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INFLUENCE OF COHERENCE OF RESONANCE ABSORPTION OF γ QUANTA IN POLYCRYSTALLINE SUBSTANCES

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Resonant absorption of γ quanta in polycrystals is considered in the general case with allowance for propagation of nuclear excitation from one nucleus to another. The collective excited state which arises in this case as a result of coherence possesses parameters that differ significantly from those in the case of an isolated nucleus (individual approximation). Consequently, the resonance absorption cross section cannot be described by the usual Breit-Wigner formula. The resonance line width and position vary. At $a \gg \lambda$ (a is the interatomic distance and λ the photon wavelength) and at high temperatures our results go over to the Breit-Wigner formula.

IN problems of γ resonance, the effective cross section of the resonant absorption of γ quanta by nuclei is calculated in accordance with the well known formula

$$\sigma(E)=\frac{2I+1}{2J+1}\frac{\lambda^2}{2\pi}\frac{\Gamma_R}{\Gamma}\frac{\Gamma^2/4}{(E-E_0)^2+\Gamma^2/4}, \tag{1}$$

where λ and E_0 are the wavelength and the energy of the resonant γ quanta, J and I are the spins of the ground and excited states of the nucleus, Γ and Γ_R are the total and radiative widths of the excited level of the nucleus.

However, as first shown by Kagan and Afanas'ev[1,2], formula (1) is not valid for a regular system of resonant nuclei (an ideal single crystal). The point is that the derivation of formula (1) is based on the assumption that the nuclear excitation has an individual character (the collective action of the resonant nuclei becomes manifest only in the decrease of the flux of incident quanta over the thickness of the absorber). In fact, by multiple resonant scattering, the excitation of any nucleus can be extended to other nuclei, and consequently the incident quanta should give rise to a collective excited state of the entire system of nuclei as a whole.

An important role is played in the production and decay (in the radiative channel) of this collective excited state by the coherence effect, which is manifest in the interference between the incident wave and the wave multiply scattered by the resonant nuclei. For these directions of the incident quanta, which do not satisfy the Bragg condition, the coherence effect hinders the process of the radiative decay, i.e., it leads to a decrease in the width of the collective excited state compared with the width of the isolated nucleus in formula (1)[2]. When the Bragg condition is satisfied, there is observed the so-called "suppression of the inelastic channel"[1-3]. In this case the amplitude of the production of the excited nuclei in the sites of the single crystal may tend to zero, and consequently the coefficient of resonant absorption decreases sharply compared with the result calculated in accordance with formula (1).

We show in this paper that a collective excited state should be produced also in polycrystals, and demonstrate by the same token that the formula (1) used heretofore is incorrect in the general case. In polycrystals, the coherence effect in the formation and decay of the collective excited state leads to a change in the parameters Γ and E_0 of formula (1), but retains its Lorentz shape. In the cases of practical interest, when $\lambda \lesssim a$, such changes are quite noticeable. The formulas obtained by us tend to (1) only in the limit when $\lambda \ll a$ and at high temperatures. The presence of a nonzero spin of the ground state of the nucleus also decreases the coherence effect. Allowance of this nonzero spin complicates the exposition, but owing to the spin-coherence phenomenon the final results turn out to differ from the case of zero spin only by simple factors.

EQUATIONS FOR THE AMPLITUDES OF THE STATES

We shall calculate the effective cross section of resonant absorption of γ quanta in crystals by the method of nonstationary perturbation theory, in analogy with the method developed in[2] (see also[4]). Let us assume that the considered nuclear levels are completely degenerate and that the γ transition between them has a pure multipolarity L. Let $C_{\mathbf{k}\sigma\langle j\rangle n}$ be the amplitude of the state at which all the nuclei are in the ground state with spin projections $j_1, j_2, \ldots, j_N = \langle j \rangle$, the lattice is in the phonon state n, with energy ΔE_n, and there exists a photon with wave vector $\mathbf{k}$, energy ΔE_n, and polarization σ ($\sigma = \pm 1$). We shall also denote by $C_{s\langle j i_s\rangle n}$ the amplitude of the state corresponding to excitation of only a nucleus at the site s with spin projection i_s, while the remaining nuclei are in the ground state with spin projections $j_1, j_2, \ldots, j_{s-1}, j_{s+1}, \ldots, j_N$ and there is no photon. The total aggregate of the spin projections of this state is denoted by $\langle j i_s \rangle$.

The equations of motion for these amplitudes of states can be written in the energy representation in the form

$$(E-E_k-\Delta E_n)C_a(E)=\sum_{si_s n'}H_{b'}{}^{a}C_{b'}(E), \tag{2a}$$

$$(E-E_0'-\Delta E_n+{}^1/{}_2 i\Gamma_C)C_b(E)=\sum_{\mathbf{k}\sigma j_s'n'} H_c{}^b C_c(E), \tag{2b}$$

where for brevity a, b, b′, c denotes the aggregates of the indices

$$a=\{\mathrm{k}\sigma\langle j\rangle n\},\quad b=\{s\langle ji_s\rangle n\},\quad b'=\{s\langle ji_s\rangle n'\},\quad c=\{\mathrm{k}\sigma\langle jj_s'\rangle n'\},$$

Γ_C is the conversion width of the excited state of the nucleus ($\Gamma_R+\Gamma_C=\Gamma$) and $\langle jj_S'\rangle$ is the aggregate of the projections of the spins $j_1, j_2, \ldots, j_{S-1}, j_S', j_{S+1}, \ldots, j_N$. The matrix element of the γ transition of the nucleus s in the crystal is diagonal in the spin projections j_S' of the nuclei $s' \neq s$. Its general form is

$$H_{sin'}^{\mathbf{k}\sigma jn}=(H_{\mathbf{k}\sigma jn}^{sin'})^*=M_i^{\mathbf{k}\sigma j}e^{-i\mathbf{k}\mathbf{r}_s}(e^{-i\mathbf{k}\mathbf{u}_s})_{nn'}; \tag{3}$$

here $\mathbf{r}_S$ are the coordinates of the s-th lattice site, $\mathbf{u}_S$ is the displacement of the s-th nucleus due to its thermal motion, and

$$M_i^{\mathbf{k}\sigma j}=i\pi\sqrt{2L+1}\,\varepsilon C(JLI,ji-j)D_{i-j,\sigma}^L(\mathbf{k})M(k), \tag{4}$$

where M(k) is the so-called reduced matrix element, which does not depend on the direction of the wave vector **k** and the projections of the spins i and j; $\epsilon=1$ for radiation of the magnetic type and $\epsilon=i\sigma$ for radiation of the electric type; the remaining notation is standard. The matrix elements (3) are connected with the parameters of the excited state of the nucleus by the relation

$$\sum_{\mathbf{k}\sigma jn''}\frac{H_{\mathbf{k}\sigma jn''}^{sin}H_{si'n'}^{\mathbf{k}\sigma jn''}}{E_0-E_k-\Delta E_{n''}+i\delta}=\left(-i\frac{\Gamma_R}{2}+\Delta E_R\right)\delta_{nn'}\delta_{ii'}, \tag{5}$$

where δ is an infinitesimally small real number and

$$\Delta E_R=E_0-E_0' \tag{6}$$

denotes renormalization of the energy of the excited state, connected with the presence of the self-field of the γ radiation of the nucleus.

SOLUTION OF THE SYSTEM (2)

Let us consider the problem of the passage of a photon with a definite wave vector $\mathbf{k}_0$ and polarization σ through a crystal. We choose as the ground state of the crystal the state described by the aggregate of occupation numbers of the phonons n_0 ($\Delta E_{n_0}=0$) and the aggregate of the spin projections $\langle j\rangle_0$:

$$\langle j\rangle_0=j_{1_0}, \quad j_{2_0}, \ldots, j_{s_0}, \ldots, j_{N_0}.$$

Let us determine the equation for the amplitude of the state $C_{a_0}(E)$, where $a_0=\{\mathbf{k}_0\sigma_0\langle j\rangle_0 n_0\}$.

To this end we express the amplitudes of the states $C_b(E)$ in terms of $C_{a_0}(E)$ and substitute them in Eq. (2a) at $a=a_0$. First, separating from the right side of (2b) the amplitude of the state $C_{a_0}(E)$ and using relation (5), we arrive at the equation

$$\left(E-E_0-\Delta E_n+i\frac{\Gamma}{2}\right)C_b(E)=H_{a_0}{}^b C_{a_0}(E)$$
$$+\sum_{\substack{\mathbf{k}\neq\mathbf{k}_0\\ \sigma j'_s i_{s'}}}\sum_{s'\neq s}\frac{H_c{}^b H_d{}^c}{E-E_k-\Delta E_{n'}+i\delta}C_d(E), \tag{7}$$

where $d=\{s'\langle jj_S' i_S'\rangle n''\}$ (the remaining indices were explained earlier).

Equation (7) demonstrates clearly the possibility of the propagation of nuclear excitation and formation of collective excited states in a system of resonant nuclei. It is easy to verify that in the approximation of the individual nuclear excitation it is necessary to retain for the subsequent analysis only the first term in the right side of (7). In the general case, on the other hand, we can regard the second term of the right side of (7) as a perturbation, and express it by the iteration method in terms of $C_{a_0}(E)$. Substitution of the equation obtained in this manner for $C_b(E)$ in (2), with allowance for the indicated property of the diagonality of the matrix elements H^S in terms of the spin indices $j_S'(s'=s)$ yields an equation describing the motion of the photon in the medium. In real cases $\Gamma \ll \Delta E_n \ll E_0$, this equation takes the form

$$(E-E_{k_0}-g_0)C_{a_0}(E)=0, \tag{8}$$

where

$$g_0=\sum_{\substack{ss'\\ i_s i_{s'}}}H_\beta{}^\alpha\left[\frac{\delta_{ss'}}{E-E_0+i\Gamma/2}+\frac{L(i_s i_{s'})}{(E-E_0+i\Gamma/2)^2}\right. \tag{9}$$

and

$$\left.+\sum_{s''i_{s''}}\frac{L(i_s i_{s''})L(i_{s''}i_{s'})}{(E-E_0+i\Gamma/2)^3}+\ldots\right]H_{\alpha'}{}^{\beta'},$$

$$L(i_s i_{s'})=\sum_{\mathbf{k}\neq\mathbf{k}_0,\sigma,n}\frac{H_\gamma{}^\beta H_{\beta'}{}^{\gamma'}}{E-E_k+i\delta} \tag{10}$$

and we put

$$\alpha=\{\mathrm{k}_0\sigma_0 j_{s_0}n_0\},\quad \alpha'=\{\mathrm{k}_0\sigma_0 j_{s_0'}n_0\},$$
$$\beta=\{i_s n_0\},\quad \beta'=\{i_{s'}n_0\},\quad \gamma=\{\mathrm{k}\sigma j_{s_0}n\},\quad \gamma'=\{\mathrm{k}\sigma j_{s_0'}n\}.$$

We note that a contribution to the expression for g_0 is made only by terms corresponding to multiple scattering without change of the spin projections of the ground states of each nucleus. This is a manifestation of the well-known spin-coherence phenomenon. In addition, inasmuch as all the spin projections of the individual nuclei are equally probable for a system of nonoriented nuclei, it is easy to see that

$$g_0\equiv\bar{g}_0=\frac{1}{(2J+1)^N}\sum_{j_{1_0},\ldots,j_{N_0}}g_0. \tag{11}$$

Since Eq. (8) is general for all specified aggregates $\langle j_0\rangle$, we can derive from it a dispersion equation for a photon in a medium

$$E/c-k_0-g_0/c=0. \tag{12}$$

Using (3), (4), and (11), we can write for g_0 the expression

$$g_0=\frac{N_0 e^{-Z(\mathbf{k}_0)}}{2J+1}\sum_{jij'i'}M_i{}^{\mathbf{k}_0\sigma_0 j}\left[\frac{\delta_{ii'}\delta_{jj'}}{E-E_0+i\Gamma/2}+\frac{1}{2J+1}\frac{L(ijj'i')}{(E-E_0+i\Gamma/2)^2}\right.$$
$$\left.+\frac{1}{(2J+1)^2}\sum_{j''i''}\frac{L(ijj''i'')L(i''j''j'i')}{(E-E_0+i\Gamma/2)^3}+\ldots\right]M_{\mathbf{k}_0\sigma_0 j'}^{i'}, \tag{13}$$

where N_0 is the number of resonant nuclei per unit volume of the absorber

$$L(ijj'i')=\sum_{\mathbf{k}\neq\mathbf{k}_0,\sigma}\frac{M_{\mathbf{k}\sigma j}^{i}M_{i'}^{\mathbf{k}\sigma j'}}{E-E_k+i\delta}.$$
$$\times\sum_{\mathbf{r}_s\neq 0}\exp\{-i(\mathbf{k}_0-\mathbf{k})\mathbf{r}_s\}(\exp\{i\mathbf{k}(\mathbf{u}_s-\mathbf{u}_0)\})_{n_0n_0}, \tag{14}$$

$e^{-Z(\mathrm{k}_0)}$ is the probability of resonant absorption of γ quanta without a change of the phonon state of the crystal[2]

$$Z(\mathbf{k})=\frac{1}{2MN_0}\sum_{\mathbf{q}\mu}\frac{|\mathbf{ke}(\mathbf{q}\mu)|^2}{\omega(\mathbf{q}\mu)}[2\bar{n}(\mathbf{q}\mu)+1], \tag{15}$$

q and μ are the wave vector and the number of the phonon branch, $e(q\mu)$ and $\omega(q\mu)$ are the polarization and frequency of the phonon, $\bar{n}$ is the average phonon occupation number, and M is the mass of the atoms in the lattice.

Formulas (7)—(14) are general for the problem of passage of γ quanta through a medium containing resonant nuclei, with the exception of the case of single crystals, when the incident quanta satisfies the Bragg condition (as is well known[1,2], in the latter case it is necessary to separate a system of equations for the amplitudes of states $C_{k_0\sigma_0}(E)$ and $C_{k_i\sigma}(E)$, where k_i are the wave vectors of the photons scattered at the Bragg angles). For single crystals, when the incident quanta do not satisfy the Bragg condition, and also for polycrystals, the series (11) becomes a geometric progression, summation of which yields the expression

$$g_0=\frac{cN_0e^{-Z(\mathbf{k}_0)}}{8\pi}\lambda^2\frac{2I+1}{2J+1}\frac{\Gamma_R}{E-E_0+i\Gamma/2+R(E)}, \tag{16}$$

where

$$R(E)=-\sum_{\mathbf{k}\neq\mathbf{k}_0}\frac{\varphi(\mathbf{k})}{E-E_k+i\delta}\sum_{\mathbf{r}_s\neq 0}\exp\{-i(\mathbf{k}_0-\mathbf{k})\mathbf{r}_s\}$$

$$\times(\exp\{i\mathbf{k}(\mathbf{u}_s-\mathbf{u}_0)\})_{n_0n_0} \tag{17}$$

and

$$\varphi(\mathbf{k})=2\pi^2\frac{2I+1}{2J+1}\sum_{\nu=0}^{L}C^2(LL2\nu,1,-1)P_{2\nu}(\cos(\widehat{\mathbf{kk}_0}))M^2(k). \tag{18}$$

DISCUSSION OF RESULTS

From the dispersion equation (12) it is easy to go over to an expression for the cross section of resonant absorption of the γ quanta in a medium with the aid of the well known relation

$$\sigma N_0=2\,\mathrm{Im}\,k_0=-2\,\mathrm{Im}\,(g_0/c). \tag{19}$$

We note first that in the individual approximation of nuclear excitation in Eqs. (7), (9), and (13) there remain only the first term, and then we have for g_0 formula (16) with $R(E) = 0$. Substitution of (16) with $R(E) = 0$ in (19) yields directly formula (1). Thus, the collective effect in nuclear excitation leads to changes in the width and position of the resonant level, amounting to

$$\Delta\Gamma=2\,\mathrm{Im}\,R(E_0),\quad \Delta E=\mathrm{Re}\,R(E_0). \tag{20}$$

Let us determine these values for polycrystalline absorbers.

We consider first the case of a rigid lattice and low temperatures. Then the phonon factor $\exp[ik(u_s - u_0)])_{n_0n_0}$ can be set equal to unity, and formulas (20) become

$$\Delta\Gamma^{(0)}=2\pi\Big[\sum_{\mathbf{k}\neq\mathbf{k}_0}\sum_{\mathbf{r}}\varphi(\mathbf{k})\exp\{-i(\mathbf{k}_0-\mathbf{k})\mathbf{r}\}-\sum_{\mathbf{k}\neq\mathbf{k}_0}\varphi(\mathbf{k})\Big]\delta(E_0-E_k). \tag{21}$$

$$\Delta E^{(0)}=\mathrm{P}\Big[\sum_{\mathbf{k}\neq\mathbf{k}_0}\sum_{\mathbf{r}}\frac{\varphi(\mathbf{k})}{E-E_k}\exp\{-i(\mathbf{k}_0-\mathbf{k})\mathbf{r}\}-\sum_{\mathbf{k}\neq\mathbf{k}_0}\frac{\varphi(\mathbf{k})}{E-E_k}\Big], \tag{22}$$

where the symbol P denotes the sums over k are taken in the sense of principal values. The summation in (21) over the lattice sites of the polycrystal can be carried out in accordance with the following scheme[5].

We first sum over the lattice sites of the polycrystal, using the relation

$$\sum_{\mathbf{r}}\exp\{-i(\mathbf{k}_0-\mathbf{k})\mathbf{r}\}=\eta\frac{(2\pi)^3}{V}\sum_{\tau}\delta(\mathbf{k}_0-\mathbf{k}-2\pi\boldsymbol{\tau})F_\tau, \tag{23}$$

where η is the concentration of the resonant isotopes in the absorber, V is the volume of the unit cell, $2\pi\tau$ is the reciprocal-lattice vector, and F_τ is a structure factor corresponding to the vector τ. Then, fixing the modulus of the vector τ, we average (23) over all possible directions of this vector with respect to the vector $k_0 - k$. This averaging does not affect the structure factor, and we can write

$$\frac{1}{4\pi}\int\delta(\mathbf{k}_0-\mathbf{k}-2\pi\boldsymbol{\tau})\,d\Omega_\tau=\frac{1}{2(2\pi)^3\tau^2}\delta(|\mathbf{k}_0-\mathbf{k}|-2\pi|\boldsymbol{\tau}|). \tag{24}$$

Finally, substituting (23), (24) in (21) and summing over k, we obtain

$$\Delta\Gamma^{(0)}=\eta\Gamma_R\frac{2I+1}{2J+1}\Big[\sum_{\nu=0}^{L}C^2(LL2\nu,1-1) \tag{25}$$

$$\times\sum_{\tau=\tau_{\min}}^{k_0/\pi}\frac{\pi\beta_\tau F_\tau}{2V\tau k_0^2}P_{2\nu}\Big(1-\frac{2\pi^2\tau^2}{k_0^2}\Big)-\frac{1}{2L+1}\Big],$$

where β_τ is the repeatability factor, indicating the existence of several vectors τ having the same absolute magnitude.

The quantity $\Delta\Gamma^{(0)}$ in (25) depends on the crystal structure and on the photon wavelength λ, and can assume both positive and negative values, depending on the ratio a/λ (we note that $|\Delta\Gamma^{(0)}| < \Gamma_R$). For a specified crystal structure, the quantity $\Delta\Gamma^{(0)}$ changes jumpwise with changing ratio a/λ, this being connected with the successive inclusion in the sum (25) of new reciprocal-lattice vectors with increasing a/λ. This effect, however, becomes rapidly smoothed out with increasing a/λ, and in the shortwave limit $a \gg \Gamma$ the quantity $\Delta\Gamma^{(0)}$ should tend to zero. For sufficiently large ratios a/λ, the function in the sum over τ in (25) can be regarded as continuous in τ, and the sum over τ can then be replaced by an integral

$$\sum_{\tau}\beta_\tau\to 4\pi V\int\tau^2\,d\tau.$$

Recognizing at the same time that $F_\tau \to 1$ and that the sum (24) contains no term with $\tau = 0$, we get from (25) the value of $\Delta\Gamma^{(0)}$ for the case of sufficiently short wavelengths:

$$\Delta\Gamma^{(0)}=\eta\Gamma_R\frac{3\lambda^2}{8\pi a^2}\frac{2I+1}{2J+1}. \tag{26}$$

Actually formula (26) can be employed with high accuracy already at $a/\lambda \gtrsim 10$.

Let us proceed to calculate the shift of the resonant energy $\Delta E^{(0)}$ in accordance with formula (22). For a polycrystal, expression (22) takes the form

$$\Delta E^{(0)}=\frac{\eta}{c}\mathrm{P}\Big[\frac{(2\pi)^3}{V}\sum_{\mathbf{k}\neq\mathbf{k}_0}\sum_{\tau}\overline{\frac{\varphi(\mathbf{k})}{k_0-k}\delta(\mathbf{k}_0-\mathbf{k}-2\pi\boldsymbol{\tau})}-\sum_{\mathbf{k}\neq\mathbf{k}_0}\frac{\varphi(\mathbf{k})}{k_0-k}\Big], \tag{27}$$

where the bar denotes averaging over all possible orientations of the single crystals (averaging over the directions of the vector τ or the vector k_0). Integration of (27) with respect to k with allowance *for (5)* yields

$$\Delta E^{(0)}=\eta\Big[\frac{1}{cV}\sum_{\tau\neq 0}\frac{\varphi(\mathbf{k}_0-2\pi\boldsymbol{\tau})}{k_0-|\mathbf{k}_0-2\pi\boldsymbol{\tau}|}-\frac{2I+1}{(2J+1)(2L+1)}\Delta E_R\Big]. \tag{28}$$

At sufficiently large ratios a/λ, the sum over τ can be replaced by an integral over the reciprocal-lattice space. It is then easy to see that the first term in the right side of (27) is exactly equal to the second and $\Delta E^{(0)} = 0$. Thus, in the short-wave limit both $\Delta\Gamma^{(0)}$ and $\Delta E^{(0)}$ tend to zero, and the usual Breit-Wigner formula (1) is valid.

Formulas (25)—(28) were obtained under the assumption that the lattice is rigid and the temperature is low ($T = 0$). In the opposite limiting case, the phonon multiplier $(\exp[i\mathbf{k}(\mathbf{u}_s - \mathbf{u}_0)])_{n_0 n_0}$ tends to zero, and with it also the quantities $\Delta\Gamma$ and ΔE.

Let us consider now the general case ($T \neq 0$). The phonon factor in (17) can be represented in the form[2,6]

$$(\exp i\mathbf{k}(\mathbf{u}_s - \mathbf{u}_0))_{n_0 n_0} = \exp\{-Z(\mathbf{k}) + Y_s(\mathbf{k})\}, \tag{29}$$

where

$$Y_s(\mathbf{k}) = \frac{1}{2MN_0} \sum_{\mathbf{q}\mu} \frac{|\mathbf{k}\mathbf{e}(\mathbf{q}\mu)|^2}{\omega(\mathbf{q}\mu)} [2\bar{n}(\mathbf{q}\mu) + 1] \cos(\widehat{\mathbf{q}\mathbf{r}_s}). \tag{30}$$

The exponential in (29) can be expanded in powers of $Y_S(\mathbf{k})$. At not too high temperatures, it suffices to consider only the phononless and single-phonon processes of transfer of nuclear excitation, which correspond to the first two terms of this expansion. Then

$$R(E) = -\sum_{\mathbf{k}\neq\mathbf{k}_0} \frac{\varphi(\mathbf{k})}{E - E_k + i\delta} \sum_{\mathbf{r}_s\neq 0} e^{-i(\mathbf{k}_0-\mathbf{k})\mathbf{r}_s} e^{-Z(\mathbf{k})}(1 + Y_s(\mathbf{k})). \tag{31}$$

It can be shown that single-phonon processes make a small contribution to Re $R(E_0)$. We therefore have for the shift of the resonant energy at $T \neq 0$, in analogy with (27),

$$\Delta E = \frac{\eta}{c} \mathrm{P} \left[\frac{(2\pi)^3}{V} \sum_{\mathbf{k}\neq\mathbf{k}_0} \sum_{\tau} \frac{\varphi(\mathbf{k})}{k_0 - k} e^{-Z(\mathbf{k})} \delta(\mathbf{k}_0 - \mathbf{k} - 2\pi\boldsymbol{\tau}) - \sum_{\mathbf{k}\neq\mathbf{k}_0} \frac{\varphi(\mathbf{k})}{k_0 - k} e^{-Z(\mathbf{k})}\right]. \tag{32}$$

To calculate the change of the level width, we can neglect the anisotropy of the crystal, i.e., we can put $e^{-Z(\mathbf{k})} = e^{-Z}$, where e^{-Z} does not depend on the direction of the wave vector $\mathbf{k}$ and is also the probability of phononless resonant absorption of γ quanta in the polycrystal. It then follows from (Fig. 1) that

$$\Delta\Gamma = e^{-Z}\Delta\Gamma^{(0)} + \Delta\Gamma^{(1)}, \tag{33}$$

where $\Delta\Gamma^{(0)}$ is determined by formulas (25) and (26), while $\Delta\Gamma^{(1)}$ is the contribution of the single-phonon processes and is given by

$$\Delta\Gamma^{(1)} = \eta\Gamma_R \frac{2I+1}{2J+1} e^{-Z} \left[\sum_{\nu=0}^{L} C^2(LL2\nu, 1-1) \times \sum_{\tau\mathbf{q}\mu} \frac{\pi^2 F_\tau \{2\bar{n}(\mathbf{q}\mu) + 1\}}{6M\omega(\mathbf{q}\mu)\,|2\pi\boldsymbol{\tau} + \mathbf{q}|} P_{2\nu}\left(1 - \frac{|2\pi\boldsymbol{\tau} + \mathbf{q}|^2}{2k_0^2}\right) - \frac{Z}{2L+1}\right]. \tag{34}$$

$\Delta\Gamma^{(1)}$ tends to zero when $\lambda/a \ll 1$.

In conclusion, it is important to note that the coherence effect in collective nuclear excitation can be revealed both by the change of the width and shift of the resonant line, and by the intensity of the resonant absorption. The coefficient of resonant absorption of γ quanta in polycrystals differs by a factor $\Gamma/(\Gamma + \Delta\Gamma)$ from the value used for it so far in the approximation of individual nuclear excitation.

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