POSSIBILITY OF GENERATING INFRARED RADIATION BY POLARIZATION REVERSAL OF A FERROELECTRIC

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The possibility of generating electromagnetic radiation in the microwave band via transitions of ferromagnets from the metastable to the stable state in a magnetic field was demonstrated recently [1].

We have evaluated a similar possibility of generating infrared radiation by uniform polarization reversal of a ferroelectric by an electric field. We consider a ferroelectric with a second-order phase transition at $T < T_c$ (T_c is the transition temperature). In an electric field, the ferroelectric can be in two states, one stable and the other metastable. The system is in the latter state if the external field $\vec{E}$ is directed opposite to the spontaneous polarization $\vec{P}$ [2].

From the known equation for the free energy of the ferroelectric

$$F = F_0 + \frac{\alpha}{2}P^2 + \frac{\beta}{4}P^4 - EP \tag{1}$$

we can readily determine the thermodynamic coercive field E_c at which the metastable state vanishes and the ferroelectric is repolarized and goes over into the stable state. Following this process, the field and the polarization have the same direction

$$E_c = \frac{2|\alpha|}{3}\sqrt{-\frac{\alpha}{3\beta}}. \tag{2}$$

In the field E_c, the equilibrium polarization values P_{01} and P_{02} corresponding to the metastable and stable states are

$$P_{01} = \sqrt{-\alpha/3\beta}; \quad P_{02} = 2\sqrt{-\alpha/3\beta}. \tag{3}$$

The transition from the vanished metastable state into the stable one should be accompanied by a transient process in which the polarization oscillates about the new equilibrium value $P = P_{02}$. The damped oscillations of P should produce a pulse of infrared electromagnetic radiation from the sample.

The duration of such a pulse is determined by the oscillation damping time γ^{-1}, which should coincide with the time of the experimentally-determined damping of the "soft" ferroelectric phonon mode. The pulse is repeated by reversing the sign of the external field.

Actually, however, the polarization reversal is not uniform through the sample, and occurs in fields much weaker than E_c. The reason is that the repolarization process proceeds at the expense of a growth of domains having dipole moments in the opposite direction, and the rate of repolarization is limited by processes such as nucleus formation and domain-wall motion [3]. To obtain a high radiation power it is necessary to cause the crystal to reverse its polarization immediately in the entire volume. This can be done in rapidly alternating fields $E \sim E_c$ or by using dc field pulses with very steep fronts.

Since the velocity of the domain walls can apparently not exceed the speed of sound u, the time of polarization reversal due to domain motion in a crystal of thickness d is $\Delta t \gtrsim d/u$. The front rise time δ_f should therefore be shorter than Δt. This is the most stringent estimate. The character of the dipole-moment oscillations can be ascertained from the time-dependent equation of motion for $\vec{P}$ [4]:

$$\ddot{\mathbf{P}} + \gamma\dot{\mathbf{P}} + \alpha\mathbf{P} + \beta\mathbf{P}^3 - \mathbf{E} = 0. \tag{4}$$

Representing P in the form $P = P_{02} + P(t)$ at $E = E_c$, we obtain for the alternating part P(t) the nonlinear-oscillator equation

$$\ddot{\mathbf{P}} + \gamma\dot{\mathbf{P}} + (\alpha + 3\beta P_{02}^2)\mathbf{P} + 3\beta P_{02}\mathbf{P}^2 + \beta\mathbf{P}^3 = 0, \tag{5}$$

with $P_{max}|_{t=0} \sim -P_{01}$.

The frequency ω of such an oscillator depends on the amplitude e of the oscillations, and if $\gamma \ll \omega$ it is given by [5]

$$\omega = \omega_0 + \Delta\omega;$$

$$\omega_0 = \sqrt{-3\alpha}; \qquad \Delta\omega = \left[\frac{3\beta}{8\omega_0^2} - \frac{5}{12}\frac{(3\beta P_{02})^2}{\omega_0^3}\right]P^2. \tag{6}$$

$\Delta\omega$ results in a small negative correction to ω_0 (on the order of 30%) even when $P \sim P_{max}$. Since the oscillator is nonlinear and is strongly excited at the initial instant, the amplitudes of the oscillations at the higher-harmonic frequencies can also be large. The coherent electromagnetic radiation power connected with the dipole-moment oscillations can be estimated for a sample with dimension $d < \lambda \sim \omega_0/c$ (c is the speed of light) by using the formula for dipole radiation

$$I = \frac{\omega_0^4}{c^3}|P_{\omega_0}|^2, \tag{7}$$

where P_{ω_0} is the dipole moment of the entire sample with volume $V \sim d^3$. The energy radiated per cycle is $W = I\gamma^{-1}$. When the dimensions of the sample greatly exceed the light wavelengths, the radiation has a different dependence on the volume and shape of the sample, but the maximum radiated energy cannot exceed the energy stored in the system at the instant of the start of the oscillations

$$W_{max} \sim P_{01}E_cV. \tag{8}$$

We present numerical estimates. In a wide range near the Curie point, the value of ω depends on the temperature and can range from 10^{13} to 10^{12} sec^{-1}. We have $\gamma \sim 10^{11}$ sec^{-1}; $P_{01} \sim 5 \times 10^4$ cgs esu, and $E_c \sim 5 \times 10^5$ V/cm. Then $W \sim 1$ erg if $d < \lambda \sim 10^{-2}$ cm, and $W_{max} \sim 10^7$ erg at $d \sim 1$ cm. Thus, if uniform polarization reversal can be accomplished in the sample, the system under consideration can serve as a source of short powerful pulses of far-infrared radiation, with a frequency that can be varied in a wide range by varying the temperature. The radiation should be completely polarized, since the oscillations are produced in a definite direction, namely the ferroelectric axis.

An appreciable fraction of W_{max} may be converted into heat. However, even if we disregard heat dissipation, the resultant temperature rise, at the characteristic values of the specific heat, amounts to several tenths of a degree.

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ROLE OF EXCITON-PHONON COMPLEXES IN THE INTERBAND MAGNETOOPTICAL ANOMALY

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Johnson and Larsen [1] observed an anomaly in the interband absorption of InSb in a strong field H when the electron cyclotron frequency ω_c approached the frequency ω of the polarized phonons. The singularity arising under these conditions in the energy spectrum of the polaron was calculated in [1, 2] and an interpretation of the observed effect was proposed on the basis of this singularity. It is indicated in [3] that this rearrangement of the spectrum corresponds to formation of electron-phonon complexes.

The interpretation of the absorption singularities at $\omega_c \simeq \omega$ on the basis of electron-phonon complexes is unconditionally correct if intraband transitions are involved (as, for example, in [4]). Since, however, transitions to the ground state of the exciton predominate in the intrinsic absorption in a strong magnetic field [5], failure to take the Coulomb interaction into account may qualitatively distort the result. We start here from the opposite approximation, i.e., we assume that the electron-phonon interaction is weak compared with the Coulomb interaction, and propose a theory of interband absorption at $\omega_c \simeq \omega$ in terms of exciton-phonon complexes.

We assume that the electron-phonon coupling is weak and that resonance $\omega_c \simeq \omega$ sets in only in the electron band; the hole-phonon interaction is therefore neglected. The magnetic field is assumed strong, i.e., $\omega_c \gg R$, where R is the exciton ionization potential at H = 0. The wave function of the system with zero total momentum in the one-phonon approximation is

$$\Psi = \Phi_{0,1,1} f_0(z)|0> + \sum_q \Phi_{-q,q,1} f_q(z) b_q^+ |0> \qquad (1)$$

where the Oz axis is parallel to H. The function Φ_{q,n_1,n_2} describes the transverse motion of the free electron and hole in the Landau bands n_1 and n_2 with total momentum $q_\perp$ and the longitudinal motion of the center of gravity with momentum q_z. The functions of longitudinal relative motion f_0 and f_q are defined by the equations ($z = z_e - z_h$):