+x^2)G_1^l=&-4\lambda(1+u)\frac{u}{u_l}\frac{\{s,k,p,k'\}}{2m(kp)}\frac{4l}{z}J_l(z)J_l'(z),\\
(1+x^2)G_0^{(l)}&\\
=&8\lambda\frac{u}{u_l}\left[\frac{(sk')}{m}+l\frac{(sk)}{m}\left(1-2\frac{u}{u_l}-\frac{x^2}{1+x^2}\frac{u}{u_l}\frac{1}{1+u}\right)\right]\\
&\times\frac{l}{z}J_l(z)J_l'(z),\\
(1+x^2)G_2^{(l)}=&-2\frac{u}{u_l}\left\{-\frac{(sk')}{m}\left(1-2\frac{u}{u_l}\right)\frac{4l^2}{z^2}J_l^2(z)\right.\\
&-l\frac{(sk)}{m}\left(1+\frac{1}{1+u}\right)\cdot\left[l^2\frac{2}{z^2}J_l^2(z)+2J_l'^2(z)\right]\\
&\left.+l\frac{(sk)}{m}\left[\frac{1}{1+u}+\frac{x^2}{1+x^2}\frac{u}{u_l}\left(2-2\frac{u}{u_l}-\frac{1}{1+u}\right)\right]l^2\frac{4}{z^2}J_l^2(z)\right\},\\
(1+x^2)G_3^{(l)}=&2\lambda\frac{u}{u_l}\left\{(1+u)\frac{(sk')}{m}-l\frac{(sk)}{m}\left[1-2\frac{u}{u_l}\right.\right.\\
&\left.\left.+\frac{x^2}{1+x^2}\frac{u}{u_l}\right]\right\}\frac{4l}{z}J_l(z)J_l'(z),
\end{aligned} \tag{4}$$

where $\{s,k,p,k'\}\equiv\varepsilon_{\mu\nu\rho\sigma}s_\mu k_\nu p_\rho k_\sigma'$, $\varepsilon_{0123}=-1$.

It is important to emphasize that the correlations of the polarizations in the inverse Compton effect with allowance for nonlinear effects have qualitatively the same character as for the usual Compton scattering. In particular, the scattered γ quanta are circularly polarized only if a circularly polarized wave collides with unpolarized electrons, as well when an unpolarized wave collides with polarized (in the scattering plane) relativistic electrons. No new polarization correlations are produced by the nonlinear effects.

Nonetheless, the differential probability (3) is now an infinite sum of terms, each with a corresponding 4-momentum conservation law:

$$lk+q=q'+k'. \tag{5}$$

From this we can obtain the following equation for the energy ω' of the produced γ quanta (valid for head-on collisions accurate to terms of order ω/m and m/ε):

$$\omega'/\varepsilon\equiv y=l\varkappa_1/[1+l\varkappa_1+x^2+(\theta\varepsilon/m)^2], \tag{6}$$

where ε is the energy of the initial electrons, $\varkappa_1=2(kp)/m$, and θ is the emission angle of the final γ quanta relative to the momentum $\mathbf{p}$ of the initial electrons. It can be seen that the maximum energy of the γ quanta produced by absorbing only one photon from a laser electromagnetic wave ($l=1$, i.e., first harmonic) is less than the corresponding energy calculated for the usual Compton effect (at $x^2=0$), but at $l>1+x^2$ this energy exceeds the maximum energy of the γ quanta produced in the usual Compton effect. It follows from the conservation law (5) that the ranges of the variable u (in terms of l) for each term in (3) are given by

$$0\leqslant u\leqslant u_l=2l(kp)/m^2(1+x^2). \tag{7}$$

2. TOTAL AND DIFFERENTIAL CROSS SECTIONS FOR SMALL VALUES OF $x^2/(1+x^2)$

At $x^2 \ll 1$ the main contribution to the probability (3) is made only by the first few harmonics. It will be made clear below that, at low values of x^2, a convenient expansion parameter is the ratio $x^2/(1+x^2)$ rather than x^2. The differential probabilities corresponding to the first two harmonics are then given by (accurate to terms proportional to $x^4/(1+x^2)$):

$$\frac{d^2W^{(i)}}{d\varphi du}=\frac{e^2m^2x^2n_e}{16\pi q_0(1+u)^2}[F_0^{(i)}+G_0^{(i)}+\xi_1G_1^{(i)}+\xi_2(F_2^{(i)}+G_2^{(i)})+\xi_3(F_3^{(i)}+G_3^{(i)})],\quad i=1,2,$$

$$F_0^{(1)}=2+\frac{u^2}{1+u}-4\frac{u}{u_1}\left(1-\frac{u}{u_1}\right)-\frac{4x^2}{1+x^2}\frac{u}{u_1}\left(1-\frac{u}{u_1}\right)\left[1+\frac{u^2}{1+u}-\frac{u}{u_1}\left(1-\frac{u}{u_1}\right)\right],$$

$$F_2^{(1)}=\lambda\left(2+\frac{u^2}{1+u}\right)\left(1-2\frac{u}{u_1}\right)\left[1-\frac{2x^2}{1+x^2}\frac{u}{u_1}\left(1-\frac{u}{u_1}\right)\right],$$

$$F_3^{(1)}=4\frac{u}{u_1}\left(1-\frac{u}{u_1}\right)-\frac{2x^2}{1+x^2}\frac{u}{u_1}\left(1-\frac{u}{u_1}\right)\left(1+2\frac{u}{u_1}-2\frac{u^2}{u_1^2}\right)\quad(1+x^2)G_0^{(1)}=2\lambda\frac{u}{u_1}\left[\frac{(sk')}{m}+\frac{(sk)}{m}\left(1-2\frac{u}{u_1}\right)\right]$$
$$-2\lambda\frac{x^2}{1+x^2}\frac{u^2}{u_1^2}\left[\frac{(sk')}{m}\left(1-\frac{u}{u_1}\right)+\frac{(sk)}{m}\left(\frac{1}{1+u}+2\left(1-\frac{u}{u_1}\right)\left(1-2\frac{u}{u_1}\right)\right)\right]$$

$$(1+x^2)G_1^{(1)}=-4\lambda(1+u)\frac{u}{u_1}\frac{\{s,k,p,k'\}}{2m(\kappa p)}\left[1-\frac{2x^2}{1+x^2}\left(1-\frac{u}{u_1}\right)\frac{u}{u_1}\right],$$

$$(1+x^2)G_2^{(1)}=2\frac{u}{u_1}\left[\frac{(sk)}{m}+\frac{(sk')}{m}\left(1-2\frac{u}{u_1}\right)\right]-\frac{2x^2}{1+x^2}\frac{u^2}{u_1^2}\left[\frac{(sk')}{m}\left(1-\frac{u}{u_1}\right)\left(1-2\frac{u}{u_1}\right)+\frac{(sk)}{m}\left(4-4\frac{u}{u_1}-\frac{u}{u_1}\frac{1}{1+u}\right)\right],$$

$$(1+x^2)G_3^{(1)}=-2\lambda\frac{u}{u_1}\left[(sk)\left(1-2\frac{u}{u_1}\right)-(sk')(1+u)\right]-2\lambda\frac{x^2}{1+x^2}\frac{u^2}{u_1^2}$$
$$\times\left[\frac{(sk)}{m}\left(-1+6\frac{u}{u_1}-4\frac{u^2}{u_1^2}\right)+2\frac{(sk')}{m}(1+u)\left(1-\frac{u}{u_1}\right)\right]\tag{8}$$

for the first harmonic and

$$F_0^{(2)}=\frac{4x^2}{1+x^2}\frac{u}{u_2}\left(1-\frac{u}{u_2}\right)\left[2+\frac{u^2}{1+u}-4\frac{u}{u_2}\left(1-\frac{u}{u_2}\right)\right],$$

$$F_2^{(2)}=4\lambda\frac{x^2}{1+x^2}\frac{u}{u_2}\left(1-\frac{u}{u_2}\right)\left(1-2\frac{u}{u_2}\right)\left(2+\frac{u^2}{1+u}\right),$$

$$F_3^{(2)}=\frac{16x^2}{1+x^2}\frac{u^2}{u_2^2}\left(1-\frac{u}{u_2}\right)^2,$$

$$(1+x^2)G_0^{(2)}$$
$$=8\lambda\frac{x^2}{1+x^2}\frac{u^2}{u_2^2}\left(1-\frac{u}{u_2}\right)\left[\frac{(sk')}{m}+2\frac{(sk)}{m}\left(1-2\frac{u}{u_2}\right)\right]$$

$$(1+x^2)G_1^{(2)}=-16\lambda\frac{x^2}{1+x^2}\frac{u^2}{u_2^2}\left(1-\frac{u}{u_2}\right)(1+u)\frac{\{s,k,p,k'\}}{2m(kp)},$$

$$(1+x^2)G_2^{(2)}$$
$$=\frac{8x^2}{1+x^2}\frac{u^2}{u_2^2}\left(1-\frac{u}{u_2}\right)\left[\frac{(sk')}{m}\left(1-2\frac{u}{u_2}\right)+2\frac{(sk)}{m}\right],$$

$$(1+x^2)G_3^{(2)}$$
$$=8\lambda\frac{x^2}{1+x^2}\frac{u^2}{u_2^2}\left(1-\frac{u}{u_2}\right)\left[(1+u)\frac{(sk')}{m}\right.$$
$$\left.-2\frac{(sk)}{m}\left(1-2\frac{u}{u_2}\right)\right]\tag{9}$$

for the second harmonic.

If the electrons have in the initial state a longitudinal polarization ($s_\parallel=\lambda_e$, $s_\perp=0$), the equations for the quantities $G_{1-3}^{(i)}$ are much simpler:

$$G_1^{(1)}=G_3^{(1)}=G_1^{(2)}=G_3^{(2)}=0,$$

$$G_0^{(1)}=\lambda\lambda_e\left[\left(u+\frac{u}{1+u}\right)\left(1-2\frac{u}{u_1}\right)\right.$$
$$\left.-\frac{x^2}{1+x^2}\frac{u}{u_1}\left(1-\frac{u}{u_1}\right)\left(1-2\frac{u}{u_1}\right)\left(2u+\frac{u}{1+u}\right)\right]$$

$$G_2^{(1)}=\lambda_e\left[u+\frac{u}{1+u}\left(1-2\frac{u}{u_1}\right)^2\right.$$
$$\left.+\frac{2x^2}{1+x^2}\frac{u}{u_1}\left(1-\frac{u}{u_1}\right)\left(\frac{u}{1+u}-2u\right)\right],\tag{10}$$

$$G_0=4\lambda\lambda_e\frac{x^2}{1+x^2}\frac{u}{u_2}\left(1-\frac{u}{u_2}\right)\left(1-2\frac{u}{u_2}\right)\left(u+\frac{u}{1+u}\right)$$

$$G_2^{(2)}=4\lambda_e\frac{x^2}{1+x^2}\frac{u}{u_2}\left(1-\frac{u}{u_2}\right)\left[u+\frac{u}{1+u}\right.$$
$$\left.-4\frac{u}{u_2}\left(1-\frac{u}{u_2}\right)\frac{u}{1+u}\right].$$

The vanishing of the functions $G_1^{(i)}$ and $G_3^{(i)}$ is due to the obvious circumstance that when longitudinally polarized electrons collide with circularly polarized laser photons no linearly polarized γ quanta can be produced.

We emphasize once more that the equations in this section are derived with the Bessel functions expanded in powers of $x^2/(1+x^2)$ rather than x^2. The equations used in this case for u_1 and u_2 were exact. This makes it possible to separate correctly the kinematic singularities. Thus, the differential probability corresponding to the l-th harmonic (at $l>1$) vanishes just at $u=u_l$.

Putting $x^2=0$ in (8) and dividing by the flux of the colliding particles we obtain an equation for the differential cross section for the collision of polarized electrons with circularly polarized photons

$$\frac{d^2\sigma}{d\varphi\, d\varkappa_2}=\frac{r_0^2}{2\varkappa_1^2}(R_0+\xi_1R_1+\xi_2R_2+\xi_3R_3),$$

$$R_0=-\frac{\varkappa_1}{\varkappa_2}-\frac{\varkappa_2}{\varkappa_1}+4\left(\frac{1}{\varkappa_1}+\frac{1}{\varkappa_2}\right)\left(1+\frac{1}{\varkappa_1}+\frac{1}{\varkappa_2}\right)$$
$$-2\lambda\left(\frac{1}{\varkappa_1}+\frac{1}{\varkappa_2}\right)\left[\frac{(sk)}{m}\left(1+\frac{2}{\varkappa_1}+\frac{2}{\varkappa_2}\right)+\frac{(sk')}{m}\right],$$
$$R_1=-4\lambda\frac{\{s,k,p,k'\}}{m^3\varkappa_2}\left(\frac{1}{\varkappa_1}+\frac{1}{\varkappa_2}\right),$$
$$R_2=-\lambda\left(1+\frac{2}{\varkappa_1}+\frac{2}{\varkappa_2}\right)$$
$$-2\left(\frac{1}{\varkappa_1}+\frac{1}{\varkappa_2}\right)\left[\frac{(sk)}{m}+\frac{(sk')}{m}\left(1+\frac{2}{\varkappa_1}+\frac{2}{\varkappa_2}\right)\right],$$
$$R_3=-4\left(\frac{1}{\varkappa_1}+\frac{1}{\varkappa_2}\right)\left(1+\frac{1}{\varkappa_1}+\frac{1}{\varkappa_2}\right) \tag{11}$$
$$+2\lambda\left(\frac{1}{\varkappa_1}+\frac{1}{\varkappa_2}\right)\left[\frac{(sk')}{m}\frac{\varkappa_1}{\varkappa_2}+\frac{(sk)}{m}\left(1+\frac{2}{\varkappa_1}+\frac{2}{\varkappa_2}\right)\right],$$

where $\varkappa_2=-2(k'p)/m^2$. Transforming in (11) to the electron rest system, we obtain known published equations.[10,11] These equations were obtained earlier in a treatment of polarization phenomena for the inverse Compton effect.[8]

None of the equations derived above contain an explicit dependence on the azimuthal angle φ, defined as the angle between the planes made up by the vectors ($\mathbf{k},\mathbf{p}$) and ($\mathbf{k}',\mathbf{p}$). This is due to the definition used by us for the Stokes parameters (relative to unit vectors that depend on φ). At small values of the angle θ, neglecting quantities of the order of m/ε and ω/m, it is convenient to choose other unit vectors, namely $\mathbf{e}^1=(\cos\varphi, -\sin\varphi, 0)$ and $\mathbf{e}^{(2)}=(\sin\varphi, \cos\varphi, 0)$ (the z axis is directed along the momentum $\mathbf{p}$). We have then in this approximation

$$\xi_1=\tilde{\xi}_1\cos 2\varphi-\tilde{\xi}_3\sin 2\varphi,\quad \xi_2=\tilde{\xi}_2,\quad \xi_3=\tilde{\xi}_1\sin 2\varphi+\tilde{\xi}_3\cos 2\varphi,$$

where $\tilde{\xi}_i$ are the Stokes parameters defined in terms of the unit vectors $\mathbf{e}^{(1)}$ and $\mathbf{e}^{(2)}$.

Separating in this manner the dependence on the angle φ we obtain after integrating with respect to the variables u and φ the following equations for the total probabilities corresponding to the first and second harmonics:

$$W_1=\frac{e^2m^2x^2n_e}{16\pi q_0}\left\{\left(1-\frac{4}{u_1}-\frac{8}{u_1^2}\right)\ln(1+u_1)+\frac{1}{2}+\frac{8}{u_1}\right.$$
$$-\frac{1}{2(1+u_1)^2}+\lambda\lambda_e\left[\left(1+\frac{2}{u_1}\right)\ln(1+u_1)-\frac{5}{2}+\frac{1}{1+u_1}\right.$$
$$\left.-\frac{1}{2(1+u_1)^2}\right]$$
$$-\frac{x^2}{1+x^2}\left[2+\frac{44}{3u_1}-\frac{16}{u_1^2}-\frac{16}{u_1^3}-\frac{2}{1+u_1}\right.$$
$$-\frac{8}{u_1}\left(1+\frac{1}{u_1}-\frac{3}{u_1^2}-\frac{2}{u_1^3}\right)\ln(1+u_1)$$
$$-\lambda\lambda_e\left(\frac{2}{3}-\frac{7}{u_1}-\frac{4}{u_1^2}+\frac{3}{1+u_1}\right. \tag{12}$$
$$\left.\left.\left.-\frac{1}{2(1+u_1)^2}+\left(\frac{9}{u_1^2}+\frac{4}{u_1^3}\right)\ln(1+u_1)\right)\right]\right\}$$

$$W_2=\frac{e^2m^2x^2n_e}{16\pi q_0}\frac{x^2}{1+x^2}\left\{2+\frac{8}{3u_2}-\frac{64}{u_2^2}-\frac{64}{u_2^3}-\frac{2}{1+u_2}\right.$$
$$-\frac{4}{u_2}\left(1-\frac{6}{u_2}-\frac{24}{u_2^2}-\frac{16}{u_2^3}\right)\ln(1+u_2)$$
$$+2\lambda\lambda_e\left[\frac{1}{3}+\frac{4}{u_2}-\frac{8}{u_2^2}-\frac{1}{1+u_2}\right.$$
$$\left.\left.+2\left(-\frac{1}{u_2}+\frac{4}{u_3^2}\right)\ln(1+u_2)\right]\right\}.$$

As they should, the total probabilities depend only on the product $\lambda\lambda_e$ of the longitudinal polarizations of the colliding particles. The dependence of this probabilities on λ or λ_e separately would be P-odd and could not occur within the framework of quantum chromodynamics.

3. ENERGY SPECTRUM OF γ QUANTA

For physical applications it is important to know the energy spectrum of the γ quanta produced in the inverse Compton effect. For frontal collision, the invariant variable u is uniquely connected with the energy of the produced γ quantum, namely $u=y/(1-y)$, where $y=\omega'/\varepsilon$. These formulas are valid also for a collision that is not head-on if the particle-encounter angle α satisfies the inequality $\alpha\gg\omega/m\approx 5\times10^{-6}$ rad and if the initial electrons are longitudinally polarized. For transversely polarized electrons account must be taken of corrections proportional to ω/m. We confine ourselves therefore to longitudinally polarized electrons. After integrating with respect to the angle φ we obtain for the energy spectra of the γ quanta

$$\frac{dW^{(i)}}{dy}=e^2m^2\frac{x^2n_e}{8q_0}[F_0^{(i)}+G_0^{(i)}+\lambda(F_2^{(i)}+G_2^{(i)})],$$

$$F_0^{(1)}=1-y+\frac{1}{1-y}-4y_1+4y_1^2$$
$$-(1-y_1)\frac{x^2}{1+x^2}\left[-y+\frac{1}{1-y}-y_1(1-y_1)\right],$$

$$F_2^{(1)}=\lambda\left(1-y+\frac{1}{1-y}\right)(1-2y_1)\left[1-2y_1(1-y_1)\frac{x^2}{1+x^2}\right],$$

$$G_0^{(1)}=\lambda\lambda_e\left[(1-2y_1)\left(1-y+\frac{1}{1-y}\right)\right.$$
$$\left.-y_1(1-y_1)\frac{x^2}{1+x^2}\left(-2+y+\frac{2}{1-y}\right)\right],$$
$$G_2^{(1)}=\lambda_e\left[\frac{1}{1-y}-1+y(1-2y_1)^2\right.$$
$$\left.+4y_1(1-y_1)^2\frac{x^2}{1+x^2}\left(1+yy_1-\frac{1}{1-y}\right)\right],$$
$$y_1\equiv\frac{y}{u_1(1-y)}$$

for the first harmonic and

$$F_0^{(2)}=4y_2(1-y_2)\frac{x^2}{1+x^2}\left[1-y+\frac{1}{1-y}-4y_2(1-y_2)\right]$$
$$F_2^{(2)}=4\lambda y_2(1-y_2)(1-2y_2)\frac{x^2}{1+x^2}\left(1-y+\frac{1}{1-y}\right),$$
$$G_0^{(2)}=4\lambda\lambda_e y_2(1-y_2)(1-2y_2)\frac{x^2}{1+x^2}\left(-1+y+\frac{1}{1-y}\right)$$
$$G_2^{(2)}=4\lambda_e y_2(1-y_2)\frac{x^2}{1+x^2}\left[-1+\frac{1}{1-y}+y(1-2y_2)^2\right],$$
$$y_2\equiv\frac{y}{u_2(1-y)}$$

for the second. It can be seen from this, first, that after averaging over the angle φ the produced γ quantum cannot have linear polarization (as is the case also in usual Compton scattering of circularly polarized γ quanta). Of course, at a fixed value of the angle φ the linear polarization differs from zero.

Since y is uniquely connected with the emission angle of the final γ quantum, the θ distribution of the produced γ quanta can be easily obtained from the equations for the y-distribution. At $\theta=0$ we have the relation $u=u_l$, therefore all the functions $F_i^{(1)}$ and $G_i^{(1)}$ vanish, while the functions $F_i^{(2)}$ and $G_i^{(2)}$ are such that all the characteristics of inverse Compton scattering of laser γ quanta coincide with those of the usual (linear) Compton effect (if the substitution $\varkappa_1\to\varkappa_1/(1+x^2)$ is made. In particular, the degree of circular polarization of the γ quanta emitted at an angle $\theta=0$ is equal to ± 1 independently of the initial-particle energy and independently of λ_e.

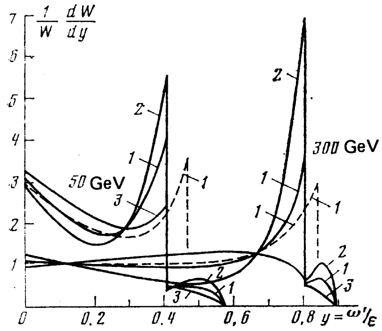

FIG. 1. Spectra of γ quanta produced in collision of relativistic electrons with a laser wave. The dashed curves correspond to the usual Compton effect, curves 1, 2, and 3 correspond to the following polarization states of the electron in the wave: $1\to\lambda_e=0, \lambda=1$; $2\to\lambda_e=1, \lambda=-1$; $3\to\lambda_e=1, \lambda=1$.

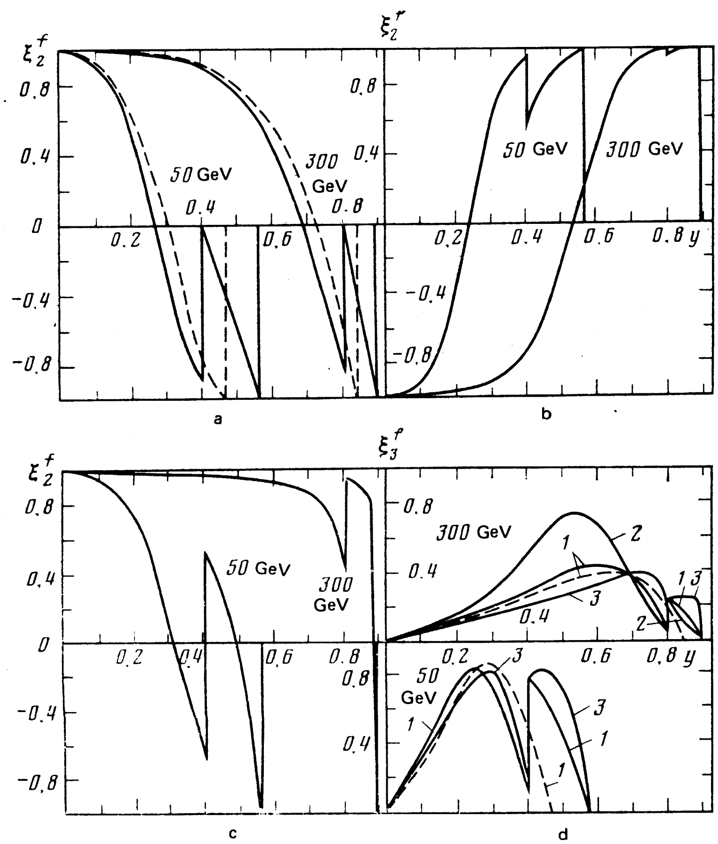

FIG. 2. Energy dependences of the Stokes parameters of the γ quanta produced in the inverse Compton effect, for different colliding-particle polarization states: $1\to\lambda_e=0, \lambda=1$; $2\to\lambda_e=1, \lambda=-1$; $3\to\lambda_e=1$, $\lambda=1$. The dashed curves correspond to the usual Compton effect. a—$\lambda_e=0, \lambda=1$; b—$\lambda_e=1, \lambda=-1$; c—$\lambda_e=1, \lambda=1$.

4. DISCUSSION OF NONLINEAR EFFECTS

In the analysis of the nonlinear effects we start from the following conditions: the incident particles have energies 50 and 300 GeV, the neodymium-laser ($\omega = 1.17$ GeV) power is 40 J, so that $x^2 = 0.3$ if the laser flash duration is 10 psec and the area of the focal spot is 10^{-5} cm^2. As seen from Fig. 1, allowance for the nonlinear effects decreases the maximum energy ω_{max} of the produced γ quanta (in the first peak) compared with usual Compton scattering, but on the other hand a noticeable number of γ quanta due to the second harmonic are produced, and with substantially higher energy. Just as for the linear Compton effect, the intensity of the produced γ quanta depends on the polarization states of the colliding particles: their intensity is larger at $\lambda\lambda_e = -1$ and smaller at $\lambda\lambda_e = +1$ than the intensity of the γ quanta produced in collisions of polarized particles (Fig. 1, curves 2 and 3).

The Stokes parameters of the produced γ quanta are determined by the following equations:

$$\xi_1=0, \quad \xi_2=(F_2+G_2)/(F_0+G_0), \quad \xi_3=(F_3+G_3)/(F_0+G_0),$$
$$F_i=\sum_l F_i^{(l)}, \quad G_i=\sum_l G_i^{(l)}.$$

In the numerical estimates of the parameters ξ_i we confine ourselves to allowance for the contribution of the first two harmonics, calculated with accuracy of the order of $x^2/(1+x^2)$. It can be seen from Fig. 2 that allowance for the nonlinear effects decreases the degree of polarization of the γ quanta in the first peak compared with the usual Compton effect. Thus, at $u = y_{max}$ we have $\xi_2 = -\lambda$ as $x^2 \to 0$, whereas $|\xi_2| \leqslant 1$ at $x^2 \neq 0$ (curves a–c of Fig. 2). We emphasize that the degree of circular polarization of the γ quanta in the second peak is substantially higher than that in the first peak (at $\theta = 0$ we have for the second harmonic $|\xi_2| = \lambda$).

The Stokes parameters of the produced γ quanta depend substantially on the polarization states of the initial electrons. Just as in the case of the usual Compton effect, the most favorable is the situation with $\lambda\lambda_e = -1$. In this case increases take place in both the degree of polarization of the final γ quanta (of all harmonics) and in the angle interval θ (meaning also the energies of the scattered quanta) in which the degree of circular polarization is large.

The degree of circular polarization of the final γ quanta increases with increasing energy of the initial electrons or photons (at a fixed value of the parameter x^2). Thus, for $\varepsilon = 300$ GeV and $\lambda\lambda_e = -1$ we have $|\xi_2| \geqslant 0.9$ at $0 \leqslant \theta \leqslant m/\varepsilon$ for first-harmonic photons and $|\xi_2| \geqslant 0.98$ for photons of the second.

We have already noted that the linear polarization of the γ quanta differs from zero only if the azimuthal angle φ is fixed. After integration with respect to the angle φ the linear polarization vanishes. These statements are valid both for the linear Compton effect and for scattering of an intense laser beam.

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Translated by J. G. Adashko