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Čerenkov and transition radiations in the γ-resonance frequency region

V. M. Samsonov

Leningrad Institute of Nuclear Physics, USSR Academy of Sciences
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Coherent (Čerenkov and transition) radiation of a charged particle moving through a bounded medium with resonantly scattering nuclei is considered. It is shown that an essential role in the formation of radiation in the γ-resonance frequency region is played by allowance for the smallness of the radiation-absorption length in the medium, and that the interference between the Čerenkov and transition radiations in the vacuum must be taken into account without fail. Attention is called to the existence of a region of emission frequencies and angles in which the waves from the sections of particle motion in the medium and in the vacuum are in phase. The possible quantum yields of coherent emission in the γ band are presented for media containing the Mössbauer isotopes ^{57}Fe, ^{119}Sn, and ^{73}Ge.

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1. It is known that resonant scattering of γ quanta by nuclei exerts an appreciable influence on the character of the frequency dispersion and on the refractive index of the medium at frequencies close to the resonant frequency of the nucleus, since the coherent amplitude of forward nuclear scattering can be comparable and may even exceed the electronic scattering amplitude. It therefore becomes possible to obtain Čerenkov radiation[2,4] and to increase the transition radiation[1,3] (in an inhomogeneous medium) in the γ band in the case of a fast charged particles moving in a medium containing resonantly scattering nuclei.

The estimates of the intensity of the Čerenkov radiation in Refs. 2 and 4 pertain to an unbounded medium and describe the additional contribution made to the energy losses of the fast charged particle in the resonantly scattering medium, whereas practical interest attaches to coherent radiation emitted into vacuum (or into another weakly absorbing medium).

For such an analysis it becomes important to consider the following singularities of the case in question: the short absorption length in the γ-resonance frequency region, and the fact that the coherence lengths in vacuum and in the medium are comparable. Account must therefore be taken of the limited region in which the Čerenkov radiation is formed in the medium, and the interference between this radiation and the transition radiation in the vacuum.

Allowance for these circumstances, and the consequences ensuing from such an analysis, are in fact the subject of the present paper.

2. The refractive index of a medium at frequencies close to the γ-transition frequency in the nucleus can be expressed in the form[5,6]

$$n=1+\delta n_{el}+\delta n_{nuc}$$
$$\delta n_{el}=-Z_0 r_e N\lambda^2/2\pi,\quad \delta n_{nuc}=-4a\Delta/(\Delta^2+1)+i\cdot 4a/(\Delta^2+1); \tag{1}$$
$$a=N(2j'+1)f^2\lambda^3/32\pi^2(2j+1)(1+\alpha_0),$$

where δn_{el} is the characteristic part of the refractive index,[1)] δn_{nuc} is the nuclear component of the refractive index, r_e is the classical radius of the electron, N is the number of atoms per unit volume, $\lambda(\omega)$ is the wavelength (frequency) of the γ quantum, Z_0 is the atomic number, j' and j are respectively the spins of the excited and ground states of the nucleus, f is the Lamb-Mössbauer factor (assumed hereafter equal to unity), α_0 is the internal conversion coefficient, ω_0 is the frequency of the resonant level of the nucleus, Γ is the total width of the nuclear level, and $\Delta=2(\omega-\omega_0)/\Gamma$.

According to (1), for example, for a medium of ^{57}Fe (nuclear-transition energy $E_\gamma=14.4$ keV, $\Gamma=4.67\times10^{-9}$ eV) we have

$$\delta n_{el}=-7.5\cdot10^{-6},\quad a=3.37\cdot10^{-5},$$

for ^{119}Sn ($E_\gamma=23.9$ keV, $\Gamma=2.47\times10^{-8}$ eV)

$$\delta n_{el}=-1.71\cdot10^{-6},\quad a=4.61\cdot10^{-6}$$

and for ^{73}Ge ($E_\gamma=67.03$ keV, $\Gamma=2.45\times10^{-7}$ eV)

$$\delta n_{el}=-2.16\cdot10^{-7},\quad a=9.2\cdot10^{-7}.$$

To calculate a we used the data of Ref. 7 on the Mössbauer isotopes.

3. Let a particle with charge e move uniformly with

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velocity **v** from a medium with refractive index n defined by Eq. (1), in a vacuum and in a direction perpendicular to the interface. The field of the emission of frequency ω produced at the point of observation in vacuum by such a particle is then determined by a superposition of the fields[2)] $\Pi_{1\omega}$ and $\Pi_{2\omega}$:

$$\Pi_{\omega}=\Pi_{1\omega}+\Pi_{2\omega},$$

where $\Pi_{1\omega}$ is the field (Fourier component of the Hertz vector with frequency ω) produced at the point of observation by the section of the particle trajectory in the medium; $\Pi_{2\omega}$ is the field from the vacuum section of the particle path. If the particle mean free path in the medium is L, the field $\Pi_{1\omega}$ can be expressed in the form[8]

$$\Pi_{1\omega}=\int_0^L \mathbf{P}_{\omega}\frac{e^{i\varphi}}{R_1}dz, \tag{2}$$

where

$$\mathbf{P}_{\omega}=-\frac{ie\mathbf{v}}{2\pi\omega v}\frac{2\cos\theta}{[n^2\cos\theta+(n^2-\sin^2\theta)^{1/2}]}, \tag{3}$$

$$\varphi=\frac{\omega}{c}R+\frac{\omega}{v}[1-\beta(n^2-\sin^2\theta)^{1/2}]z+\frac{\omega z^2\sin^2\theta\cos^2\theta}{2c(n^2-\sin^2\theta)R}, \tag{4}$$

$$R_1=R\left[1-\frac{z\cos^2\theta}{2(n^2-\sin^2\theta)^{1/2}R}\right]. \tag{5}$$

Here R is the distance from the interface to the observation point, θ is the angle between the velocity vector **v** and the direction of R, and $\beta=v/c$.

The field $\Pi_{1\omega}$ can also be regarded as the result of the interference of waves emitted by the so-called "dipole images"[8] placed in the vacuum with a distribution of the dipole moment $\mathbf{P}_{\omega}$ along the particle trajectory; in this case φ is the phase of the wave emitted by an oscillator with coordinate z.

If we put

$$R\gg L^2\frac{\sin^2\theta\cos^2\theta}{\lambda(n^2-\sin^2\theta)}, \tag{6}$$

a relation easily satisfied in the γ-resonance frequency region, since the possible sample thickness is limited by the radiation-absorption length, then we can discard the term of (4) quadratic in z. Such a condition (6) means that within the limits of the particle mean free path in the medium, given R, the change of the phase of the waves on account of the quadratic term is much less than π and therefore allowance for this term does not alter significantly the summary amplitude $\Pi_{1\omega}$.

Since the energy range of the γ radiation in this case is 10^4-10^5 eV, the refactive index is in this case close to unity ($\delta n=\delta n_{el}+\delta n_{nuc}\approx 10^{-4}-10^{-7}$), and the energy of the particles (electrons) exceeds 10^7, the expression for $\Pi_{1\omega}$ (with (6) taken into account) assumes a particularly simple form

$$\Pi_{1\omega}=-\frac{ev}{\pi\omega^2R}\frac{\exp(i\omega R/c)}{2\alpha-\beta(\delta n-\theta^2)}\left\{1-\exp\left[-i\frac{\omega L}{2v}(2\alpha-\beta(2\delta n-\theta^2))\right]\right\}, \tag{7}$$

where $\alpha=1-\beta$.

We can analogously obtain the field $\Pi_{2\omega}$, which, at a particle mean free path in vacuum much larger than the coherence length,[3)] takes the form[8]

$$\Pi_{2\omega}=\frac{ev}{\pi\omega^2R}\frac{\exp(i\omega R/c)}{(2\alpha+\theta^2)}.$$

Knowing $\Pi_{1\omega}$ and $\Pi_{2\omega}$, we can find the expression for the spectral-angular distribution of the γ-radiation intensity:

$$\frac{d^2W}{d\omega\,d\Omega}=\frac{\omega^4R^2\theta^2}{c^3}|\Pi_{1\omega}+\Pi_{2\omega}|^2, \tag{8}$$

where $d\Omega=2\pi\theta d\theta$.

4. Let us see how the field amplitude (7) varies with the particle mean free path in the medium. This can be easily done by representing (7) in the form

$$\left|\frac{\Pi_{1\omega}(L)}{\Pi_{1\omega}(\infty)}\right|=2\exp\left(-\frac{\pi}{2}L'\right)\left(\sin^2\frac{\pi}{2}kL'+\mathrm{sh}^2\frac{\pi}{2}L'\right)^{1/2} \tag{9}$$

$$k=\frac{2\alpha+\beta(\theta^2-2\,\mathrm{Re}\,\delta n)}{2\beta\,\mathrm{Im}\,\delta n}$$

$$L'=L/L_{abs},\quad L_{abs}=\lambda/2\,\mathrm{Im}\,\delta n. \tag{10}$$

Here $\Pi_{1\omega}(L)$ is the field amplitude at mean free path L of the particle in the medium, $\Pi_{1\omega}(\infty)$ is the field amplitude (7) at an infinite mean free path of the particle in the medium. Figure 1 shows a plot of (9) for ^{119}Sn against the relative mean free path L' at different values of the angles and frequencies of the emission (in terms of the parameter k). It is seen from this figure that the field amplitude (7) remains practically unchanged with increasing $L>L_{abs}$ at any value of k, and consequently L_{abs} in the form defined in (10) is a good characteristic of the "visible" part (limited by the γ-ray absorption) of the particle trajectory, and is called the absorption length.

The absorption length depends on the particle of the emitted γ quantum in the following manner:

$$L_{abs}=L_0(\Delta^2+1),$$

where Δ is the difference between the γ-quantum frequency and the resonant frequency, and is expressed in terms of the half-widths of the nuclear level; $L_0=\lambda/4a$ characterizes the absorption length at the resonant frequency. The length L_0 is 1.55×10^{-4}, 3.2×10^{-5}, and 2.5×10^{-4} cm for ^{119}Sn, ^{57}Fe, and ^{73}Ga, respectively. It is seen from these figures that the possible sample thickness is very small and the condition (6) will certainly by satisfied. To obviate the need for introducing

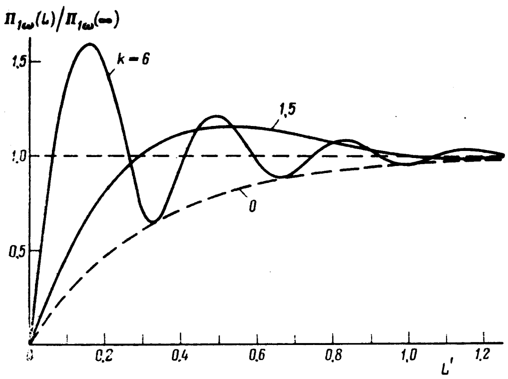

FIG. 1. Variation of the field amplitude $|\Pi_{1\omega}|$ as a function of the relative mean free path L' at different values of the parameter k.

the radiation from the second boundary, we shall henceforth assume the sample to be thicker than L_{abs} in the entire frequency band of interest. We can then discard in (7) the L-dependent term.

An important feature that determines the amplitude of the field produced by the charge at the observation point is the coherence length (see footnote 3). For an absorbing medium, however, it is convenient to introduce an "effective" coherence length, determined from (7) as follows:

$$L_{coh}^{(med)}=\frac{\lambda\beta}{\{[2\alpha+\beta(\theta^2-2\,\mathrm{Re}\,\delta n)]^2+4(\mathrm{Im}\,\delta n)^2\}^{1/2}}. \quad (11)$$

This takes into account the absorption of the radiation in the medium. That is to say, if, for example, the condition

$$2\alpha+\beta(\theta^2-2\,\mathrm{Re}\,\delta n)=0 \quad (12)$$

is satisfied, then the coherence length is equal to the absorption length. In other words, the quantity $L_{coh}^{(med)}$ as defined in (11) assumes the role of the size of the radiation source at the given frequency and at the given observation angle. The larger the source, the larger the summary radiation amplitude. Figure 2 shows plots of the coherence length in the medium as functions of the radiated frequency and of the radiation direction in the region of the γ-resonance frequencies of ^{119}Sn (electron energy $E_e=500$ MeV; for convenience we have introduced the relative emission angles θ' in units of $2\alpha(\theta'=\theta^2/2\alpha)$). In all cases in which condition (12) is not satisfied, the coherence length is substantially smaller than the absorption length. It is this which explains the oscillations of the amplitude on Fig. 1 (the curves $k\neq 0$) due to inclusion in the "visibility" range, of regions that emit waves having phase shifts larger than π (i.e., several Fresnel zones are included[9]). In the entire range of frequencies and angles that satisfy the condition (12), the radiation arrives at the observation point in phase from the entire "visible" section of particle motion in the medium. This is a reflection of the the monotonic course of the curve at $k=0$ on Fig. 1.

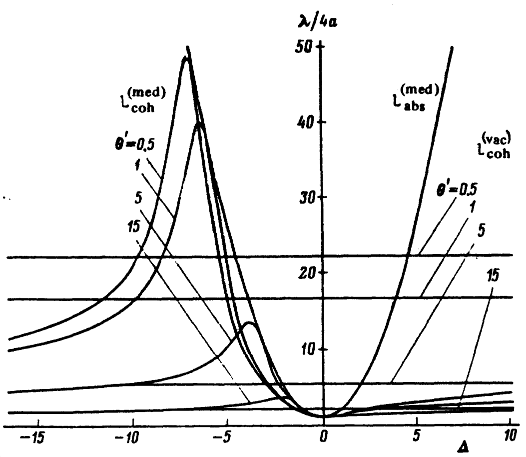

FIG. 2. Changes of coherence length of the radiation in the medium and in vacuum ($L_{coh}^{(med)}$ and $L_{coh}^{(vac)}$) against the emission frequency Δ and the observation angle θ' (^{119}Sn, $E_e=500$ MeV); $L_{abs}^{(med)}$ is the absorption length in the medium.

If we introduce the coherence lengths for the vacuum

$$L_{coh}^{(vac)}=\lambda\beta/(2\alpha+\beta\theta^2),$$

then we see that in the region of the resonance for ^{119}Sn they are comparable with the coherence length in the medium. Therefore the formation of the summary field is greatly influenced by the vacuum section of the particle motion.

Figure 3 shows the spectral distribution of the radiation intensity (8) when the electron ($E_e=500$ MeV) crosses the interface between a medium consisting of ^{119}Sn and vacuum, for different observation angles (solid curves). For $\theta'=1$ we show by way of example (dashed curve) the spectral distribution of the intensity of the radiation from the medium, a distribution defined by

$$\frac{d^2W}{d\omega\, d\Omega}=\frac{\omega^4R^2\theta^2}{c^3}|\Pi_{1\omega}|^2.$$

Comparing with the curve $\theta'=1$ for the summary intensity (8), we see that at $\Delta<-6$ the intensity at the observation point decreases, and at $\Delta>-6$, conversely, the intensity increases as a result of the addition of the amplitudes of the fields $\Pi_{1\omega}$ and $\Pi_{2\omega}$. It is easily seen from (7) that the intensity can increase only in the frequency and angle region where

$$[2\alpha+\beta(\theta^2-2\,\mathrm{Re}\,\delta n)]\leqslant 0. \quad (13)$$

The maximum of the angular distribution of the γ-ray intensity in ^{119}Sn is at $\theta'_{max}\approx 1-2$, i.e., $\theta_{max}\approx(1-2)\sqrt{2\alpha}$, and at $E_e=1$ GeV we have $\theta_{max}\approx 5\times 10^{-4}$.

5. We illustrate the foregoing in the following manner. The field of the radiation that emerges from the medium (7) can be regarded as the result of the emission of an image charge e_1 moving in vacuum with velocity v_1 on the segment L_1[8]:

$$e_1=ev/v_1,\quad v_1=v(1+\delta n-{}^1/_2\theta^2)/(1-{}^1/_2\theta^2),$$
$$L_1=L(1+\delta n-{}^1/_2\theta^2)/(1-{}^1/_2\theta^2).$$

Introduction of the image-charge parameter simplifies the analysis of the radiation emerging to the vacuum. By way of example, Fig. 4, shows the regions of the γ-photon angles and frequencies in which $v_1<v$, $v<v_1<c$, and $v_1>c$ for a medium of ^{119}Sn ($E_e=500$ MeV). The curve in the $v_1>c$ region separates the angle and fre-

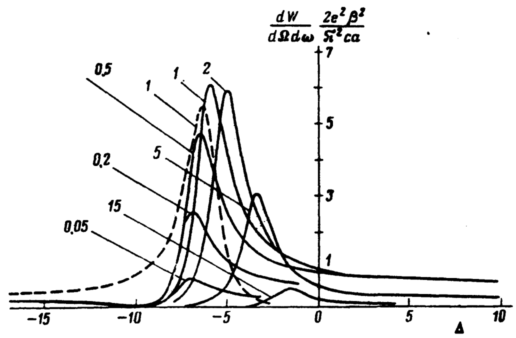

FIG. 3. Spectral distribution of the emission intensity for different observation angles θ' (solid lines; the figures on the curves are the values of θ') and for the angle $\theta'=1$ from the medium, without allowance for the interference with the vacuum radiation (dashed line (^{119}Sn, $E_e=500$ MeV)).

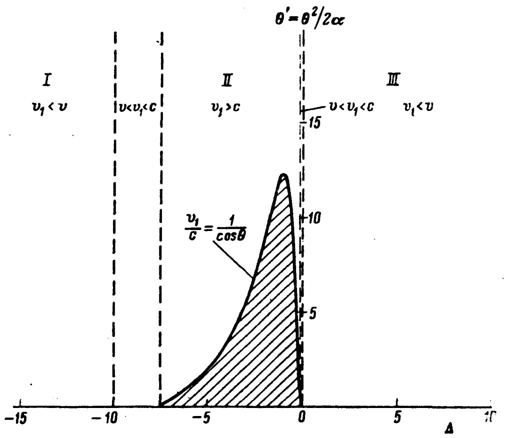

FIG. 4. Classification of the type of coherent radiation from the medium in terms of the velocity v_1 of the image charge (^{119}Sn, E_e = 500 MeV); region II—Cerenkov radiation, remaining regions—transition radiation.

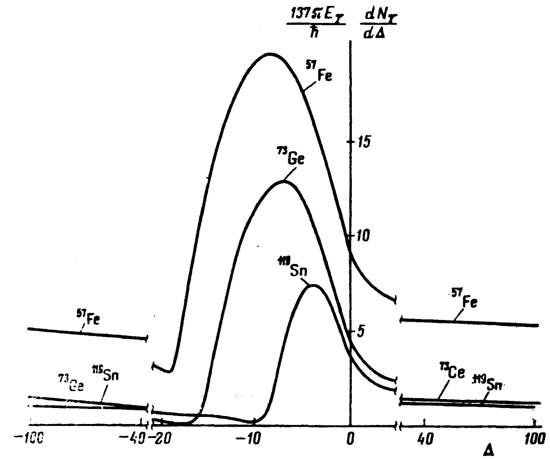

FIG. 5. Integral-spectral distribution of the coherent-radiation intensity for ^{57}Fe (E_γ=14.4 keV, E_e = 5 GeV), ^{73}Ge (E_γ=67 keV, E_e = 3 GeV), and ^{119}Sn (E_γ=24 keV, E_e=1 GeV).

quency region in which is satisfied the relation

$$n'\beta_1 = 1/\cos\theta,$$

which assumes in this case the form (12) (here n' is the refractive index of the vacuum and is equal to unity, while $\beta_1 = v_1/c$). This corresponds to the region where Cerenkov radiation propagates in the form of a cylindrical wave at small R (or where there is no absorption at large R). Since we assume that condition (6) is satisfied, where $L \approx L_{\text{abs}}$, it follows that the field $\Pi_{1\omega}$ exists at such distances in the form of a spherical wave, i.e., it fills the entire angle region (a situation analogous to the wave propagation after diffraction by a slit). This is seen also in Fig. 4 in the vicinity $v_1 > c$ for the field (7). Consequently, the region $v_1 > c$ is the region of frequencies and angles in which the Cerenkov radiation emerges from the region. The shaded region under the curve $\beta_1 = 1/\cos\theta$ on Fig. 4 corresponds to satisfaction of condition (13), i.e., this is the region of values of frequencies and angles where the Cerenkov emission and the transition radiation on the vacuum section arrive in phase at the observation point. It can therefore not be stated that the spectral-angular distributions of the intensity are due to Cerenkov radiation or only to transition radiation.

6. Expression (8) can be integrated over the angle θ; this yields the integral-spectral distribution of the radiation intensity

$$\frac{dW}{d\omega} = \frac{e^2}{\pi c}\left\{\frac{1}{2}\left(1 - \frac{2\alpha \operatorname{Re}\delta n}{|\delta n|^2}\right)\ln\frac{(\alpha - \operatorname{Re}\delta n)^2 + (\operatorname{Im}\delta n)^2}{\alpha^2} + \left(\frac{\operatorname{Re}\delta n}{\operatorname{Im}\delta n} + \frac{\alpha}{\operatorname{Im}\delta n}\,\frac{(\operatorname{Im}\delta n)^2 - (\operatorname{Re}\delta n)^2}{|\delta n|^2}\right)\left[\frac{\pi}{2} - \operatorname{arctg}\frac{\alpha - \operatorname{Re}\delta n}{\operatorname{Im}\delta n}\right] - 1\right\}.$$

Figure 5 shows the spectral distributions, plotted in accordance with this formula, in the region of the γ-resonance frequencies for a number of isotopes: ^{57}Fe (E_e = 5GeV), ^{119}Sn (E_e = 1 GeV), and ^{73}Ge (E_e = 3 GeV). The choice of the isotopes and energies of the electrons was determined for convenience in the comparison with the estimates in Refs. 1–4.

As seen from Fig. 5, the maximum of the radiation intensity is shifted away from the resonant frequency and does not coincide with the maximum of the refractive index ($\Delta = -1$). Another characteristic feature is the appreciable width of the distribution. These distributions can be used to estimate the quantum yields in the region of the γ-resonance frequencies. In the case of ^{57}Fe we have $N_\gamma = 1.2 \times 10^{-13}$ quanta/electron (the Cerenkov losses are estimated in Ref. 2 at 6.3×10^{-15} photon/electron over a length $L \approx 10^{-5}$ cm), for ^{119}Sn we get 7.3×10^{-14} photon/electron (a value 10^{-15} photon/electron is cited in Refs. 1–3, but the possibility of Cerenkov radiation is not considered there), and for ^{73}Ge we get $N_\gamma = 6.8 \times 10^{-13}$ photon/electron (in Ref. 4 the Cerenkov loss over a length $L \approx 10^{-3}$ is estimated at 1.6×10^{-13} photon/electron). At a current 10 μA (6.3×10^{13} electron/sec) the quantum yield can amount to several dozen quanta per second, and this is of interest from the experimental point of view.

Attention must be called here to the fact that if the sample thickness is made of the order of the absorption length at the frequency up to which we still have $v_1 > c$, then the rear boundary makes no contribution to the summary radiation in this region ("thick" sample), whereas in other frequency regions, with increasing distance from the resonance frequency, this influence will become more and more noticeable (the sample becomes "thin"). As a result of this influence, the background due to the transition radiation on Fig. 5 outside the region of the γ resonance will turn out to be practically suppressed. By the same token, we attune ourselves, as it were, to a definite spectral region of the radiation.

7. The foregoing analysis shows that the formation of coherent radiation of a fast charged particle in a medium with resonantly scattering nuclei, in the γ-resonance frequency region, is strongly influenced by the short radiation-absorption length, so that the range of angles at which the Cerenkov radiation emerges from the medium is broadened. Interference between the Cerenkov radiation and the transition radiation in vacuum causes realignment of the spectral distribution

of the radiation intensity. A frequency and angle region can then exist, in which the vacuum radiation and the radiation from the medium are in phase. We wish to call particular attention to this circumstance, since it can take place both in oblique incidence of the particle on the sample and when the particle moves in the channel.

This analysis is not restricted to the γ-resonance region, and can hold also at frequencies close to the natural frequencies of the medium (optical-transition frequencies, absorption edges,[10] and others)[4)].

In conclusion, the author thanks V. V. Fedorov for useful remarks, A. F. Tyunis for help with the numerical calculations, and T. B. Mezentseva for help in readying this paper for publication.

[1)]We neglect the imaginary part, which is smaller by almost two orders of magnitude than the real part in the frequency band of interest to us.

[2)]In this case we omit the third component—the field of the dipole images in the medium as the particle moves in vacuum, since this component is small compared with $\Pi_{1\omega}$ and $\Pi_{2\omega}$.

[3)]The trajectory region in which the phases of the waves emitted from it differ by not more than π.

[4)]The γ-resonance region was chosen because it was simplest to describe.

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Interaction of two-level system with a strong monochromatic wave and with a thermostat

T. K. Melik-Barkhudarov

Erevan State University
(Submitted 8 February 1978)
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An expression is obtained for the stationary distribution of a two-level system, produced as a result of interaction with the thermostat, in the field of a strong monochromatic wave; the times to establish this distribution are also determined. It is shown that when the external-field frequency is lower than a certain critical value the state of the two-level system is described by a Boltzmann distribution in the quasi-energy, and at high frequencies it contains matrix elements of the interaction with the thermostat. It is also shown that the corresponding relaxation constants are half as large in the resonant case as in the nonresonant one, so that it can be concluded that the presence of the signal suppresses the noise in a quantum amplifier.

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Progress in the theory of irreversible processes raises the question of the behavior of quantum-mechanical objects that are coupled to a thermostat, i.e., to a system that has in the limit an infinite number of degrees of freedom. Since this problem is extremely complicated, solutions can be expected for only the very simplest models, as was first done by Weisskopf and Wigner in their classical work on radiative damping.[1] In 1963, Gordon, Walker, and Louisell[2] solved the problem for a damped harmonic oscillator coupled to a set of independent oscillators distributed at the initial instant in accord with a canonic ensemble.[2] Glauber has shown later[3] that "large systems not consisting, of course, of harmonic oscillators have very frequently collective-excitation modes whose amplitudes behave dynamically like oscillator amplitudes," so that such a thermostat model is quite general. In the present paper, an attempt is made to develop further the Weisskopf-Wigner theory to include a two-level system in the field of a strong monochromatic wave, using the aforementioned model of the thermostat.

The first questions raised when this problem is formulated are:

1. Are there any memory effects whatever in the stationary distribution of a two-level system?

2. Does this distribution depend on the details of the mechanism of the interaction with the thermostat, or is the latter only a temperature, as is postulated in equilibrium statistical physics?

We describe our system by the Hamiltonian

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