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Effect of a strong electromagnetic wave on the radiation emitted by weakly excited electrons moving in a magnetic field

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A quantum-mechanical treatment is given of the spontaneous emission of radiation by electrons moving in a constant uniform magnetic field and the field of a circularly polarized plane wave propagating in the direction of the magnetic field. Electrons occupying low-lying levels $n = 0, 1$ are considered. An analysis of the emission probability is presented. The behavior of spin during the emission process is considered and it is shown that the $n = 1$ state with the electron spin lying along the magnetic field is stable against a transition to the $n = 0$ state.

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There has been undoubted recent interest in the interaction between a strong radiation field and electrons moving in a magnetic field. Spontaneous emission by electrons in fields of this configuration is of particular interest since such problems arise in calculations involving lasers, electron excitation, behavior of electron spin, and so on. In many respects, these problems are analogous to those arising in connection with the interaction between radiation and systems of atoms or molecules. The latter subject was reviewed in[1a]. The spontaneous emission of electrons moving in a constant uniform magnetic field and the field of a plane wave propagating in the direction of the magnetic field can be investigated in the greatest detail. This is so because the Dirac equation for an electron in a field of this kind has been solved exactly.[2]

Individual problems connected with the properties of radiation emitted by a relativistic electron in fields of this kind have been treated by the methods of classical electrodynamics in[1b, 3]. In particular, Bagrov and Khalilov[3] have derived the combination spectrum of frequencies emitted by an electron, and have obtained expressions for the spectral and angular distribution of the emitted radiation, and for the total radiated power. They also investigated the polarization of the emitted radiation. An analysis of the total cross section for the scattering of a strong wave by plasma electrons in a magnetic field, taking into account deceleration by radiation, is given in[1b]. Some aspects of this problem were considered quantum-mechanically in[4–11], in which equations were obtained for the radiation frequencies, and the problem of scattering of a weak wave by a relativistic electron in a magnetic field was investigated. In addition to the relativistic case, there is considerable interest in the interaction between a strong electromagnetic wave and electrons occupying low-lying energy states in a magnetic field. This interaction and the associated electromagnetic emission by electrons are investigated in the present paper. In particular, we report an analysis of the properties of the radiation emitted when an electron undergoes transitions between the $n=0,1$ levels. The effect of the field due to a strong wave on the behavior of the electron spin in this process is also discussed.

Consider a charge e moving in a uniform magnetic field H parallel to the z axis and the field of a circularly polarized plane wave propagating along the z axis with frequency $\omega_0 = c\varkappa_0$ and electric-field amplitude E_0 ($g=1$ and $g=-1$ will correspond to right-handed and left-handed polarizations, respectively). The wave function that is the exact solution of the Dirac equation is known[2] and can be written in the following form in terms of the two-dimensional Pauli matrices $\boldsymbol{\sigma}$:

$$\Psi = NL^{-1}\exp(-iS)\begin{pmatrix} G(1) \\ \varepsilon g G(-1)\end{pmatrix} v, \quad \varepsilon = \frac{eg}{|e|},$$
$$S = c\lambda t - kx - \frac{(\gamma y + \varepsilon g k) k_0 \gamma_0 \cos\varkappa_0\xi}{\lambda\varkappa_0(1+\varepsilon\delta)} + \frac{\beta\xi}{1-\beta} - \frac{\varepsilon\delta k_0^2\gamma_0^2 \sin 2\varkappa_0\xi}{4\varkappa_0\lambda(1+\varepsilon\delta)^2},$$
$$G(s=\pm1) = [(\lambda A^*\sigma_1 + \lambda + sk_0)U_{n-1}(a) + (2\gamma n)^{1/2}\sigma_1 U_n(a)](1-\varepsilon g\sigma_3) - s[(\lambda A\sigma_1 - \lambda - sk_0)U_n(a) + (2\gamma n)^{1/2}\sigma_1 U_{n-1}(a)](1+\varepsilon g\sigma_3), \quad (1)$$
$$a = \sqrt{\gamma}\, y + \frac{\varepsilon g k}{\gamma^{1/2}} + \frac{\varepsilon\sqrt{\gamma}\, k_0\gamma_0 \sin\varkappa_0\xi}{\lambda\varkappa_0(1+\varepsilon\delta)}, \quad \xi = ct - z,$$
$$k_0 = \frac{mc}{\hbar}, \quad \gamma = \frac{|e|H}{c\hbar}, \quad \gamma_0 = \frac{|e|E_0}{mc\omega_0}, \quad \delta = \frac{\gamma}{\lambda\varkappa_0}, \quad A = \frac{\varepsilon k_0\gamma_0 \exp(i\varepsilon\varkappa_0\xi)}{i\lambda(1+\varepsilon\delta)}.$$

In these expressions, v is an arbitrary two-component constant spinor, the specific choice of which determines the orientation of the electron spin. In particular, assuming that v is the eigenfunction of σ_3 ($\sigma_3 v = \zeta v$, $\zeta = \pm 1$), we take the spin to be parallel ($\zeta = 1$) or antiparallel ($\zeta = -1$) to the magnetic field. We also assume that v is normalized by the condition $v^+ v = 1$. At the same time, the wave function (1) is itself normalized to unity provided

$$|N|^2 = \frac{(1-\beta)\sqrt{\gamma}}{16\lambda^2}, \quad \beta = \frac{k_0^2 + 2\gamma n + k_0^2\gamma_0^2(1+\varepsilon\delta)^{-2} - \lambda^2}{k_0^2 + 2\gamma n + k_0^2\gamma_0^2(1+\varepsilon\delta)^{-2} + \lambda^2}.$$

The quantity c is the mean velocity of an electron along the z axis, $c\hbar\lambda = \mathscr{E} - cp_z$ is a constant of the motion in this particular field ($\mathscr{E}$ is the energy and p_z the z component of the electron momentum), and $U_n(x)$ are Hermite functions which are related to the Hermite polynomials $H_n(x)$ by the formula

$$U_n(x) = (2^n n! \sqrt{\pi})^{-1/2} \exp(-x^2/2) H_n(x), \quad n = 0, 1, 2 \ldots .$$

It is readily seen that, in the $n = 0$ state, the wave function differs from zero only when the electron spin ($e < 0$, $\varepsilon = -g$) is antiparallel to the magnetic field; the spin of a positron ($e > 0$, $\varepsilon = g$) can only be parallel to the magnetic field in this state. The presence of a plane wave does not, therefore, modify the particular spin properties of the $n = 0$ state, which are known in the case of a purely magnetic field.[12]

The emitted radiation can be calculated in first-order perturbation theory in the radiation field by the standard methods of quantum electrodynamics. The calculations are very laborious, but essentially simple, and will not be reproduced here. We merely note the following. The matrix elements of the emitted radiation can be evaluated exactly and can be expressed in terms of the Bessel functions for any given transition $n \to n'$, but we have not succeeded in obtaining a compact version of the necessary general expressions. We shall confine our attention here to a consideration of transitions from the $n = 0$ and 1 levels. Moreover, we note that the resulting exact expressions are still quite complicated, and we shall therefore carry out an expansion in terms of $\hbar$, retaining only the first nonvanishing terms. This does not signify a simple transition to the classical treatment because the very fact that we can isolate the $n \to n'$ transitions is possible only in quantum mechanics. The main argument in favor of this expansion is that the higher-order corrections in $\hbar$ are always small in practice. The characteristic parameters of the problem are the dimensionless quantities γ_0 and δ, where δ can be written in the following equivalent forms

$$\delta = \frac{\omega_H}{\omega_0} = (1-q^2)^{1/2} \frac{H}{H_0} \frac{mc^2}{\hbar\omega_0} \sqrt{\frac{1+\beta}{1-\beta}}, \quad q = \frac{\gamma_0}{(\Delta^2 + \gamma_0^2)^{1/2}},$$

$$H_0 = \frac{m^2c^3}{|e|\hbar}, \quad \Delta = 1 + \varepsilon\delta, \quad \omega_H = \frac{|e|Hc}{\mathscr{E} - cp_z}.$$

The fact that the quantum corrections are small is ensured by satisfying the following conditions: $H/H_0 \ll 1$ and $\hbar/mc^2 \ll 1$. Moreover, the case of exact resonance, $\Delta = 0$, is also excluded from our analysis. If we confine our analysis to magnetic fields such as are used in accelerators ($H \sim 10^4$ G), then for optical frequencies ω_0 we have $\delta \sim 0.01 - 0.0001 \ll 1$.

If we consider a transition from the state $k = 0$, λ, n, ζ to the state k', λ', n', ζ', with the emission of a photon of frequency $c\varkappa$ at an angle θ_0 to the z axis, we find that $\varkappa$ is given by the following equations:

$$2\varkappa_0\lambda\lambda' l = \lambda(k_0^2 + 2\gamma n' + k_0^2\gamma_0^2\Delta'^{-1}) - \lambda'(k_0^2 + 2\gamma n + k_0^2\gamma_0^2\Delta^{-1}) + \lambda\lambda'\varkappa(1+\cos\theta_0), \quad \lambda' = \lambda - \varkappa(1-\cos\theta_0), \tag{2}$$

derived in[4]. It is then a simple matter to show that equation for $\varkappa$ is a quartic. In the above expressions, l is an integer such that $\varkappa > 0$. If we solve (2) in the approximation $\hbar \to 0$, we find that

$$\varkappa = \frac{(1-\beta)\varkappa_0(l + \nu\delta)}{(1-\beta\cos\theta_0)}, \quad \nu = n - n', \tag{3}$$

and hence it follows that $l + \nu\delta > 0$, which defines the domains of l and ν.

When $l > 0$, the process can be interpreted as proceeding with the absorption of l photons from the wave; when $l < 0$, the photons are emitted into the wave. When $l > 0$ and $\nu \neq 0$, this process is usually referred to as combination (or shift) scattering. The emission probability can always be expressed in terms of the Bessel function $J_l(x)$ and its derivative $J_l'(x)$, where, in the $\hbar \to 0$ approximation, the argument x is given by

$$x = q(l + \nu\delta)\sin\theta, \quad \cos\theta = \frac{\cos\theta_0 - \beta}{1 - \beta\cos\theta_0}, \quad 0 \leqslant \theta \leqslant \pi. \tag{4}$$

Now consider the case $n = n' = 0$ (coherent, unshifted scattering). In this case, $\nu = 0$, the transition is possible only for $\zeta = \zeta' = \varepsilon g$, and its probability in the $\hbar \to 0$ limit is given by

$$W = \frac{e^2\omega_0}{\hbar c}(1-\beta)F_1, \quad F_1 = \sum_{l=1}^{\infty} l \int_0^{\pi} \sin\theta \, d\theta [q^2 J_l'^2 + \mathrm{ctg}^2\theta J_l^2], \tag{5}$$

which is exactly the same as the corresponding expression for the classical limit of the probability of emission by an electron in a plane circularly-polarized wave.[13] However, the quantity q depends both on the wave parameters and on the magnetic field and, moreover, this dependence is very important. For the function F_1 in (5), we can readily show, using the properties of Bessel functions and their approximations in terms of Macdonald functions,[14] that, in the limit:

$$F_1 = \frac{2}{3}q^2, \quad q \ll 1; \quad F_1 = \frac{55\sqrt{3}}{6}(1-q^2)^{-1/2} = \frac{55\sqrt{3}}{6}\left(1 + \frac{\gamma_0^2}{\Delta^2}\right)^{1/2}, \quad 1 - q^2 \ll 1; \tag{6}$$

where, for $q \ll 1$, the main contribution to F_1 is the term with $l = 1$, and for $1 - q^2 \ll 1$, the main contribution is due to terms with $l \sim (1 - q^2)^{-3/2} \approx \gamma_0^3 |\Delta|^{-3}$ (see[14]). It is also known that F_1 is a monotonically increasing function of q.

It is clear that the condition $1 - q^2 \ll 1$ is equivalent to $\gamma_0^2 \gg \Delta^2$, which can be realized in two cases, namely, a) in the region near resonance ($\varepsilon = -1$, $\delta \sim 1$, Δ close to zero), and b) in the case of a strong wave, i.e., when the parameter γ_0 is large. It is clear from (5) that, as $q \to 1$, the probability of emission of a given harmonic l ceases to depend on q and tends to a finite limit which is a function of l but, at the same time, according to (6), the sum diverges as $(1-q^2)^{-1/2}$.

In the other limiting case, $q \ll 1$, which is equivalent to $\gamma_0^2 \ll \Delta^2$, we see that the scattering cross section

$$\sigma = \frac{4\pi\hbar\omega_0 W}{cE_0^2} = \frac{8\pi}{3}(1-\beta)\frac{r_0^2}{\Delta^2}, \quad r_0 = \frac{e^2}{mc^2} \tag{7}$$

is independent of the wave amplitude and, naturally, in the absence of a magnetic field ($\Delta = 1$), it is identical with the Thomson cross section. It is clear that this situation is analogous to that described in[1a] for systems determined by the quasienergy. We also note the following. When $\varepsilon = 1$ (nonresonant motion), the probability given by (5) decreases with increasing magnetic field H for fixed wave parameters. Consequently, in this case, the coherent scattering cross section is less than the Thomson cross section. If, on the other hand, $\varepsilon = -1$ (resonance is possible), both the probability (5) and the cross section (7) increase with increasing H up to the point $\delta < 1$, where the cross section (7) becomes formally infinite (we note that, near the resonance $\delta = 1$, the damping theory must be used in calculating physical quantities such as the cross section), and then decrease with increasing H. Thus, even when $\delta > 1$, the coherent scattering cross section of electrons in a magnetic field is substantially lower than the Thomson cross section.

Let us now consider combination emission (spontaneous emission of a quasienergy system). In the limit of $\hbar \to 0$, the $n=1 \to n'=0$ $(\nu = 1)$ and $n = 0 \to n' = 1$ $(\nu = -1)$ transitions occur without spin flip. We then have $\zeta = \varepsilon g$, and the probability of such transitions is

$$W = BF_2, \quad F_2 = \sum_{l=l_0}^{\infty} (l+\nu\delta) \int_0^{\pi} \sin\theta\, d\theta \left\{ \left[\left(\frac{\varepsilon\delta l(l-\nu\varepsilon)}{\Delta(l+\nu\delta)(1+\cos\theta)} + (l+\nu\delta)\cos\theta - \frac{\varepsilon\delta l}{\Delta} \right) J_l - \frac{\nu\delta}{\Delta}\left(1 - \frac{l-\nu\varepsilon}{(l+\nu\delta)(1+\cos\theta)}\right) xJ_l' \right]^2 \right.$$
$$\left. + \left[\left(1 - \frac{\varepsilon\delta(l-\nu\varepsilon)}{\Delta(l+\nu\delta)(1+\cos\theta)}\right) xJ_l' - \varepsilon\delta\left(1 - \frac{\nu\varepsilon(x^2-l^2+\nu\varepsilon l)}{\Delta(l+\nu\delta)(1+\cos\theta)}\right) J_l \right]^2 \right\}. \tag{8}$$

$$B = \frac{\sqrt{1-\beta^2}}{2T_0}\frac{H_0}{H}\left(\frac{\hbar\omega_0}{mc^2}\sqrt{\frac{1-\beta}{1+\beta}}\right)^3 = \frac{\sqrt{1-\beta^2}}{2T}\frac{\tilde{H}}{H}\left(T\omega_0\sqrt{\frac{1-\beta}{1+\beta}}\right)^3,$$
$$T_0 = \frac{\hbar^2}{e^2mc}, \quad T = \frac{r_0}{c}, \quad \tilde{H} = \frac{e}{r_0^2}.$$

In these expressions, $l_0 = -\nu[\delta]$, where $[\delta]$ is the integral part of δ. It is readily seen that this probability is, in general, independent of Planck's constant $\hbar$. When $\nu = 1$, $l = 0$, and $\gamma_0^2 \ll \Delta^2$, it follows from (8) that the spontaneous emission probability in a purely magnetic field[12] is

$$W = 2T_{\rm cl}^{-1}, \quad T_{\rm cl} = \frac{3}{2}T\left(\frac{\tilde{H}}{H}\right)^2 = \frac{3}{2}\frac{m^3c^5}{e^4H^2}, \tag{9}$$

where $T_{\rm cl}$ is known from classical theory and is defined as the electron emission time in a magnetic field. When $l \neq 0$, we find that, in the most interesting case when $\delta \ll 1$, F_2 is given by (8) as follows:

$$F_2 = \sum_{l=1}^{\infty} l^2 \int_0^{\pi} \sin\theta\, d\theta (\cos^2\theta J_l^2 + q^2 \sin^2\theta J_l'^2),$$

and hence

$$F_2 = \frac{2}{5}\gamma_0^2, \quad \gamma_0 \ll 1; \quad F_2 = \frac{55\sqrt{3}}{72}\gamma_0^7, \quad \gamma_0 \gg 1. \tag{10}$$

Consequently, when $\delta \ll 1$, transitions with $\nu = \pm 1$ are roughly equally probable, and this leads to the following conclusion: when a wave is incident on an electron in a magnetic field, this results in an effective excitation of the electron. It is clear that the number of excited levels is $\Delta n \sim \delta^{-1}$.

The emission probability decreases with increasing γ_0 when $l = 0$, and increases when $l \neq 0$. These probabilities are equal for a certain value of γ_0. It follows from (9) and (10) that this is achieved (for $\delta \ll 1$) when $10\delta^3 = 3\gamma_0^2$. Since for modern lasers $\gamma_0 \sim 0.01$, we find that $\delta \sim 0.01$. The excitation time $T_{\rm cl}$ for fields $H \sim 10\,000$ G is then ~ 5 sec.

It follows from (8) that, when $\delta \gg 1$, transitions involving the excitation of the $n = 1$ $(\nu = -1)$ level are possible only if accompanied by the absorption of a large number of photons $l \sim \delta$ from the wave. The probability of such transitions is small for $\gamma_0 \ll 1$.

The above analysis of the spontaneous emission of the system can be extended to stimulated processes. Some results in this field can be found in[1a, 15].

There is particular interest in the behavior of the electron spin during the emission of radiation.

It follows from[12] that, in a time $t > T_{\rm cl}$, all the electrons moving in a magnetic field will emit their energy by radiation, and all the electrons with spins antiparallel to the magnetic field will be in the ground state with $n = 0$. If, on the other hand, the spins are parallel to the magnetic field, the electrons end in the $n = 1$ state. Transition from the $n = 1$ level to the ground state $n = 0$ is then of low probability (it occurs in a time $t > \tau = T_0(H_0/H)^3$, $\tau \sim 10^{11}$ sec for $H \sim 10\,000$ G), since this transition must be accompanied by spin flip. Consequently, if the magnetic field contains a system of excited electrons that, on the average, are unpolarized, then, after emission of radiation, i.e., after a time $t > T_{\rm cl}$, the system will divide into two phases, namely, electrons with spins antiparallel to the field in the ground state $n = 0$ and those with spins parallel to the field in the first excited state $n = 1$.

It is interesting to consider how this situation will change in the presence of a strong wave. The probability of transitions with spin flip (initial spin $\zeta = -\nu\varepsilon g$) for $\hbar \to 0$ is

$$W = B\frac{H}{H_0}(1-q^2)^2 F_3,$$
$$F_3 = \sum_{l=l_0}^{\infty} (l+\nu\delta)^3 \int_0^{\pi} (1-\cos\theta)^2 J_l^2 \sin\theta\, d\theta. \tag{11}$$

When $l = 0$, $\nu = 1$, and $\delta \ll 1$, this yields

$$W=W_0=\frac{4\sqrt{1-\beta^2}}{3T_0}\left(\frac{H}{H_0}\right)^3(1-q^2)^{7/2},$$

so that, when $\gamma_0 \to 0$ and $\beta=0$, we have the result given in[12].

Assuming that $q \ll 1$, $\delta < 1$, we find from (11) that

$$W=W_0+W_1,$$

$$W_1=\frac{B}{5}\frac{H}{H_0}q^2(1+\nu\delta)^5.$$

For fields $H \sim 10\,000$ G and $\gamma_0 \sim 0.01$, we have $W_0 \sim 10^{-11}$ sec^{-1}, $W_1 \sim 10^{-5}$ sec^{-1}. Thus, the presence of a plane wave reduces the electron spin-flip time by five or six orders of magnitude, but this time remains very long ($\sim 10^5$ sec) so that the $n=1$, $\zeta=-\varepsilon g$ state remains stable against transition to the $n=0$ state.

In the other limiting case, when $1-q^2 \ll 1$, we find from (11) that, when $\delta<1$,

$$W=B\frac{H}{H_0}\frac{35\sqrt{3}}{72}(1-q^2)^{-1/2}\left(1+\frac{192\sqrt{3}}{175}\nu\delta\sqrt{1-q^2}\right),$$

i.e., the $n=1$, $\zeta=-\varepsilon g$ state is no longer stable in the presence of a strong wave.

Thus, when the wave is present, there is a range of levels $\Delta n \sim \delta^{-1}$ ($\delta<1$) for $t>T_{c1}$ which is filled with electrons but, if the electron spin is antiparallel to the field, the lowest level is $n=0$; when the electron spin is parallel to the field, the lowest level is $n=1$.

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Mechanism of isotopically selective dissociation of SF_6 molecules by CO_2 laser radiation

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An experimental study is reported of the selective multiphoton dissociation of SF_6 in a strong infrared laser field. The frequency characteristics of the rate of dissociation and of the enrichment coefficient, the threshold charcteristics of the dissociation process, the quantum yield, and the effect of collisions on the rate and selectivity of SF_6 dissociation have all been determined. The mechanism of selective dissociation of the SF_6 molecule by a strong infrared laser field is discussed. A new method is proposed for the selective dissociation of molecules by a two-frequency infrared laser field.

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1. INTRODUCTION

The phenomenon of isotopically selective dissociation of polyatomic molecules by strong CO_2 laser pulses has recently been discovered (BCl_3,[1] SF_6,[2] OsO_4,[3] etc.). This discovery was preceded by studies of the visible and ultraviolet luminescence of molecular gases (C_2F_3Cl, SiF_4, BCl_3, etc.) under the action of a focused CO_2 laser pulse.[4–6] The desire to understand the nature of this emission and to demonstrate that it can appear as a result of the resonance interaction between a strong infrared field and a molecule, without the necessary participation of collisions with other molecules, has led to systematic studies of luminescence in isotopic mixtures of molecules[1] and to the discovery of the above phenomenon. Subsequent experiments have shown that, on the one hand, luminescence does not necessarily occur during the primary dissociation process but is frequently the result of secondary processes (see, for example, the experiments with trans-dichloroethylene[7])