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POSSIBILITY OF CONTROLLING THE PARAMETERS OF PHONONLESS LINES WITH THE AID OF ULTRASOUND

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In the case when electrons localized on impurities interact strongly with lattice vibrations, multiphonon processes leading to broad electron-vibrational light-absorption and luminescence bands become probable [1]. If the constant of the coupling with the oscillations is small, then at low temperatures the spectrum constitutes a phononless line with a weakly pronounced vibrational structure. A rise in temperature can lead to a strong release of heat even in the case of weak coupling, owing to stimulated processes in the phonon subsystem.

In the present paper we propose a method for directed variation of the shape of the optical curves by heating the phonon subsystem of the crystal in a narrow spectral region; this heating is produced, for example, by a powerful ultrasonic wave or by stimulated Mandel'shtam-Brillouin scattering of laser radiation. It turns out that this superheating can become strong enough to make it possible to observe ultrasound-stimulated multiphonon absorption of light.

We write the light-absorption coefficient in the form [2]:

$$K(\Omega) = \frac{1}{2\pi}\int_{-\infty}^{\infty} e^{-i\Omega t - \Gamma|t|} \langle\langle d^{+}(0)d(t)\rangle\rangle\, dt,$$

where d(t) is the operator of the dipole moment in the Heisenberg representation, and the operator of the interaction with the ultrasound is chosen in an approximation that is linear in the phonon operators $b_{\kappa\sigma}$ and $b^{+}_{\kappa\sigma}$,

$$\hat{H}' = \frac{1}{\sqrt{2}}\sum_{i} v_{ii\vec{\kappa}\sigma}\,(b^{+}_{\vec{\kappa}\sigma} + b_{\vec{\kappa}\sigma})$$

with account taken of only the polaron effect; i numbers of electronic states, which are assumed to be nondegenerate; $\vec{\kappa}$ and σ are the wave vector and the polarization index of the ultrasonic wave of frequency ω; Γ is the damping constant of the discrete lines of the optical spectrum. The results of the averaging in (1) depends on the statistical properties of the vibrations introduced in the crystals, and we consider here two limiting cases: 1) an absolutely coherent source and 2) a thermal (Gaussian) source [3]. The calculation procedure is the same as in the problem of calculating K(Ω) in the presence of electromagnetic radiation of high intensity [3].

For an absolutely coherent ultrasound source we obtain:

$$K(x) = \frac{1}{2\pi}\int_{-\infty}^{\infty} dt \exp(ixt - \Gamma|t|) I_0\left(\frac{4B_{K\sigma}\sqrt{\bar{n}}}{\hbar\omega}\sin\frac{\omega t}{2}\right), \tag{2}$$

where $\bar{n}\hbar\omega$ is the volume energy of the ultrasound and I_0 is the Bessel function:

$$B_{\vec{K}\sigma} = v_{ii\vec{K}\sigma} - v_{kk\vec{K}\sigma}, \quad \hbar x = \hbar\Omega - \epsilon_i + \epsilon_k, \quad |d_{ik}|^2 = 1,$$

$\varepsilon_{i(k)}$ is the energy of the electronic terms. Using the well-known expansion

$$I_0\left(z\sin\frac{\omega t}{2}\right) = \sum_{p=-\infty}^{\infty} I_p^2(z/2)e^{ip\omega t},$$

we can represent (2) in the form

$$K(x) = \frac{1}{2\pi}\sum_{p=-\infty}^{\infty} I_p^2\left(\frac{2B_{\vec{K}\sigma}\sqrt{\bar{n}}}{\hbar\omega}\right)\frac{2\Gamma}{(x+p\omega)^2+\Gamma^2}, \tag{3}$$

which shows clearly the multiquantum nature of the band, in which the fine structure can be resolved at sufficiently small natural widths Γ, and the "effect of suppression" of the optical absorption [3] can be observed. For sufficiently powerful ultrasound sources $(B\bar{n}^{1/2}/\hbar\omega) >> 1$ (estimates are given below) we can use the quasiclassical approximation, replacing $\sin(\omega t/2)$ by $\omega t/2$, and obtain

$$K(x) = \frac{\left[\sqrt{\left(\Gamma^2 - x^2 + \frac{4B^2\bar{n}}{\hbar^2}\right)^2 + 4x^2\Gamma^2} + \Gamma^2 - x^2 + \frac{4B^2\bar{n}}{\hbar^2}\right]^{1/2}}{\sqrt{\left(\Gamma^2 - x^2 + \frac{4B^2\bar{n}}{\hbar^2}\right)^2 + 4x^2\Gamma^2}}. \tag{4}$$

Expression (4) describes a symmetrical spectral curve, which is bell-shaped if $\Gamma^2 > B^2\bar{n}/\hbar^2$ (but not Gaussian), and is double-humped in the opposite case. Curve (4) represents a multiphonon band stimulated by ultrasonic phonons whose frequencies are $\omega << \Gamma, \omega_D$. In this case the structure of the band cannot be resolved in practice, and the effect can become manifest experimentally in the form of stimulated ultrasonic broadening and a change in the contour of the phononless line; formula (1) therefore takes into account the effect of the natural width Γ of the phononless line. It should be borne in mind that Γ, generally speaking, depends also on the ultrasound power.

In the case of a Gaussian (thermal) source [3] we obtain

$$K(x) = \frac{1}{2B\sqrt{\pi\bar{n}}}\operatorname{Re}\left\{\exp\left[\frac{\hbar^2(\Gamma+ix)^2}{4B^2\bar{n}}\right]\left[1 - \Phi\left(\frac{\hbar(\Gamma+ix)}{2B\sqrt{\bar{n}}}\right)\right]\right\}, \tag{5}$$

where Φ is the error function. The curve (5) always has a bell-shaped form, with a maximum at $\Omega = \Omega_0$ (the absence of a Stokes shift in the case of strong heat release should not be surprising, since the constant of coupling with long-wave ultrasonic oscillations is very small, and terms of the type $B << B\bar{n}^{1/2}$ have been discarded). Comparison of curves (3) and (5) can serve as the basis for an experimental establishment of the statistical properties of ultrasonic phonons in crystals.

A favorable situation for the observation of the indicated effects is one in which the electron-phonon interaction in the impurity center is not too small, but at the same time it is not large enough to hide the phononless line by the vibrational structure. Such objects, for example, are the phononless line of the U-band of ruby ($Cr^{3+}:Al_2O_3$) and also the phononless line and its vibrational replicas of the $4f \rightarrow 5d$ transitions in $Ce^{3+}:CaF_2$ [4]. For numerical estimates we shall use the procedure for calculating the parameters of the electron-phonon interaction [5, 6]; this procedure results in good quantitative agreement with experiment. For the phononless line of the U band (the transition ${}^4A_{2g}(t_2^3) \rightarrow {}^4T_{2g}(t_2^3)$), we obtain the broadening δ:

$$\delta = (100\sqrt{2\ln 2}/3)D_q\sqrt{P/\rho v^3}, \qquad (6)$$

where D_q is the parameter of the theory of the crystalline field, ρ is the density of the crystal, D is the density of the ultrasonic energy flux. Using for ruby D_q = 1800 cm^{-1}, ρ = 4g/cm^3, v = 10^6 cm/sec, and P = 100 W/cm^2 (this power corresponds to a frequency $\nu \sim 10^6$ sec^{-1}), we get $\delta \sim 1$ cm^{-1}. For the highest attainable ultrasonic powers ($P \simeq 10^5$ W/cm^2, $\nu = 2 \times 10^4$ sec^{-1} [7]) we obtain $\delta \simeq 30$ cm^{-1}.

The proposed effect can be important for applications, since the width of the phononless line is one of the main parameters determining the operating conditions of a laser.

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MECHANISM OF EXCITATION AND CONTROL OF RELAXATION OSCILLATIONS IN A PLASMA-BEAM SYSTEM

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Relaxation oscillations in a beam-plasma system placed in a magnetic field (these were observed by many investigators and described in detail in [1]) are customarily called low-frequency (LF) oscillations, resulting from the rapid diffusion of the plasma from the region where the beam is located. The diffusion of the plasma is due to ionic oscillations. When the diffusion causes the plasma density in the region of the beam to become lower than a critical value, the interaction between the beam and the plasma stops, and the power of the high-frequency (HF) electronic oscillations excited by the beam decreases, and with them also the power of the ionic oscillations; the diffusion stops, the