$b \sim 1$ and $G \sim 10^{-15}$ erg/deformation unit. Such a spin-phonon interaction mechanism was observed in quartz [2-4], and the dependence of α_0 on θ obtained by us recalls the dependence of α_0 on θ for paramagnetic centers in quartz.

Comparison of G for Cr^{3+} in Al_2O_3 and $LiNbO_3$ shows that the presence of electric domains strongly affects the character of the spin-phonon interaction.

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RADIATION OF AN ATOM (MOLECULE) IN AN ABSORBING MEDIUM

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1. We shall estimate the possible role played by one simple effect when a source emits in an absorbing medium (an effect to which apparently no attention has been paid before) - the absorption of radiation and its conversion into heat in the non-wave zone of the emitter. This is essentially a collective interaction between the emitter and a large number of particles that fall into this zone, when the dipole oscillations are inevitably accompanied by quasistationary currents. This effect takes place if $[(4\pi/3)(N/k^3)] \gg 1$, where $\lambda = \lambda_0/\sqrt{|\epsilon|} = 2\pi/(k_0\sqrt{|\epsilon|}) = 2\pi/|k|$ is the wavelength in the medium, $\epsilon = \epsilon' + i\epsilon''$ is the complex dielectric constant, N is the number of absorbing particles per unit volume. The effect has a classical nature and is estimated classically.

The energy absorbed per unit time, dW/dt, referred to the initial emitter energy $\hbar\omega = \hbar c k_0$, gives the following line width (without allowance for the microscopic structure of the medium, i.e., without the impact and Stark widths, and also without the Doppler width):

$$\gamma = \frac{1}{\hbar\omega}\frac{dW}{dt} = \frac{1}{\hbar\omega}\int \sigma |E|^2 dV = \frac{1}{\hbar\omega}\frac{\omega\epsilon''(\infty)}{4\pi}\int_{(V_0)} |E|^2 dV \qquad (1)$$

(the notation is standard), where the integration over the volume must be carried out outside a volume V_0 of radius R_0 such that on the average the volume V_0 contains one particle of the medium (the binary broadening theory is valid inside the volume):

$$\frac{4}{3}\pi R_0^3 N = 1. \qquad (2)$$

Since $|\vec{E}| \sim R^{-3}$, the integral depends strongly on R_0 when $P < \lambda$ and therefore we cannot obtain an exact estimate.

2. We substitute in (2) the complete expression for the field of the dipole moment, an expression valid in the non-wave zone. We put, further, $\xi = k_2/|k| = \mathrm{Im}\sqrt{\epsilon}/\sqrt{|\epsilon|}$ (ξ measures the damping on the path λ), $\epsilon'' = |\epsilon| 2\xi\sqrt{1-\xi^2}$. An interesting case, of course, is when $\xi \ll 1$. Therefore, $\epsilon'' \simeq 2\xi\epsilon$. We get:

$$\frac{\gamma}{\gamma_0} = |\epsilon| \exp(-2\xi|k|R_0)\left(1 + \frac{2\xi}{|k|R_0} + \frac{4\xi^2}{|k|^2 R_0^2} + \frac{2\xi}{|k|^3 R_0^3}\right). \tag{3}$$

Here $\gamma_0 = (2/3)\omega p^2/\hbar c$ is the value that γ retains in vacuum, $k_2 \to 0$, $k \to k_0$, and $\xi \to 0$, i.e., natural width. We see that the transition to vacuum is realized when $2\xi \ll |k|^3 R_0^3$ or if $|\epsilon| \sim 1$, when $\epsilon'' \ll k_0^3 R_0^3$. On the other hand, if $\epsilon''/|k|^3 R_0^3 \gtrsim 1$, we have

$$\frac{\gamma}{\gamma_0} \approx \frac{\gamma'}{\gamma_0} = \frac{2\xi}{k_0^3 R_0^3} \approx \frac{\epsilon''}{k_0^3 R_0^3} \approx \epsilon'' \frac{\lambda_0^3 N}{6\pi^2} \tag{4}$$

and the entire initial energy is not radiated, but is converted into heat already in the non-wave zone (of course, a more detailed analysis should reveal also a frequency shift). Here $\gamma' = \gamma - \gamma_0$ is the width dure to this absorption. It measures the effect of interest to us, and γ' is small compared with the impact width etc. so long as $\gamma'/\gamma_0 < 1$. In essence, (4) and the condition $\gamma'/\gamma_0 \sim 1$ determine the threshold beyond which emission of radiation by an atom (molecule) becomes impossible.

3. Let us apply this result to an emitter in a quasineutral plasma. The number of electrons N, λ_0, and the temperature T will be expressed by means of the quantities N^*, λ_0^*, and T^*, in fact measuring T in keV, λ_0 in microns, and N in terms of pressure in atmospheres at normal temperature:

$$N = 2.7 \times 10^{19} N^*, \quad \lambda_0 = 10^{-4}\lambda_0^*, \quad T = 10^4 T^*. \tag{5}$$

Neglecting collisions between electrons and neutral atoms, we obtain for the effective electron collision frequency

$$\nu_{eff} = 1{,}5\,\frac{N^*}{T^{*3/2}}\,L \cdot 10^{14}\,\mathrm{sec}^{-1}, \quad L = \ln\frac{0{,}73\,T^*}{N^{*1/3}}. \tag{6}$$

The main case is when $\omega^2 \gg \nu_{eff}^2$, i.e., $N^*\lambda_0^{*2}/T^{*3/2} \ll (2.4/L) \times 10^{16}$. Then

$$\frac{\gamma'}{\gamma_0} \approx \frac{\epsilon''}{k_0^3 R_0^3} = 78\,L\,\frac{\lambda_0^{*6} N^{*3}}{T^{*3/2}} \tag{7}$$

(when $\omega^2 \ll \nu_{eff}^2$ we obtain $\gamma'/\gamma_0 = 4.3 \times 10^{-6} L^{-1} \lambda_0^{*4} N^* T^{*3/2}$).

Thus, the effect depends very strongly on λ_0 and N. When $N^* \sim 10^{-1}$ and $T^* \sim 1$ in the infrared region ($\lambda_0^* \gtrsim 1$), γ'/γ_0 can be very large. Such radiation will all be absorbed in the non-wave zone. At a lower temperature and somewhat larger values of N^*, the atom will not emit even optical lines.

4. However, we still did not take into account here the absorption, in the same zone, by meutral and ionized atoms and molecules (in particular, of the same kind as the emitter). In general the effect can take place in any medium with $\epsilon'' \neq 0$. It is therefore more convenient to express γ' more phenomenologically - in terms of the range of the photon absorption ℓ in the medium, $\exp(-\ell \text{ Im } k) \sim 1$, i.e., $\ell = (2|k|\xi)^{-1} \sim (|k|\epsilon'')^{-1}$. Equation (4) yields

$$\frac{\gamma'}{\gamma_0} = 2{,}7\,\frac{N^* \lambda_0^{*4}}{\ell}, \tag{8}$$

where ℓ is in centimeters. With such an approach N is no longer the number of electrons, but a parameter which determines, in accord with (2), the minimum distance, starting with which the collective quasistationary interaction is realized between the emitter and the absorbing particles (resonant atoms which lose in their excited state energy to friction and to impacts of the second kind, and also electrons producing the ohmic losses).

The conditions for the realization of the effect are worse in radioastronomic conditions, where N is relatively small. However, it may be encountered for rotational frequencies, and especially for vibrational frequencies, as well as for transitions between highly-excited levels.

It must be emphasized once more that owing to the strong dependence on the not-too-well-defined quantity R_0, the presented numerical estimate is quite rough. However, the effect itself is real. Thus, assuming, for the sake of reliability, a value of R_0 larger than (2) by 2 - 3 times, we obtain, to be sure, a smaller value of γ', but the correctness of the calculation becomes obvious. As noted by V. L. Ginzburg during the discussions, in the case of very small R_0 an important role may be assumed by spatial dispersion, which possibly automatically cuts off the integral, and the introduction of the parameter R_0 is unnecessary.

Indeed, if the Debye radius D turns out to be larger than R_0, then it can be assumed that one should choose as the lower limit in (1) not R_0 but D. In terms of our variables, $D/R_0 \sim 0.4 T^{*1/2} N^{*-1/6}$. In such an approach, if $D/R_0 > 1$, we get in lieu of (7) $\gamma'/\gamma_0 \sim 90 \cdot L\lambda_0^{*6} N^{*7/2}/T^{*3}$. Of course, a more detailed analysis is required here.

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SINGULARITIES OF THE PHONON SPECTRUM OF CRYSTALS WITH EXTENDED DEFECTS

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The purpose of the present note is to call attention to several very general peculiarities of the spectral functions of crystals with linear and two-dimensional defects (edge dislocations, stacking faults, flat boundaries, and others). If the defect has the form of a straight line or a plane (assuming that the defect concentration is low and they can be re-