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SCATTERING OF ELECTROMAGNETIC WAVES BY NEUTRAL FERMI PARTICLES

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The scattering (Compton effect) of plane electromagnetic waves by uncharged fermions with an anomalous magnetic moment is studied. The scattering cross section, probability, and intensity of the polarized scattered radiation are found. Spin effects are considered, and it is found that the particle spin may take on a predominant orientation as the result of scattering. Effects which depend on the intensity of the wave and which lead to the emission of secondary harmonics are considered.

1. MOTION OF A NEUTRAL FERMION IN A PLANE ELECTROMAGNETIC WAVE

THE motion of an uncharged Dirac particle (for example, a neutron) having a negative anomalous magnetic moment—$c\hbar\mu$ can be described by the generalized Dirac equation:[1]

$$i\hbar\partial\Psi/\partial t = \mathscr{H}\Psi, \quad \mathscr{H} = c(\boldsymbol{\alpha}\mathbf{p}) + \rho_3 mc^2 + c\hbar\mu\{\rho_3(\boldsymbol{\sigma}\mathbf{H}) + \rho_2(\boldsymbol{\sigma}\mathbf{E})\}, \tag{1}$$

where α, ρ, σ are the Dirac matrices, $\mathbf{p} = -i\hbar\nabla$, and $\mathbf{E}$ and $\mathbf{H}$ are the external electric and magnetic fields.

Let us consider the motion of a particle in a plane electromagnetic field propagating along a unit vector $\mathbf{n}$. In this case, the fields $\mathbf{E}$ and $\mathbf{H}$ satisfy the relations*

$$\mathbf{E} = \mathbf{E}(\xi), \quad \mathbf{H}(\xi) = [\mathbf{nE}],$$

$$(\mathbf{nE}) = (\mathbf{nH}) = (\mathbf{EH}) = 0, \quad \xi = ct - z.$$

as is well known.[2] It is easy to see that the integrals of motion in the case of motion in such fields will be the operators.

$$\hbar\hat{\mathbf{k}} = \mathbf{p} - \mathbf{n}(\mathbf{pn}), \quad c\hbar\hat{\lambda} = \mathscr{H} - c(\mathbf{pn}).$$

Consequently, in addition to (1), the wave function Ψ must satisfy the equations

$$\hat{\mathbf{k}}\Psi = \mathbf{k}\Psi, \quad (\mathbf{kn}) = 0, \quad \hat{\lambda}\Psi = \lambda\Psi. \tag{2}$$

With account of (2), one must seek the wave function in the form

$$\Psi = L^{-3/2}\exp\{-ic\lambda t + i(\mathbf{kr}) - ik_3\xi\}\psi(\xi),$$
$$k_3 = \frac{k_0^2 + k^2 - \lambda^2}{2\lambda}, \quad k_0 = \frac{mc}{\hbar}. \tag{3}$$

Substituting (3) in Eq. (1), we get a set of equations for the components of the spinor $\psi(\xi)$. The spinor $\psi(\xi)$ can be sought in a form similar to the form of the wave function of an electron in a plane wave:[3, 4]

$$\psi(\xi) = N\begin{pmatrix} k_0 + \lambda + (\boldsymbol{\sigma}\mathbf{n})(\boldsymbol{\sigma}\mathbf{k}) \\ (\boldsymbol{\sigma}\mathbf{n})(k_0 - \lambda) + (\boldsymbol{\sigma}\mathbf{k}) \end{pmatrix} U. \tag{4}$$

In Eq. (4) and below, $\boldsymbol{\sigma}$ are the Pauli matrices and U is a two-component spinor, and N is a normalizing factor. If U satisfies the condition

$$\frac{1}{L}\int U^+ U d\xi = 1,$$

then, for a choice of N in the form $N = [2(k_0^2 + k^2 + \lambda^2)]^{-1/2}$, the wave functions (3) are normalized to unity.

In contrast with the electronic waves functions, the spinor U is not constant and satisfies the equation

$$dU/d\xi = -\mu(\boldsymbol{\sigma}\mathbf{n})(\boldsymbol{\sigma}\mathbf{E})U. \tag{5}$$

If we assume the quantity μ to be small, then Eq. (5) has the solution

$$U = \{1 + \mu(\boldsymbol{\sigma}\mathbf{n})(\boldsymbol{\sigma}\mathbf{A})\}U_0, \tag{6}$$

in the approximation that is linear in μ, where $\mathbf{A}$ is the vector potential of a plane wave in the Lorentz gauge, $(\mathbf{A}\cdot\mathbf{n}) = 0$, U_0 is an arbitrary constant spinor which characterizes the orientation of the spin of the particle. Actually, if we assume the average value $\langle\sigma\rangle$ of the Dirac matrices according to the functions (3), (4), (6) in the nonrelativistic approximation $\lambda \sim k_0$, $\mathbf{k} = 0$, then we get

*$[\mathbf{mE}] \equiv \mathbf{m}\times\mathbf{E}$, $(\mathbf{mH}) \equiv \mathbf{m}\cdot\mathbf{H}$.

$$\langle\sigma\rangle = U_0^+\sigma U_0.$$

By making the spinor U_0 satisfy the equation

$$(\sigma l) U_0 = \zeta U_0, \quad \zeta = \pm 1, \tag{7}$$

where $\mathbf{l}$ ($\sin\eta\cos\Phi$, $\sin\eta\sin\Phi$, $\cos\eta$) is the unit vector of the direction characterized by the angles η and Φ, we get a particle with spin oriented along the direction of $\mathbf{l}$. The solution of Eq. (7) with account of normalization has the form

$$U_0 = \frac{1}{\sqrt{2}}\begin{pmatrix}\zeta\sqrt{1+\zeta\cos\eta}\, e^{-i\Phi/2}\\ \sqrt{1-\zeta\cos\eta}\, e^{i\Phi/2}\end{pmatrix}$$

2. RADIATION OF AN UNCHARGED DIRAC PARTICLE MOVING IN A PLANE ELECTROMAGNETIC WAVE

The spectral angular distribution of the intensity of the polarized radiation of a neutral Fermi particle having an anomalous moment has, as is well known, the form

$$dW = \frac{c^3\hbar^2\mu^2}{2\pi}\frac{\varkappa^4\sin\theta\, d\theta\, d\varphi}{|1-(\varkappa_0\mathbf{n})+\partial\lambda'/\partial\varkappa|}|l_2A_2 - il_3A_3|^2, \tag{8}$$
$$\lambda' = \lambda - \varkappa(1-\varkappa_0\mathbf{n}),$$

where $\kappa_0 = (\sin\theta\cos\varphi,\ \sin\theta\sin\varphi,\ \cos\theta)$ is the unit wave vector of the photon, the quantities A_2 and A_3 are connected with the matrix elements of $\mathbf{B}$:

$$W = {}^1/_8 W_0 p\alpha^2 (E_0/K_0)^2\{{}^1/_2 l_2^2 + {}^1/_2 l_3^2 + l_2 l_3 \sin 2\psi\}(1+\beta_3).$$
$$\langle\mathbf{B}\rangle = \int \Psi^{+\prime}\exp\{-i\varkappa\mathbf{r} + ic\varkappa t\}\mathbf{B}\Psi d^3x,$$
$$\mathbf{B} = \begin{pmatrix}[\sigma\varkappa_0] & -i\sigma\\ i\sigma & -[\sigma\varkappa_0]\end{pmatrix}. \tag{9}$$

In Eqs. (8) and (9), the primes denote quantities referring to final states, and $\boldsymbol{\beta}_2$ and $\boldsymbol{\beta}_3$ are unit vectors of linear polarization, orthogonal to one another and to the vector $\boldsymbol{\kappa}_0$. A different choice of the quantities l_2 and l_3 corresponds to different polarizations of the radiation: $l_2 = 1$ and $l_3 = 0$, or $l_2 = 0$ and $l_3 = 1$ to two components of linear polarization; $l_2 = l_3 = 1/\sqrt{2}$ to right circular polarization, $l_2 = -l_3 = 1/\sqrt{2}$ left circular polarization; $l_2^2 = l_3^2 = 1$ and $l_2 l_3 = 0$ correspond to unpolarized radiation of light. The appearance in (8) of the factor $1 - (\boldsymbol{\kappa}_0 \cdot \mathbf{n} + \partial\lambda'/\partial\kappa)$ and the "conservation law" $\lambda' = \lambda - \kappa(1 - \boldsymbol{\kappa}_0 \cdot \mathbf{n})$ is connected with the nonstationarity of the problem (see [4]).

Assuming μ to be a small quantity (i.e., using the solution (6)), and choosing the vector $\mathbf{n}$ along z and the vectors $\boldsymbol{\beta}_i$ in the form

$$\beta_2 = \frac{[\varkappa_0\mathbf{n}]}{\sqrt{1-(\varkappa_0\mathbf{n})^2}}, \quad \beta_3 = \frac{\varkappa_0(\varkappa_0\mathbf{n}) - \mathbf{n}}{\sqrt{1-(\varkappa_0\mathbf{n})^2}},$$

it is not difficult to compute the intensity of the radiation in the general case. If the potential $\mathbf{A}(\xi)$ is periodic (with period $T = 2\pi c/\omega$), then Eq. (8) takes the form (in the initial state $\mathbf{k} = 0$)

$$dW = \frac{3W_0}{2}\frac{\alpha^2}{2\pi}\sum_{m=1}^{\infty}\frac{p^2(1+\beta_3\cos\theta')\sin\theta' d\theta' d\varphi}{[1+p(1-\cos\theta')]^4}|l_2S_2 - il_3S_3|^2, \tag{10}$$

$$S_2 = \frac{(\varkappa_0\mathbf{E}_m)\bar{\sigma}_3\cos\theta'}{\sin\theta} + (\bar{\sigma}\mathbf{E}_m)\sin\theta', \quad S_3 = \frac{(\mathbf{n}[\varkappa_0\mathbf{E}_m])\bar{\sigma}_3}{\sin\theta}.$$

Here the following notation has been introduced ($m = 1, 2, 3, \ldots$):

$$p = 2m\alpha\frac{\omega}{\omega_0}, \quad \alpha = \frac{\lambda}{k_0}, \quad \beta_3 = \frac{1-\alpha^2}{1+\alpha^2},$$

$$\mathbf{E}_m = \frac{1}{TQ_0}\int_0^T \mathbf{E}(\xi)\, e^{im\omega\xi/c}d\xi, \quad \bar{\sigma} = U_0^{+\prime}\sigma U_0, \tag{11}$$

$$\omega_0 = 2ck_0 = 2.856\cdot 10^{24}\,\text{sec}^{-1}; \quad Q_0 = k_0/\mu = 1.558\cdot 10^{20}\,\text{Oe}$$
$$W_0 = {}^{16}/_3 c^3\hbar^2\mu^2 k_0^4 = 7.672\cdot 10^{19}\ \text{erg/sec}$$

(the numerical values are shown for the neutron).

The angle θ' is connected with the angle θ by the transformation

$$\cos\theta = \frac{\cos\theta' + \beta_3}{1+\beta_3\cos\theta'}.$$

The frequency of radiation of the photon ω' is connected with the wave frequency by the relation

$$\omega' = c\varkappa = \frac{(1-\beta_3)m\omega}{1-\beta_3\cos\theta + 2m\dfrac{\omega}{\omega_0}\sqrt{1-\beta_3^2}\,(1-\cos\theta)}$$
$$= \frac{1+\beta_3\cos\theta'}{1+\beta_3}\frac{m\omega}{1+p(1-\cos\theta')} \tag{12}$$

and for $m = 1$ corresponds to the usual Compton effect. The probability of emission is obtained by division of Eq. (10) by $\hbar\omega'$, summation over the polarizations of the photon and integration over the angles:

$$w = \frac{3}{4T_0}\frac{\alpha^3}{1+\alpha^2}\sum_{m=1}^{\infty}\frac{p}{2\pi}\int\frac{\sin\theta' d\theta' d\varphi}{[1+p(1-\cos\theta')]^3}(|S_2|^2 + |S_3|^2),$$

$$T_0 = \frac{3}{32c^2\hbar\mu^2k_0^3} = 2.941\cdot 10^{-13}\ \text{sec}. \tag{13}$$

The scattering cross section can be obtained from (13) by dividing the probability by the photon flux density in the incident wave, equal to $\overline{E^2}c/4\pi\hbar\omega$.

It follows from Eq. (11) that only those harmonics are radiated which are present in the incident wave. Therefore, it is natural to consider the isolated monochromatic component E. Since the monochromatic wave is elliptically polarized (see [2]), we then choose the axes x and y along the axes of the polarization ellipse of the incident wave. We shall characterize the polarization by the angle ψ:

$$\mathbf{E} = \sqrt{2}E_0\left(\mathbf{i}\cos\psi\cos\frac{\omega}{c}\xi + \mathbf{j}\sin\psi\sin\frac{\omega}{c}\xi\right), \quad \overline{E^2} = E_0^2.$$

In this case, only the Fourier component $\mathbf{E}_1$ is different from zero, m = 1:

$$\mathbf{E}_1 = \frac{E_0}{\sqrt{2}Q_0}(\mathbf{i}\cos\psi + i\mathbf{j}\sin\psi).$$

Integrating over the angle φ, summing over the final and averaging over the initial spins, we get

$$W = \frac{3}{16}W_0\alpha^2p^2\left(\frac{E_0}{Q_0}\right)^2\int_0^\pi \frac{(1+\beta_3\cos\theta')\sin\theta'\,d\theta'}{[1+p(1-\cos\theta')]^4} \times \{l_2^2[\cos^2\theta' + 2\sin^2\theta'] + l_3^2 + 2l_2l_3\sin 2\psi\cos\theta'\}. \quad (14)$$

Integration over θ' can be carried out exactly; however, the expressions obtained are very cumbersome and unclear. In the limiting case of small p, we have

$$W = W_0\alpha^2p^2\left(\frac{E_0}{Q_0}\right)^2\left\{\frac{5}{8}l_2^2 + \frac{3}{8}l_3^2 + \frac{1}{4}l_2l_3\beta_3\sin 2\psi\right\}.$$

In the case of large p, we obtain

$$W = {}^1/_8 W_0 p\alpha^2(E_0/Q_0)^2\{{}^1/_2\, l_2^2 + {}^1/_2\, l_3^2 + l_2l_3\sin 2\psi\}(1+\beta_3).$$

It is then seen that in the case of small p the radiation is appreciably linearly polarized. With increase in p the degree of linear polarization decreases. There is circular polarization of the radiation at any p. Linear polarization of the radiation does not depend on the polarization of the incident wave, while the circular polarization depends on it in significant fashion.

We shall divide the probability of radiation into parts corresponding to transitions with spin reversal and without spin reversal:

$$w = \frac{3}{16T_0}\frac{\alpha^3}{1+\alpha^2}p\left(\frac{E_0}{Q_0}\right)^2\int_0^\pi \frac{\sin\theta'\,d\theta'}{[1+p(1-\cos\theta')]^3} \times \left(\frac{1+\zeta\zeta'}{2}R_1 + \frac{1-\zeta\zeta'}{2}R_2\right),$$

$$R_1 = (1+\cos^2\theta')\cos^2\eta + (1+\cos 2\Phi\cos 2\psi)\sin^2\eta\sin^2\theta',$$

$$R_2 = (1+\cos^2\theta')\sin^2\eta + 2\sin^2\theta' \times \left[1 - \frac{1}{2}\sin^2\eta(1+\cos 2\Phi\cos 2\psi) - \zeta\cos\eta\sin 2\psi\right]. \quad (15)$$

It is then seen that the probability of transitions with spin reversal depends on the initial spin orientation. Thus, as a result of scattering, the spin obtains a preferential orientation. As is well known, an analogous effect occurs for electrons moving in a magnetic field.[6] In our case, this effect is most clearly illustrated in the case in which the incident wave is circularly polarized ($\psi = g\pi/4$, g = 1 for the right circular polarization and g = −1 for the left) and the spin of the electron is oriented along the direction of motion of the wave ($\eta = 0$). In this case we have

$$R_1 = 1 + \cos^2\theta', \quad R_2 = 2(1-\zeta g)\sin^2\theta'. \quad (16)$$

It is then seen that, as the result of scattering, the spin of the particle is oriented along the direction of the spin of the incident photon. Integration over the angles in (16) leads to the expression

$$w = \frac{3}{16T_0}\frac{\alpha^3}{1+\alpha^2}\left(\frac{E_0}{Q_0}\right)^2\left\{\frac{1+\zeta\zeta'}{2} \times \left[\frac{2(1+p)(2p^2-2p-1)}{p(1+2p)^2} + \frac{\ln(1+2p)}{p^2}\right] + \frac{1-\zeta\zeta'}{2}\frac{8(1-\zeta g)p(1+p)}{3(1+2p)^3}\right\}.$$

By assuming that there were n_0 particles in the initial state, of which n_1 had $\zeta = 1$, and n_{-1} corresponded to $\zeta = -1$ ($n_1 + n_{-1} = n_0$), we obtain the result that in the passage of time the numbers n_i change in the following fashion (see [6]):

$$n_1(t) = \frac{n_0}{2}\left[1 - g + \left(\frac{n_1 - n_{-1}}{n_0} + g\right)e^{-t/\tau}\right],$$

$$n_2(t) = \frac{n_0}{2}\left[1 + g - \left(\frac{n_1 - n_{-1}}{n_0} + g\right)e^{-t/\tau}\right],$$

$$\tau = T_0\frac{1+\alpha^2}{\alpha^3}\left(\frac{Q_0}{E_0}\right)^2\frac{(1+2p)^3}{p(1+p)}$$

It follows from the last formula that the effective time for spin reversal τ has a minimum as a function of p for $p \sim 1$:

$$\tau_{\min} = 6\sqrt{3}T_0\frac{1+\alpha^2}{\alpha^3}\left(\frac{Q_0}{E_0}\right)^2$$

i.e., the spin has a point of maximum instability at some frequency ω of the incident wave.

The total transition probability, averaged over the spin, does not depend on the polarization of the incident wave and has the form

$$w = \frac{3}{16T_0}\frac{\alpha^3}{1+\alpha^2}\left(\frac{E_0}{Q_0}\right)^2 \times \left[\frac{2(1+p)(12p^3+18p^2+12p+3)}{3p(1+2p)^3} - \frac{\ln(1+2p)}{p^2}\right]$$

In the limiting cases, we get

$$w = \frac{p}{T_0}\frac{\alpha^3}{1+\alpha^2}\left(\frac{E_0}{Q_0}\right)^2, \quad (p \sim 0),$$

$$w = \frac{3}{16T_0}\frac{\alpha^3}{1+\alpha^2}\left(\frac{E_0}{Q_0}\right)^2, \quad (p \sim \infty).$$

It then follows that for $p \gg 1$ the probability of emission ceases to depend on the frequency of the incident wave and is determined only by the dynamic characteristics of the particles and by the wave amplitude.

After determining the scattering cross section, we find that it does not depend on the wave amplitude and is determined by the expression

$$\sigma = \frac{3r_0^2}{16}\frac{\alpha^2}{1+\alpha^2}\left[\frac{2(1+p)(12p^3+18p^2+12p+3)}{3(1+2p)^3} - \frac{\ln(1+2p)}{p}\right]$$

$$r_0 = 8\sqrt{2\pi/3c\hbar k_4}\mu^2 = 1{,}626\cdot 10^{-15}\text{ cm}.$$

In the limiting cases, we have

$$\sigma = r_0^2\frac{p^2\alpha^2}{1+\alpha^2}, \quad (p\sim 0),$$

$$\sigma = \frac{3r_0^2}{16}\frac{p\alpha^2}{1+\alpha^2}, \quad (\infty \sim d) \tag{17}$$

3. EFFECTS WHICH DEPEND ON THE WAVE INTENSITY IN THE SCATTERING OF ELECTROMAGNETIC WAVES BY NEUTRAL PARTICLES

Equation (5) has an exact solution for certain special cases of the incident wave. Let us consider the wave polarized linearly along the x axis,

$$A_x = A(\xi), \quad A_y = A_z = 0.$$

In this case, there is an exact solution of (5) which undergoes transition to the solution of (6)

$$U = \frac{1}{2}\{(1+\sigma_2)e^{i\mu A} + (1-\sigma_2)e^{-i\mu A}\}U_0.$$

as $\mu \to 0$.

Assuming the incident wave to be monochromatic

$$A = -\frac{cE_0}{\omega}\sin\frac{\omega}{c}\xi,$$

we get for the intensity of radiation averaged over the spins and integrated over φ:

$$W = \frac{3W_0}{32}\sum_{m=1} p^4 J_m^2(q)\int_0 \frac{(1+\beta_3\cos\theta')\sin\theta'\,d\theta'}{[1+p(1-\cos\theta')]^4}$$

$$\times\ \{l_2^2[\cos^2\theta' + 2\sin^2\theta'] + l_3^2\},$$

where $J_m(q)$ is a Bessel function, $q = \omega_0 E_0/\omega Q_0$. As $q \to 0$, which corresponds to $\mu \to 0$, we get Eq. (14) if we set $\psi = 0$ there. We then see that the exact solution leads to the possibility of radiation of higher harmonics. A similar situation occurs in the ordinary electron Compton effect (see [4]). Equation (12) remains valid.

Equation (5) also admits an exact solution in the case of a monochromatic wave that is circularly polarized:

$$\mathbf{E} = E_0\left(\mathbf{i}\cos\frac{\omega}{c}\xi + g\mathbf{j}\sin\frac{\omega}{c}\xi\right).$$

The solution of (5), which goes over into (6) as $\mu \to 0$, has the form

$$U = \left(\frac{1+\sqrt{1+q^2}}{8\sqrt{1+q^2}}\right)^{1/2}\Big\{(1+\sigma_3)e^{-i\delta_1\xi} + (1-\sigma_3)e^{i\delta_1\xi} - \frac{igq(\sigma_1+i\sigma_2)}{1+\sqrt{1+q^2}}e^{-i\delta_2\xi} - \frac{igq(\sigma_1-i\sigma_2)}{1+\sqrt{1+q^2}}e^{+i\delta_2\xi}\Big\}U_{0_1}$$

$$\delta_1 = \frac{g\omega}{2c}(1-\sqrt{1+q^2}), \quad \delta_2 = \frac{g\omega}{2c}(1+\sqrt{1+q^2}).$$

In this case, the intensity averaged over the spins and integrated over φ has the form

$$W = \frac{3W_0}{64}\frac{q^2}{1+q^2}\sum_{m=1}^{4} p_m^4\int_0^\pi \frac{(1+\beta_3\cos\theta')\sin\theta'\,d\theta'}{[1+p_m(1-\cos\theta')]^4}S_{m_1}$$

$$p_m = 2\alpha\frac{\omega_m}{\omega_0}. \tag{18}$$

Here, we use the notation

$$S_1 = (l_2\cos\theta' + gl_3)^2, \quad \omega_1 = \omega;$$

$$S_2 = 2l_2^2\sin^2\theta', \quad \omega_2 = \omega\sqrt{1+q^2};$$

$$S_3 = \frac{q^2(l_2\cos\theta' + gl_3)^2}{2(1+\sqrt{1+q^2})^2}, \quad \omega_3 = \omega(1+\sqrt{1+q^2});$$

$$S_4 = \frac{(1+\sqrt{1+q^2})^2}{2q^2}(l_2\cos\theta' - gl_3)^2, \ \omega_4 = \omega(\sqrt{1+q^2}-1).$$

As $q \to 0$, we again get the result (14) ($\psi = g\pi/4$). The frequency of radiation of the photon is as before determined by Eq. (12), where in the place of $m\omega$ we must write ω_m in this case.

Thus, in the case of circular polarization of the incident wave, four different harmonics are radiated; these are not multiples of the fundamental frequency. For example, at low q, we have $\omega_1 \sim \omega_2 \sim \omega$, $\omega_3 \sim 2\omega$, $\omega_4 \sim q^2\omega/2$. This feature has no analog with the ordinary electron Compton effect.

It is easy to obtain the intensity of the polarized radiation of a neutron from Eqs. (18) and (19) as $\omega \to 0$ for constant fields $\mathbf{E}$ and $\mathbf{H}$ that are equal in magnitude and orthogonal. Assuming μ to be small, we find

$$W = \frac{3}{4}\alpha^4 W_0\left(\frac{E_0}{Q_0}\right)^4\int_0^\pi (1+\beta_3\cos\theta')$$

$$\times\ \sin\theta'\,d\theta'[l_2^2(2-\cos^2\theta') + l_3^2].$$

Integrating over θ', we get

$$W = 4\alpha^4 W_0\left(\frac{E_0}{Q_0}\right)^4\left\{\frac{5}{8}l_2^2 + \frac{3}{8}l_3^2\right\}.$$

Similar results were obtained in [5,7] for a purely magnetic field.

We now draw some brief conclusions. The presence of an anomalous magnetic moment in neutral Dirac particles leads to the possibility of Compton scattering of electromagnetic waves by these particles. The scattered radiation possesses significant linear polarization that is independent of the polarization of the incident wave and circular polarization that is significantly connected with the polarization of the wave. The scattering cross section does not depend on the amplitude of the incident wave and its polarization, but does depend on its frequency and the dynamic characteristics of the particle. At high frequencies ($\omega \sim 10^{23}$ sec^{-1}), the scattering cross section is only a few orders of magnitude smaller than the ordinary electron Compton cross section and evidently is capable of experimental measurement (see Eq. (17)).

As a result of the scattering, there is a preferential orientation of the neutron spin; however, even the minimal effective orientation time for reasonable amplitudes of the field is very large and is not yet experimentally measurable. The existence of a minimal spin orientation time demonstrates the presence of a region of maximum spin instability.

Expansion in the moment μ is equivalent to expansion in the quantity $q = \omega_0 E_0/\omega Q_0$, i.e., it is possible either for low intensities of the incident wave or for its high frequencies. In the optical range, such an expansion is possible for all reasonable E_0. Effects which depend on the intensity of the wave appear in the emission of higher harmonics and of harmonics which are not multiples of the fundamental frequency, which is a very interesting fact. For very low frequencies, an expansion in μ is impossible and one must use the exact solution.

It must be noted that consideration has been given as to the dependence of the momentum of the neutron on the frequency of the incident wave (see [8], p. 131). The results obtained can obviously serve as the basis of a test of this hypothesis.

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