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THEORY OF PARAMETRIC EXCITATION OF SURFACE WAVES BY LASER RADIATION IN NARROW-GAP SEMICONDUCTORS

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1. The parametric resonance phenomenon in narrow-gap semiconductors represents particular interest because of the presence of a new specific mechanism of energy transfer from the driving field to conducting electrons through the vibrations of their effective mass m determined by the pseudorelativistic energy dispersion formula:

$$E(\vec{p}) = [(mc^{*2})^2 + c^{*2}p^2]^{1/2}, \tag{1}$$

where $\vec{p}$ is the electron canonical momentum and $c^* = (E_g/2m)^{1/2}$ (E_g is the value of the band gap).

A complete quantum consideration of such an excitation mechanism for bulk semiconductors was performed in /1/. The present work is devoted to the treatment of surface waves. Let our system consist of two half-infinite spaces filled up by "pseudorelativistic" solid state plasmas with different physical parameters (E_g, c^*, n, etc...) and divided by the boundary xy-plane. The driving field is supposed to be incident normally to this plane and is represented in the well-known dipole approximation by a homogeneous oscillatory electric field

$$\vec{E}_0(t) = \vec{E}_0 \sin\omega_0 t. \tag{2}$$

To investigate such a system we shall proceed from the basic equations of cold-plasma hydrodynamics and the complete system of Maxwell equations that can be written in the form:

$$\left[\frac{\partial}{\partial t} + (\vec{v}\cdot\vec{\nabla})\right]\frac{\vec{v}}{\sqrt{1-v^2/c^{*2}}} = \frac{e}{m}\left\{\vec{\nabla}\phi + \frac{1}{c}\frac{\partial\vec{A}}{\partial t} - \frac{1}{c}[\vec{v}\times\mathrm{rot}\,\vec{A}]\right\}, \tag{3}$$

$$\frac{\partial n}{\partial t} + \mathrm{div}\, n\vec{v} = 0, \tag{4}$$

$$\Delta\vec{A} - \frac{1}{c^2}\frac{\partial^2\vec{A}}{\partial t^2} = \frac{1}{c}\frac{\partial\vec{\nabla}\phi}{\partial t} + \frac{4\pi e}{c}n\vec{v}, \quad \mathrm{div}\,\vec{A} = 0, \tag{5}$$

$$\Delta\phi = 4\pi e(n - n_{0i}) . \tag{6}$$

Here, $\vec{A}$ and ϕ are the vector and scalar potentials of the self-consistent perturbation field, $\vec{v}$ and e are the electron velocity and charge, respectively, c is the light velocity in vacuum, n the density of electrons and n_{0i} the equilibrium ion density that helps to conserve the neutrality of plasma in the equilibrium state.

The linearization procedure for the system of equations (3)-(6) must be carried out with respect to the deviations of physical quantities from the equilibrium state characterized by the vibration velocity

$$\vec{v}_0(t) = \frac{\vec{v}_E \cos\omega_0 t}{\sqrt{1+\beta^2\cos^2\omega_0 t}}, \quad \beta = \frac{v_E}{c^*}, \quad \vec{v}_E = \frac{e\vec{E}_0}{m\omega_0} . \tag{7}$$

2. As is well known, in the chosen geometry surface waves are always of TM-type. In the absence of external fields, if the retardation effect is neglected, these are surface plasmons of frequency $\omega_{ps}^2 = [(\omega_{p1}^2 + \omega_{p2}^2)/2]^{1/2}$ ($\omega_{pj}^2 = 4\pi e^2 n_j/m$, index $j = 1,2$ indicates the two contacting media).

The procedure described above (with $E_0 \neq 0$ and $c = \infty$) leads to the system of equations for Fourier harmonics of the scalar potential ϕ to which appropriate boundary conditions must be applied. Attention should be paid here to the fact that the dielectric function $\epsilon(\omega)$ now is modified by the boundary effect in the presence of the laser field. In the resonant situation when $\omega_0 = \omega_{ps}$ calculations show that the surface plasmon mode ω_{ps} in the case of $\vec{k} \perp \vec{E}_0$ grows with the rate γ expressed by:

$$\gamma = \frac{1}{16}\,\frac{\beta_1^2\omega_{ps1}^2 + \beta_2^2\omega_{ps2}^2}{\sqrt{\omega_{ps1}^2+\omega_{ps2}^2}} . \tag{8}$$

For semiconductor (1)-metal (2) systems γ has the form:

$$\gamma = \frac{1}{16}\,\frac{\beta_1^2\omega_{ps1}^2}{\sqrt{\omega_{ps1}^2+\omega_{ps2}^2}} . \tag{9}$$

In the semiconductor-vacuum case we have:

$$\gamma = (1/16)\,\beta^2\omega_{ps} . \tag{10}$$

In the case of $\vec{k} \parallel \vec{E}_0$ we have instead of (8):

$$\gamma = \left\{\left[\frac{1}{16}\,\frac{\beta_1^2\omega_{ps1}^2 + \beta_2^2\omega_{ps2}^2}{\sqrt{\omega_{ps1}^2+\omega_{ps2}^2}}\right] - \left[J_1^2(\lambda_1)\omega_{ps1}^2 + J_1^2(\lambda_2)\omega_{ps2}^2\right]\right\}^{1/2} . \tag{11}$$

We see that the second term expresses the damping due to the effect of $\vec{k} \parallel \vec{E}_0$-orientation.

The physical picture is more complete when the retardation effect is included. The electromagnetic surface TM-waves described by the initial equations (3)-(6), if E_0 vanishes, undergo the well-known dispersion law:

$$\omega_\perp^2(\vec{k}) = c^2k^2 + \frac{1}{2}(\omega_{p1}^2+\omega_{p2}^2) - c^2k^2\left[1 + \left(\frac{\omega_{p1}^2-\omega_{p2}^2}{2c^2k^2}\right)^2\right]^{1/2} . \tag{12}$$

The main feature of the wave behavior here is excitation under the action of the incident field of a bulk plasma wave propagating along the boundary plane and interacting with the tangential component of the electromagnetic surface wave. Taking into account this phenomenon by introducing an additional boundary condition of $E_z = 0$ for the plasmon field (see /2/, § 10 and /3/), we arrive in the case of $\vec{k} \parallel \vec{E}_0$ at the following system of equations for the Fourier components A_x, A_z of the vector potential:

$$\left[\frac{\partial^2}{\partial z^2} - k^2\Delta(\omega)\right]A_{x,z}(\omega) = \frac{\omega_p^2}{c^2}\sum_{n,\ell}{}' \left\{J_n J_{n+\ell}\left[1 + \frac{(n+\ell)\omega_0}{\omega - n\omega_0}\,\frac{1}{k^2}\,\frac{\partial^2}{\partial z^2}\right] - \frac{1}{8}\beta^2 J_n\left[J_{n+\ell+2} + J_{n+\ell-2}\right]\right\}A_{x,z}(\omega + \ell\omega_0), \tag{13}$$

where $\Delta(\omega) = 1 - (\omega^2/c^2k^2)\,\epsilon_0(\omega)$, $\epsilon_0(\omega) = 1 - \omega_p^2/\omega^2$, J_n is the Bessel function of the first kind of the argument $\lambda = e(\vec{k}\vec{E}_0)/m\omega_0^2 \ll 1$; the prime on the summation means that the term with $n = \ell = 0$ is excluded.

The procedure of obtaining the dispersion equation by using the boundary condition of continuity of E_x and H_y is standard, and we have after lengthy calculations the following expression for the growth rate of the surface polariton mode when $\omega_0 = \omega_\perp(\vec{k})$:

$$\gamma = \{[\frac{1}{4}[1+(\frac{\omega_{p1}^2-\omega_{p2}^2}{2c^2k^2})^2]^{-1/2}\frac{\omega_\perp(\vec{k})}{\omega_{p1}^2-\omega_{p2}^2}\theta(\vec{k},\beta)]^2 -$$

$$-\frac{1}{2}[1+(\frac{\omega_{p1}^2-\omega_{p2}^2}{2c^2k^2})^2]^{-1/2}\frac{\omega_\perp^2(\vec{k})}{\omega_{p1}^2-\omega_{p2}^2}\eta(\vec{k},\lambda)\}^{1/2}, \qquad (14)$$

where

$$\eta(\vec{k},\lambda) = (\omega^2/c^2k^2)\tilde{\eta} + \eta', \qquad \theta(\vec{k},\beta) = (\omega^2/c^2k^2)\tilde{\theta} + \theta',$$

$$\tilde{\eta} = J_1^2(\lambda_2)\omega_{p2}^2\Delta_2(\omega)\epsilon_{01}^2(\omega) - J_1^2(\lambda_1)\omega_{p1}^2\Delta_1(\omega)\epsilon_{02}^2(\omega),$$

$$\eta' = 2\Delta_1(\omega)\Delta_2(\omega)[J_1^2(\lambda_2)\omega_{p2}^2\epsilon_{02}^2(\omega) - J_1^2(\lambda_1)\omega_{p1}^2\epsilon_{01}^2(\omega)],$$

$$\tilde{\theta} = \frac{1}{8}[\beta_2^2\omega_{p2}^2\epsilon_{01}^2(\omega) - \beta_1^2\omega_{p1}^2\epsilon_{02}^2(\omega)],$$

$$\theta' = \frac{1}{4}[\beta_2^2\omega_{p2}^2\epsilon_{02}(\omega)\Delta_1(\omega) - \beta_1^2\omega_{p1}^2\epsilon_{01}(\omega)\Delta_2(\omega)].$$

In the $c \to \infty$ limit eq.(14) tends to the formula (11) as expected. In the case of $\vec{k} \perp \vec{E}_0$ instead of eq.(13) we have:

$$[\frac{\partial^2}{\partial z^2} - k^2\Delta(\omega)]A_{x,z}(\omega) = -\frac{1}{8}\beta^2\frac{\omega_p^2}{c^2}[A_{x,z}(\omega+2\omega_0) + A_{x,z}(\omega-2\omega_0)] \qquad (15)$$

and the growth rate for the surface polariton $\omega_\perp(\vec{k})$ is:

$$\gamma = \frac{1}{4}[1+(\frac{\omega_{p1}^2-\omega_{p2}^2}{2c^2k^2})^2]^{-1/2}\frac{\omega_\perp(\vec{k})}{\omega_{p1}^2-\omega_{p2}^2}\theta(\vec{k},\beta). \qquad (16)$$

It is not difficult to see that in the $c \to \infty$ limit the expression (16) tends to the formula (8).

3. Resuming the performed analysis we can conclude that the excitation condition is more advantageous in the $\vec{k} \perp \vec{E}_0$- geometry than in the $\vec{k} \parallel \vec{E}_0$ one. To show the real values of the threshold field E_{0th} required for the amplification process to begin (this will take place when γ exceeds the linear damping of the mode considered), some numerical estimation was performed for an In Sb sample using the formula (10). The result shows that

$$E_{0th} \approx 3.9\cdot10^4 \text{ v/cm}. \qquad (17)$$

The data employed here are (see /4/): $n = 4\cdot10^{17}$ cm^{-3}, $c^* = 10^8$ cm$\cdot$s^{-1}, $\omega_p = 2{,}7\cdot10^{14}$ s^{-1}, $m = 0{,}016\ m_e$ (m_e is the free electron mass), and the effective electron collision frequency $\nu_{eff} = 0{,}6\cdot10^{12}$ s^{-1} (at T = 77 K).

It is of interest to find what will happen with the same surface plasmon mode for the external field intensity of an order of (17) when $\vec{k} \parallel \vec{E}_0$. For this semiconductor-vacuum case the formula (11) takes the simple form:

$$\gamma = \frac{1}{16}\beta^2\omega_{ps}[1-(\frac{8\xi}{\beta})^2]^{1/2}, \qquad (18)$$

where $\xi = kc^*/\omega_0$. With the data presented above we have $\xi \sim 10^{-1}$, $\beta \sim 10^{-3}$. It is clear that the mode cannot be amplified in this case.

The field values of an order of (17) are undangerous for solids (they are much smaller than atomic fields) and are available by current lasers. Then, such surface waves amplified parametrically in the considered way could be revealed, for example, in some neutron inelastic scattering experiment (see /5/). Such an experiment could serve simultaneously as a method for indirect changing of neutron spectra by laser irradiation /6/.

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