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Emission of soft photons and electron form factors in the two-dimensional approximation of quantum electrodynamics

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A method is developed for the calculation of the cross section for the emission of soft photons in an arbitrary strong scattering process in a strong magnetic field, when the electrons are on the Landau ground level. The form factors of the electrons are calculated in the two-dimensional approximation of quantum electrodynamics, and the corresponding elastic-scattering cross section is calculated with allowance for the radiative corrections. It is shown that there is no infrared divergence in the summary cross section.

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INTRODUCTION

The discovery of pulsars (neutron stars) which presumably have strong magnetic fields has stimulated interest in the investigation of electrodynamic processes in such fields. An essentially new circumstance that can influence the evolution of similar objects[1] is the suppression of the transverse degrees of freedom of the charged particles (we shall deal here with electrons and positrons). This is possible if the characteristic electron energies ε, which are connected with the field induction B by the relation, are given by

$$\varepsilon^2-m^2<m^2(2B/B_0),\quad B_0=m^2/e=4.41\cdot 10^{13}\text{ G},\tag{1}$$

when the electron is on the Landau ground level. In a number of cases the result is that the quantum electrodynamics degenerates into a practically two-dimensional theory in the sense of the character of the "motion" of the electron (along the field plus a temporal coordinate). This greatly simplifies the calculations, since the electron wave function and the corresponding "two-dimensional" representation of the electron Green's function[2] have an exceedingly simple form in comparison with the usual case (the z axis is directed opposite to the field):

$$\psi(x)=\frac{(\gamma/\pi)^{1/4}}{(2\varepsilon L_2L_3)^{1/2}}\exp\left\{-\frac{\xi^2}{2}+i(p_2x_2+p_3x_3)\right\}u(p),\tag{2a}$$

$$\xi=(x_1\sqrt{\gamma}-p_2/\sqrt{\gamma}),\quad \gamma=|eB|,\quad p=(\varepsilon,\ p_3);$$

$$G(x,y)=\varphi(x,y)G(x-y),\quad \varphi(x,y)=\exp\{{}^1/{}_2 i\gamma(x_1+y_1)(x_2-y_2)\};\tag{3a}$$

$$G(x-y)=-\frac{\gamma}{(2\pi)^3}\frac{1+i\gamma_1\gamma_2}{2}\exp\left\{-\frac{\gamma}{4}[(x_1-y_1)^2+(x_2-y_2)^2]\right\}$$

$$\times\int d^2p\, e^{ip(x-y)}\frac{\hat{p}+m}{p^2-m^2}.\tag{3b}$$

All the scalar products here and below are two-dimensional (0, 3), and the spinor $u(p)$ satisfies the equations

$$(\hat{p}-m)u(p)=0,\quad {}^1/{}_2(1+i\gamma_1\gamma_2)u(p)=u(p),\quad \bar{u}(p)u(p)=2m,\tag{2b}$$

with a density matrix

$$\rho={}^1/{}_2(\hat{p}+m),\tag{2c}$$

and with no summation or averaging over the electron spin states in the square of the matrix element, since the spin projection in the ground state is fixed [Eq. (2b)].

We have shown earlier that in the case of loop diagram the matter reduces to a calculation of the Feynman integrals with respect to the two-dimensional momentum of the loop, and in addition to (1) it is necessary to satisfy the condition $B\gg B_0$.[1,2] On the other hand, if there is no excitation of the vacuum, the problem likewise degenerates to a two-dimensional one, in which the electron is simply nonrelativistic, in accord with (1).

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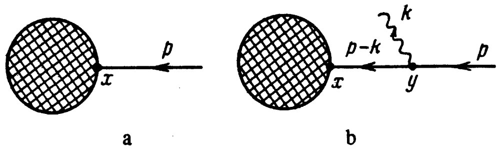

FIG. 1.

In the present paper we consider the emission of soft photons in an arbitrary scattering process with participation of charged particles whose energies satisfy the condition (1). Just as in four-dimensional electrodynamics, the cross section contains an infrared divergence and depends on a cutoff parameter, namely the photon mass λ. We calculate next the electron form factors, which depend in the developed "two-dimensional" formalism on the 2-momentum squared $k^2-k_0^2-k_3^2$. The essential departure from ordinary electrodynamics is that the form factor $g(k^2)$ also contains a part that diverges as $\lambda\to 0$, and depends, like $f(k^2)$, on the cutoff parameter λ. Using scattering by a pointlike immobile potential as an example, we show that the radiative corrections to the elastic-scattering cross sections cancel out the contribution from the soft-photon emission, so that the proposed "two-dimensional" approximation of quantum electrodynamics in strong fields does not contain, in final analysis, an infrared divergence.

1. CROSS SECTION FOR THE EMISSION OF SOFT PHOTONS IN AN ARBITRARY SCATTERING PROCESS

Let us establish the connection between the matrix of diagrams a and b on Fig. 1. Diagram a corresponds to an arbitrary scattering process, where one of the incoming external electron lines is singled out, while diagram b differs from it in that an emitted soft photon of frequency $\omega\ll m$ is added to this line. Separating in the matrix element M_a in explicit form only the factors that pertain to the incoming electron line, we express this element in the form

$$M_a=\int dx_1\,dx_2 f(x)\exp\left\{-\frac{1}{2}\left(x_1\sqrt{\gamma}-\frac{p_2}{\sqrt{\gamma}}\right)^2+ip_2x_2\right\}u(p), \tag{4}$$

where the function $f(x)$ combines all the remaining factors. Using expressions (2) and (3) for the wave function and the Green's function, we find that diagram b of Fig. 1 corresponds to the expression

$$M_b=-\frac{\gamma e\sqrt{4\pi}}{2\pi}\int dx_1\,dx_2 f(x)\frac{\hat{p}-\hat{k}+m}{(p-k)^2-m^2}\hat{e}^*$$
$$\times\int dy_1\,dy_2\exp\left\{-\frac{1}{2}\left(y_1\sqrt{\gamma}-\frac{p_2}{\sqrt{\gamma}}\right)^2+ip_2y_2\right.$$
$$\left.+\frac{i\gamma}{2}(x_1+y_1)(x_2-y_2)-\frac{\gamma}{4}\left[(x_1-y_1)^2+(x_2-y_2)^2\right]\right\}, \tag{5}$$

where $f(x)$ denotes the same block of the diagram. Integrating with respect to dy_1dy_2, we obtain

$$M_b=-e\sqrt{4\pi}\int dx_1\,dx_2 f(x)\frac{\hat{p}-\hat{k}+m}{(p-k)^2-m^2}\hat{e}^*$$
$$\times\exp\left\{-\frac{1}{2}\left(x_1\sqrt{\gamma}-\frac{p_2}{\sqrt{\gamma}}\right)^2+ip_2x_2\right\}u(p). \tag{6}$$

Using also the equality

$$(\hat{p}+m)\hat{e}^*u(p)=2(pe^*)u(p),$$

we obtain, with allowance for the restriction $\omega\ll m$,

$$M_b=M_a e\sqrt{4\pi}\frac{(pe^*)}{(pk)}. \tag{7}$$

Adding the analogous term for the outgoing line and assigning a mass λ to the photon, we obtain for the cross section for scattering with emission of a soft photon the expression

$$d\sigma_b=-d\sigma_0 e^2\left(\frac{p'}{(p'k)}-\frac{p}{(pk)}\right)^2\frac{d^3k}{(2\pi)^2\omega},\quad \omega^2=\mathbf{k}^2+\lambda^2, \tag{8}$$

where we have summed over three polarizations, and $d\sigma_0$ is the cross section of the process without the emission of the soft photon. This formula does not differ in fact in any way from the corresponding expression in the absence of fields.[3] For example, the total cross section for the emission of a photon with frequency $|\mathbf{k}|\leq\omega_m$ in elastic scattering ($p_s'=\pm p_3$) is obtained from the formulas of Sec. 117 of the book[3] by making the substitution $|\mathbf{q}|\to|p_3-p_3'|$, and also contains a cutoff parameter—the photon mass.

We show next that when the radiation corrections to the cross section of the pure elastic scattering are taken into account the dependence on the photon mass can be eliminated. To this end it is necessary to calculate first the correction term of order α for the vertex function in the two-dimensional approximation.

2. CALCULATION OF THE ELECTRON FORM FACTORS IN THE TWO-DIMENSIONAL APPROXIMATION

The form of the correction term of order α to the zeroth-approximation vertex function γ^μ can be determined by comparing the elements of the S matrix for the effective diagram a of Fig. 2:

$$\langle|S|\rangle^\mu=-ie\int\bar{\psi}_-(x)(\Gamma-\gamma)^\mu\psi_+(x)e^{ikx}d^4x \tag{9a}$$

and for the triangular diagram b:

$$\langle|S|\rangle^\mu=i(-ie)^3\int d^4x\,d^4y\,d^4z\,\bar{\psi}_-(y)\gamma^\nu G(y,x)\gamma^\mu G(x,z)\gamma^\lambda\psi_+(z)D(z-y)_{\nu\lambda}e^{ikx}, \tag{9b}$$

which correspond to the decay of a virtual photon into an e^+e^- pair at $t=k^2>0$. Calculating the integrals with

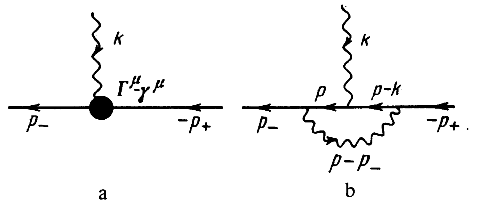

FIG. 2.

respect to the coordinates with allowance for formulas (2) and (3) and comparing these expressions, we get

$$\Gamma^\mu-\gamma^\mu=-\frac{\alpha}{(2\pi)^2}\int dq_\perp{}^2\exp\left\{-\frac{q_\perp{}^2}{2\gamma}\right\}$$
$$\times\int d^2p\,\gamma^\nu\frac{\hat{p}+m}{p^2-m^2}\gamma^\mu\frac{\hat{p}-\hat{k}+m}{(p-k)^2-m^2}\gamma_\nu\frac{i}{(p-p_-)^2-q_\perp{}^2-\lambda^2},\tag{10}$$

where the outer integral is with respect to the square of the transverse momentum of the photon, and the inner one with respect to the two-dimensional momentum of the loop (λ is the mass of the virtual photon). Denoting the last integral by I^μ, we have for its imaginary part (we leave out for the time being the λ^2 in the denominator):

$$\operatorname{Im}I^\mu=-2\pi^2\int d^2p\,\delta^{(+)}(p^2-m^2)\delta^{(+)}((p-k)^2-m^2)2m\frac{2p^\mu-\hat{k}\gamma^\mu}{(p-p_-)^2-q_\perp{}^2},\tag{11}$$

where we have used the obvious relations

$$\gamma^\nu\hat{p}\gamma_\nu=0,\quad \gamma^\nu\gamma_\nu=2.$$

It is easy to show that

$$I_1=\int d^2p\frac{\delta^{(+)}(p^2-m^2)\delta^{(+)}((p-k)^2-m^2)}{(p-p_-)^2-q_\perp{}^2}$$
$$=-\frac{t-4m^2+2q_\perp{}^2}{2q_\perp{}^2[t(t-4m^2)]^{1/2}(q_\perp{}^2+t-4m^2)},\tag{12a}$$
$$\int d^2p\frac{\delta^{(+)}(p^2-m^2)\delta^{(+)}((p-k)^2-m^2)}{(p-p_-)^2-q_\perp{}^2}p^\mu$$
$$=\frac{1}{2}I_1k^\mu+(p_--p_+)^\mu\frac{4m^2-t}{4q_\perp{}^2[t(t-4m^2)]^{1/2}(q_\perp{}^2+t-4m^2)}.\tag{12b}$$

Taking into account also the relations

$$\bar{u}(p_-)(p_--p_+)^\mu u(-p_+)=\bar{u}(p_-)[\sigma^{\mu\nu}k_\nu+2m\gamma^\mu]u(-p_+),$$
$$k^\mu=-\sigma^{\mu\nu}k_\nu+\gamma^\mu\hat{k},\quad \sigma^{\mu\nu}={}^1/_2(\gamma^\mu\gamma^\nu-\gamma^\nu\gamma^\mu),\tag{13}$$

we obtain

$$\operatorname{Im}I^\mu=(2\pi)^2\frac{m^2}{q_\perp{}^2}\left\{\frac{t-4m^2}{[t(t-4m^2)]^{1/2}(t-4m^2+q_\perp{}^2)}\gamma^\mu-\frac{1}{2m}\frac{(-2)}{[t(t-4m^2)]^{1/2}}\sigma^{\mu\nu}k_\nu\right\}\tag{14}$$

which demonstrates conclusively the gauge invariance of the "two-dimensional" approximation to the vertex function

$$\bar{u}(p_-)\Gamma^\mu u(-p_+)k_\mu=0.$$

We denote the factors of γ^μ and $-\sigma^{\mu\nu}k_\nu/2m$ in the curly brackets by $\operatorname{Im}\tilde{f}(t)$ and $\operatorname{Im}\tilde{g}(t)$, respectively. The functions $\tilde{f}(t)$ and $\tilde{g}(t)$ correspond, apart from a common multiplier, to the form factors $f(t)-1$ and $g(t)$ of ordinary electrodynamics, i.e., we should have $\tilde{f}(0)=0$. The value of $\tilde{g}(0)$, however, is not proportional to the anomalous magnetic moment of the electron in a strong magnetic field, since the concept of anomalous moment becomes meaningless at a fixed orientation of the electron spin on the Landau ground level.[4] Thus, the dispersion relations for the quantities $\tilde{f}(t)$ and $\tilde{g}(t)$ should be written in the form

$$\tilde{f}(t)=\frac{t}{\pi}\int_{4m^2}^{\infty}\frac{\operatorname{Im}\tilde{f}(t')}{t'(t'-t)}dt',\tag{15a}$$

$$\tilde{g}(t)=\frac{1}{\pi}\int_{4m^2}^{\infty}\frac{\operatorname{Im}\tilde{g}(t')}{t'-t}dt'.\tag{15b}$$

As usual, the integrals are calculated in the spacelike region $t<0$, introducing new variables

$$t/m^2=-(1-\xi)^2/\xi,\quad t'/m^2=(1+\xi')^2/\xi'\tag{16}$$

The result takes the form

$$\tilde{f}(t,q_\perp'^2)=-\frac{1}{\pi m^2}\left(\frac{1}{2}+\frac{\xi\ln\xi}{1-\xi^2}\right)-\frac{q_\perp'^2}{2\pi m^2(4-q_\perp'^2)}(1-2M(q_\perp'^2))$$
$$-\frac{q_\perp'^2\xi}{\pi m^2[(1+\xi)^2-\xi q_\perp'^2]}\left(\frac{\xi\ln\xi}{1-\xi^2}+M(q_\perp'^2)\right),\tag{17a}$$

where

$$q_\perp'^2=\frac{q_\perp{}^2}{m^2},\quad M(x)=\int_0^1\frac{d\xi'}{1+\xi'(x-2)+\xi'^2},$$
$$\tilde{g}(t,q_\perp'^2)=-\frac{2}{\pi m^2}\frac{\xi\ln\xi}{\xi^2-1}.\tag{17b}$$

We note that the function $\tilde{f}(t,q_\perp'^2)$ has in fact no poles at $q_\perp'^2=4$ and $q_\perp'^2=(1+\xi)^2/\xi$ because the poles are canceled by the zeros of the numerators.

Combining formulas (10), (11), (14), (17a), and (17b) and calculating the integral with respect to $q_\perp^2$ in (10), we obtain the final expression for the vertex function in first order in α. We recognize here that is suffices to retain the square of the photon mass only in the common factor $1/q_\perp^2-1/(q_\perp^2+\lambda^2)$ in expression (14), since only this factor leads to a logarithmic divergence of the integral as $\lambda\to0$. We write down the result in a form similar to the corresponding formula of the "four-dimensional" electrodynamics:

$$\Gamma^\mu-\gamma^\mu=(f(t)-1)\gamma^\mu-\frac{1}{2m}g(t)\sigma^{\mu\nu}k_\nu,\tag{18}$$

$$f(t)-1=\frac{\alpha}{\pi}\left\{\left(\frac{1}{2}+\frac{\xi\ln\xi}{1-\xi^2}\right)\left(-C-2\ln\frac{\lambda}{m}+\ln(2\tilde{B})\right)\right.$$
$$\left.+\int_0^\infty dx\,e^{-x/2\tilde{B}}\left[\frac{1-2M(x)}{2(4-x)}+\frac{\xi}{(1+\xi)^2-\xi x}\left(\frac{\xi\ln\xi}{1-\xi^2}+M(x)\right)\right]\right\},\tag{19}$$

$$g(t)=-\frac{2\alpha}{\pi}\frac{\xi\ln\xi}{1-\xi^2}\left(-C-2\ln\frac{\lambda}{m}+\ln(2\tilde{B})\right).\tag{20}$$

Here $\tilde{B}=B/B_0$, C is the Euler constant, and the integral with respect to $q_\perp'^2=x$ was calculated for the case $B\gg1$. As already noted, this is the necessary condition for the applicability of the representation (3) of the electron Green's function in the case of loop diagrams. It is seen that the logarithmic dependence on the photon mass is present in both form factors, which have also a logarithmic dependence on the field.

In the nonrelativistic approximation $|t|/m^2\ll1$ ($\xi\to1$) the expressions (19) and (20) take the form

$$f(t)-1\approx-\frac{\alpha t}{4\pi m^2}\left(K-\frac{2}{3}\ln\frac{\lambda}{m}\right),$$

$$K=\int_0^4\frac{dx}{x^2}\left\{2+\frac{x}{3}-\frac{4}{[x(4-x)]^{1/2}}2\operatorname{arctg}[x/(4-x)]^{1/2}\right\}+\tag{19a}$$

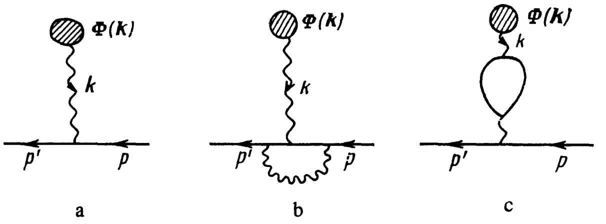

FIG. 3.

$$+\int_0^\infty dx\left\{\frac{1}{x^2}\left[2-\frac{x}{3}-\frac{4}{[x(x+4)]^{1/2}}\ln\left(\frac{x}{2}+1+\frac{1}{2}[x(x+4)]^{1/2}\right)\right]+\frac{1}{3(x+1)}\right\};$$

$$g(t)\approx\frac{\alpha}{\pi}\left(-C-2\ln\frac{\lambda}{m}+\ln(2\tilde{B})\right). \qquad (20a)$$

Thus, the form factor $f(t)-1$ is independent of the field in the nonrelativistic approximation, whereas $g(t)$ increases logarithmically.

3. RADIATIVE CORRECTIONS TO THE ELASTIC-SCATTERING CROSS SECTION IN THE TWO-DIMENSIONAL APPROXIMATION

The scattering of an electron by a potential $\Phi_0(\mathbf{k})$ in the principal and next orders of perturbation theory is represented by the diagrams of Fig. 3, and the sum of the corresponding matrix elements is of the form

$$M_a+M_b+M_c=(-e)\bar{u}(p')\left[\gamma^0+\gamma^0(f(t)-1)-\frac{1}{2m}g(t)\gamma^0\gamma^3k_3+\gamma^0\mathcal{P}(t)\right]u(p)\Phi_0(\mathbf{k}), \qquad (21)$$

$$\Phi_0(\mathbf{k})=\frac{1}{2\pi}\int dk_1\left(\frac{1}{t-k_\perp^2}\right)\Phi_0(\mathbf{k})\exp\left\{-\frac{k_\perp^2}{4\gamma}+\frac{i}{2\gamma}k_1(p_2+p_2')\right\},$$

where $\mathcal{P}(t)$ is the photon polarization operator in the two-dimensional approximation and was calculated by us earlier.[2] The square of the matrix element is calculated, according to (2c), by the usual rules:

$$|\bar{u}(p)Au(p')|^2={}^1\!/_4\mathrm{Sp}\{(\hat{p}'+m)\bar{A}(\hat{p}+m)A\},$$
$$\bar{A}=\gamma^0A^+\gamma^0,$$

and the corresponding interference term in the cross section is written in the form

$$d\sigma_{\rm rad}=d\sigma_{\rm rad}^{(b)}+d\sigma_{\rm rad}^{(c)} \qquad (22a)$$

$$d\sigma_{\rm rad}^{(b)}=2\left[(f(t)-1)+\frac{t}{4\varepsilon^2+t}g(t)\right]d\sigma_0\,, \qquad (22b)$$

$$d\sigma_{\rm rad}^{(c)}=\frac{2}{t-k_\perp^2}\mathcal{P}(t)\,d\sigma_0\,. \qquad (22c)$$

The cross section $d\sigma_{\rm rad}^{(b)}$, which depends on the cutoff parameter, differs from zero, just as $d\sigma_{\rm ph}$, only at $p_3'=-p_3$, and in this case it is readily been that ($v=p_3/\varepsilon$)

$$d\sigma_{\rm rad}^{(b)}=\frac{\alpha}{\pi}\left\{\left[\frac{1}{2}(\ln 2\tilde{B}-C)+\ln\frac{m}{\lambda}\right]\left(2-\frac{1+v^2}{v}\ln\frac{1+v}{1-v}\right)\right.$$
$$\left.+\int_0^\infty dx\,e^{-x/2\tilde{B}}\left[\frac{1-2M(x)}{4-x}+\frac{2(1-v^2)}{4-x(1-v^2)}\left(-\frac{1-v^2}{4v}\ln\frac{1+v}{1-v}+M(x)\right)\right]\right\}d\sigma_0. \qquad (23)$$

Putting $|\mathbf{q}|=2v/\sqrt{1-v^2}$, in the formulas of Sec. 17 of[3], we arrive at the conclusion that the physically meaningful sum of the cross sections

$$d\sigma_{\rm tot}=d\sigma_{\rm ph}+d\sigma_{\rm ph}+d\sigma_0$$

does in fact not contain the photon mass, since

$$-\ln\frac{2\omega_m}{\lambda}+\ln\frac{m}{\lambda}=\ln\frac{m}{2\omega_m}.$$

In the nonrelativistic approximation, (23) takes the form

$$d\sigma_{\rm rad}^{(b)}\approx-\frac{2\alpha}{\pi}v^2\left\{\frac{4}{3}\ln\frac{m}{\lambda}+\ln(2\tilde{B})-C-K\right\}d\sigma_0. \qquad (24)$$

CONCLUSION

It is clear from the foregoing that the contribution of processes with participation of charged leptons to the electromagnetic radiation of neutral stars that have strong magnetic fields can be estimated in simple fashion. It suffices for this purpose to know the cross sections of the corresponding process without photon emission. Foremost among them are, apparently, the scattering of electrons by photons and neutrons, as well as e^-e^- scattering. These cross sections, in turn, can be easily determined within the framework of the "two-dimensional" approximation.

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