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ELECTRIC PROPERTIES OF A SEMICONDUCTOR IN THE FIELD OF A STRONG ELECTROMAGNETIC WAVE

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1. In [1, 2] we investigated the electric and magnetic properties of a semiconductor in the field of a strong electromagnetic wave $\vec{A}(t) = \vec{A}\cos\omega t$ with a frequency ω exceeding the width Δ of the forbidden band. It was shown that the stationary state of the system is a state of saturation, wherein the strong-field absorption coefficient is equal to zero (in the absence of recombination) (see [3]). In addition, it was shown that in the electron (hole) spectrum there is a gap near the quasimomentum p_0 (see Fig. 1)

$$\lambda_p = \frac{1}{2} e A v_{cv, p_0} = \sqrt{m(\omega - \Delta)}\,, \quad \hbar = c = 1, \qquad (1)$$

where $\vec{v}_{cv}$ is the matrix element of the transition between the conduction and valence bands; the effective masses of the electrons and holes are assumed to be equal.

We consider in this paper the absorption of a weak electromagnetic field $\vec{A}_1(t) = \vec{A}\cos\omega_1 t$ with frequency $\omega_1 \sim \omega$,1) and also direct current in a semiconductor in the saturation state.

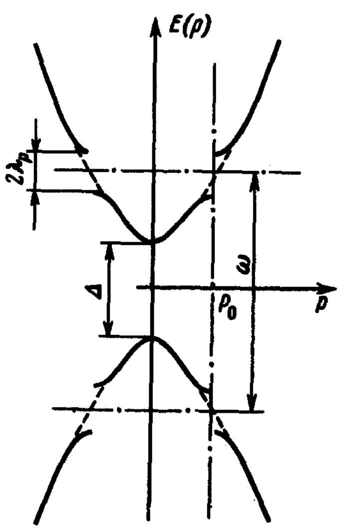

Fig. 1

2. Using standard perturbation-theory methods, we obtain the average Fourier component of the current in the first approximation in the weak field $\vec{A}_1$

$$j_i(\omega_1 q) = K_{ij}(\omega_1 q)\, A_{1j}(\omega_1 q)$$

where the imaginary part K_{ij}, describing the absorption is given at $Q \to 0$ and $T = 0$ by (see Eq. (26) of [1])

$$K_{ij}(\omega_1) = \frac{\pi e^2}{2} \sum_p v^i_{cv}(p)\, v^j_{cv}(p) \left[u_p^4\, \delta(2\epsilon_p + \omega - \omega_1) - v_{-p}^4\, \delta(2\epsilon_p - \omega + \omega_1)\right], \qquad (2)$$

where

$$u_p^2,\ v_{-p}^2 = \frac{1}{2}\left(1 \pm \frac{p^2 - p_0^2}{2m\epsilon_p}\right), \quad \epsilon_p = \sqrt{\left(\frac{p^2 - p_0^2}{2m}\right)^2 + \lambda_p^2}\,.$$

1) A similar method is used in gases, as pointed out to us by I. I. Sobel'man.

If the frequency difference is larger than the gap, $|\omega - \omega_1| \gg \lambda$, then we get from (2)

$$K_{ij} = \delta_{ij} K(\omega_1), \quad K(\omega_1) = \pm \frac{e^2}{24\pi} |v_{cv}|^2 m^{3/2} \sqrt{\omega_1 - \Delta}, \quad \omega_1 \gtrless \omega, \tag{3}$$

which coincides with the results of [3, 4], where the absorption of the additional field was considered without allowance for the gap. The minus sign corresponds to negative absorption, since population inversion is obtained for the frequency $\omega_1 < \omega$.

In the region $|\omega - \omega_1| < 2\lambda$ the absorption of the field A_1 differs noticeably from that predicted in [3, 4]. In the case of an isotropic gap $\lambda_p = \lambda = \text{const}$, the absorption coefficient

$$K(\omega_1) = 0, \quad |\omega - \omega_1| \leqslant 2\lambda \tag{4}$$

vanishes in the region of 4λ. If we put $\lambda_p = \lambda \cos\theta$ [1], we obtain for $p_0^2/2m \gg \lambda$

$$K(\omega_1) = \frac{5\pi e^2 m p_0 |v_{cv}|^2}{64\omega_1} [a\cos^2\phi + b\sin^2\phi], \tag{5}$$

$$a = x^3, \quad b = \frac{6}{5} x \left(1 - \frac{5}{12} x^2\right), \quad x = \frac{\omega_1 - \omega}{2\lambda}$$

ϕ is the angle between $\vec{A}$ and $\vec{A}_1$. Thus the absorption becomes anisotropic and at $\phi = 0$ it practically vanishes in the interval λ. Comparison of the frequency dependence of the absorption (4) and (5) with the result of [3] is made in Fig. 2. The presence of a gap in the spectrum causes the absorption coefficient to vanish in the region of frequency variation (of the order of λ), whereas the results of [3, 4] call for vanishing only at the point $\omega = \omega_1$. This result may be important for the theory of the semiconductor laser.

3. Let us consider the conductivity of a semiconductor in a constant homogeneous field $\vec{E}_0$ at $T = 0$. It is obvious that in the case of an isotropic gap λ the current is equal to zero if $\vec{E}_0$ is small. Indeed, the electrons cannot acquire energy for the same reason as in a completely filled crystal band. For an anisotropic gap, the current is equal to zero when $\vec{E}_0$ is parallel to $\vec{A}$. This can be shown by using formula (33) of [1] and taking into account the fact that when $\vec{q} = 0$ the longitudinal and transverse conductivities coincide. For example when $\vec{E}_0$ is parallel to $\vec{A}$,

$$\sigma = \lim_{\Omega \to 0} \frac{K(\Omega)}{\Omega} \sim \Omega^2 = 0,$$

where Ω is the frequency of the field $\vec{E}_0$. In the case of a finite temperature, the conductivity becomes different from zero because of the thermal spilling of the electrons through the gap. Determining the density n_{eff} of such electrons, we can expect the conductivity to take the form

$$\sigma = \tau e^2 n_{eff} / m,$$

where τ is the momentum relaxation time, and

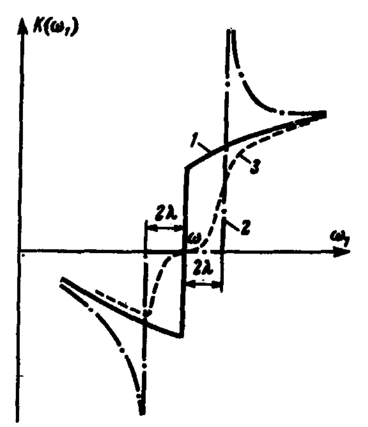

Fig. 2

$$n_{\text{eff}} \sim p_0^3 \begin{cases} \exp(-\lambda/T), & \lambda_p = \lambda \\ (T/\lambda)^3, & \lambda_p = \lambda\cos\theta,\ E_0 \parallel A \end{cases} .$$

When the field E_0 is increased, "breakdown" occurs, wherein the electrons jump through the gap as a result of the energy acquired from E_0.

We note in conclusion that an experimental investigation of the effects under consideration is feasible at the present time. It calls for fields $E \approx 3 \times 10^4 - 10^5$ V/cm with $\lambda \sim 3 \times 10^{-3} - 10^{-2}$ eV.

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GRAVITATIONAL RESONANT DETECTOR WITH TWO DEGREES OF FREEDOM

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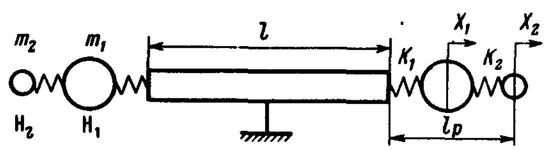
The sensitivity of resonant detectors for gravitational waves can be increased, as noted in [1], by increasing the initial displacement, this being accomplished by introducing a rigid (at the given frequency) rod in a break of the resonant coupling. In addition, one can use as a mechanical amplifier of resonant oscillations (as will be shown below) a system with two degrees of freedom. Such a detector is illustrated schematically in the figure.



Let us estimate first the permissible rod length in the field of a gravitational wave defined by a parameter h and a frequency ω, i.e., the length at which the condition $\Delta\ell \ll \zeta$ is still satisfied (here $\Delta\ell = F\ell/ES$ is the quasi-static displacement of the ends of the rod, and $\zeta = h\ell/2$ is the reduction of the spatial distance ℓ in the field of the wave). Rewriting the condition in the form $\zeta = n\Delta\ell$ ($n \gg 1$) and substituting in $\Delta\ell$ the expression for the force $F = \omega^2 m\ell h/4$ [2], we obtain for ℓ after simple transformations

$$\ell = \sqrt{\frac{2}{n}}\ \frac{v_s}{\omega}, \tag{1}$$

where v_s is the velocity of propagation of the transverse oscillations in the rod.

Let us consider now a resonant system with two degrees of freedom (corresponding to one half of the detector, since the center of mass of the system is at rest in the field of the wave). The damping in the system is determined by the friction forces acting on the masses:

$$\begin{cases} m_1\ddot{x}_1 + (k_1 + k_2)x_1 - k_2x_2 + H_1\dot{x}_1 = F_1 \\ m_2\ddot{x}_2 + k_2x_2 - k_2x_1 + H_2\dot{x}_2 = F_2 \end{cases} . \tag{2}$$