Interaction between two-dimensional plasmons and acoustic waves caused by the deformation potential

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The interaction produced by the deformation potential between a two-dimensional electron plasma and Rayleigh waves is investigated. The cases considered are a plasma layer on the surface of an isotropic dielectric and the inversion channel in a metal-insulator-semiconductor (MIS) structure. In the latter case, the amplitude of the hypersound wave emitted by the plasmons is calculated. It is shown that the two-dimensional character of the plasma leads to significant singularities in the dispersion law and in the absorption of the coupled plasma-acoustic waves. The possibility of observing the considered effects in experiment is discussed.

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Coupled plasma-acoustic waves in piezoelectric crystals were considered in Ref. 1–3. The effect of a plasma layer on the characteristics of an elastic surface wave was considered and the feasibility of hypersound generation by two-dimensional plasmons was demonstrated. The experiments reported to date[4,5] on the observation of two-dimensional plasmons pertain to metal-insulator-semiconductor (MIS) structures on silicon, in which there is no piezoelectric effect. It is therefore of interest to consider the analogous problem for the case when the plasma and acoustic oscillations are coupled via the deformation potential. We recall in this connection that the deformation interaction, unlike the piezoelectric one, increases with frequency. Since we shall be dealing with the hypersound band, it is important to determine the relative contributions of these two electron-phonon coupling mechanisms in the particular situation of "two-dimensional" electrons and a three-dimensional elastic continuum.

It is known[6] that coupling via a deformation potential leads in the homogeneous three-dimensional case both to a change of the sound dispersion in the short-wave region and to the onset of instability. This instability develops starting with a certain value of the wave number ($k > k_0$), but only if the deformation-potential constant Λ exceeds a critical value Λ_c that depends on the electron concentration. In our problem, analogous changes occur in the Rayleigh wave, but the peculiarities of the Coulomb interaction in a two-dimensional system lead to a unique situation. On the one hand, there is no threshold for the coupling constant, i.e., instability sets in formally at any Λ if $k > k_0$. On the other, the numerical value of k_0 is of the order of the reciprocal-lattice vector in all real cases. In this region the macroscopic elasticity theory on which the entire analysis is based no longer holds. An instability of this type does not occur in a two-dimensional system (at least for a monopolar conductor).

In the first section of the article, the problem is considered for the simple model of a plasma layer on the free surface of an elastically isotropic medium. We obtain also the plasmon damping due to hypersound emission. In the second section we calculate the amplitude of the hypersound wave emitted by the plasmons in a typical experimental situation, viz., an MIS structure with an amorphous insulator on the (001) surface of a cubic crystal.

1. TWO-DIMENSIONAL PLASMA ON THE SURFACE OF AN ISOTROPIC ELASTIC MEDIUM

We use the simplest isotropic variant of a deformation potential. The force acting on the electron because of the deformation is

$$\mathbf{F} = \nabla \Lambda \operatorname{div} \mathbf{u},$$

where u is the elastic-medium displacement vector. It can be assumed that Λ is independent of the electron momentum. Neglecting the relatively small change of the elastic properties near the surface, we assume Λ to be equal to its bulk value.

When fluctuations of the electron density $\tilde{n}$ set in, a force $f_v = \nabla \Lambda \tilde{n}$ acts on a unit volume of the elastic medium. We assume the electrons to be two-dimensional in the sense that their density (at equilibrium and with the fluctuations taken into account) is proportional to $|\Psi(z)|^2$, where $\Psi(z)$ is the normalized wave function corresponding to the first quantum level of the transverse motion (the z axis is perpendicular to the surface). In addition, the wavelengths of all the considered oscillations are assumed to be much larger than the thickness of the layer in which the electrons are concentrated. Therefore, calculating the surface density of the force f applied to the elastic medium, and the force acting on the two-dimensional electrons, we shall average the values of F and f with weight $|\Psi(z)|^2$ and take outside the integral with respect to z the slowly varying function $\partial \mathbf{u}/\partial x, \partial \mathbf{u}/\partial z$, etc. We then obtain

$$\tilde{n} = \tilde{n}_s |\Psi(z)|^2, \quad \mathbf{f}_s = \nabla \tilde{n}_s,$$

$\tilde{n}_s$ is the nonequilibrium increment to the plasma surface density and depends only on the coordinates along the surface.

Let the elastic medium fill the half-space $z \leq 0$, and let the plasmon propagate along the x axis with wave

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number k. Since the problem is obviously homogeneous in y, only the components u_x and u_z differ from zero. The solution of the equations of motion of the elastic medium jointly with the Poisson equation is then written in the form

$$u_x = kB\exp(\varkappa_l z) + \varkappa_t A\exp(\varkappa_t z),\quad u_z = -i\varkappa_l B\exp(\varkappa_l z) - ikA\exp(\varkappa_t z),$$
$$\varphi = Ce^{kz},\ z<0;\quad \varphi = De^{-kz},\ z>0;$$
$$\varkappa_l^2 = k^2 - \omega^2/c_l^2,\quad \varkappa_t^2 = k^2 - \omega^2/c_t^2, \tag{1}$$

where c_l and c_t are the velocities of the longitudinal and transverse shear waves, φ is the electrostatic potential, $A, B, C,$ and D are arbitrary constants, and the factor $\exp(ikx - i\omega t)$ was left out.

The boundary conditions at $z=0$ take the form

$$T_{xz} - f_x = 0,\quad T_{zz} = 0,$$
$$-\frac{\partial\varphi}{\partial z}\Big|_{z=+0} + \varepsilon\frac{\partial\varphi}{\partial z}\Big|_{z=-0} = 4\pi e\tilde{n}_s,\quad \varphi|_{z=+0} = \varphi|_{z=-0}. \tag{2}$$

Here T_{ik} is the stress tensor, ε is the dielectric constant, and e is the electron charge.

The value of $\tilde{n}_s$ is obtained from the kinetic equation

$$\frac{\partial f}{\partial t} + \left(v_x\frac{\partial}{\partial x}\right)f - \left(e\frac{\partial\varphi}{\partial x} + \frac{\partial}{\partial x}\Lambda\,\mathrm{div}\,\mathbf{u}\right)\frac{1}{m}\frac{\partial f}{\partial v} = -\frac{f-\bar{f}}{\tau}, \tag{3}$$

where f is the distribution function, m and v are the effective mass and velocity of the electron, τ is the relaxation time, and the superior bar denotes averaging over the momenta.

From (3) we get ($\hbar=1$):

$$\tilde{n}_s = -\frac{ik^2\sigma_{xx}}{e\omega(1-R_x)}\left[\varphi + \frac{\Lambda}{e}\left(iku_x + \frac{\partial u_z}{\partial z}\right)\right],$$
$$\sigma_{xx} = -\frac{2e^2\tau}{(2\pi)^2}\int\frac{v_x^2}{1-i\omega\tau+ikv\tau}\frac{\partial f_0}{\partial E}dp, \tag{4}$$
$$R_x = -\frac{2}{(2\pi)^2}\frac{k}{\omega}\frac{\partial\mu}{\partial N_s}\int\frac{v_x}{1-i\omega\tau+ikv\tau}\frac{\partial f_0}{\partial E}dp.$$

Here μ is the chemical potential, $E=p^2/2m$ is the electron energy, and f_0 is the equilibrium distribution function.

Substituting (1) and (2) and using (4), we obtain the dispersion equation for the coupled plasma-acoustic waves:

$$[1 + i4\pi k\sigma_{xx}/\omega(1-R_x)(\varepsilon+1)][4k^2\varkappa_l\varkappa_t - (\varkappa_t^2+k^2)^2]$$
$$= 2i\Lambda^2k^4\sigma_{xx}\omega\varkappa_l/e^2\rho c_t^2c_l^2(1-R_x), \tag{5}$$

where ρ is the density of the elastic medium.

We investigate first the frequency region typical of acoustic waves, for which we can put $\omega \ll kv$. In the limiting case of strong scattering we have at $kv\tau \ll 1$

$$\sigma_{xx} = \sigma = e^2N_s\tau/m,\quad R_x = -4i\pi\sigma k^2r_0/(\varepsilon+1)\omega,$$
$$r_0 = \frac{\varepsilon+1}{4\pi e^2}\frac{\partial\mu}{\partial N_s} = \frac{\varepsilon+1}{4me^2}\left[1 - \exp\left(-\frac{\pi N_s}{mT}\right)\right]^{-1}, \tag{6}$$

where r_0 is the screening radius of the Coulomb interaction in the two-dimensional system, N_s is the equilibrium concentration of the surface charges, and T is the temperature in energy units. The initial section of the sound branch is of the form $\omega = sk$, where s is the velocity of the Rayleigh wave. The small dispersion corrections to the function $\omega(k)$ are proportional to k^3 at $kr_0 \ll 1$ and to k^2 at $kr_0 \gg 1$. The sound damping by the electrons $\omega'' \equiv -\mathrm{Im}\,\omega$ will be reduced (in the case most accessible to experiment, $\sigma \gg s$, $kr_0 \ll 1$) to the form

$$\omega'' = \frac{[(\varepsilon+1)s\Lambda k]^2}{32\pi^2\rho(\sigma e c_t)^2}\frac{\omega x_t}{R},\quad x_{t,l}^2 = 1 - \frac{s^2}{c_{t,l}^2},$$
$$R = \frac{x_t}{x_l} + \frac{c_t^2x_l}{c_l^2x_t} - x_t^2 - 1 > 0. \tag{7}$$

If the spatial dispersion is strong ($kv\tau \gg 1$) the damping is described by the equations

$$\omega'' = \frac{[(\varepsilon+1)\Lambda v\tau]^2x_t\omega^5}{2(8\pi)^2e^2c_t^2s^3\sigma R\rho} \propto \frac{\omega^5\tau}{N_s},\quad (kv\tau)^2 \ll \frac{\sigma}{s}, \tag{7a}$$

$$\omega'' = \frac{\Lambda^2x_t\sigma\omega}{\rho c_t^2e^2sR(v\tau)^2} \propto \frac{\omega N_s}{\tau},\quad (kv\tau)^2 \gg \frac{\sigma}{s}. \tag{7b}$$

The experiment is easier in case (7a), which is characterized by a very strong frequency dependence ($\sim\omega^5$) of the damping.

An investigation of Eq. (5) with (6) taken into account shows that at sufficiently large k the dispersion law is greatly distorted and a region of anomalously frequency is produced. Assuming that $\omega \ll c_tk$, we expand the second factor of (5) in powers of $\omega^2/c_{t,l}^2k^2$ and reduce the dispersion equation to the form

$$\left(1 + kr_0 - i\frac{\omega}{\omega_0}\right)\left[1 - \frac{\omega^2}{4k^2}\frac{c_l^4+c_t^4}{c_t^2c_l^2(c_l^2-c_t^2)}\right] = \frac{k^2}{Q^2},$$
$$Q^2 = \frac{4\pi e^2\rho(c_l^2-c_t^2)}{\Lambda^2(\varepsilon+1)}, \tag{8}$$

where $\omega_0 \equiv 4\pi\sigma k/(\varepsilon+1)$ is the Maxwellian frequency in the two-dimensional case. As is seen from (8), the positive root of the equation $1+kr_0 = k^2/Q^2$ determines the point k_0 in which $\omega(k)$ vanishes. At the same point, an anomaly occurs in the second solution of the dispersion equation (5), a solution closely connected with the plasma degrees of freedom. At small k this solution describes the Maxwellian relaxation of the charge fluctuation. Near the point k_0 where $\omega \ll c_{t,l}k_0$, the following equation holds:

$$\omega = -i\omega_0(1+kr_0-k^2/Q^2),\quad kv\tau \ll 1,$$
$$\omega = -\frac{2i\omega_0}{(kv\tau)^2}(1+kr_0-k^2/Q^2),\quad kv\tau \gg 1. \tag{9}$$

At $k=k_0$ the sign of ω'' is reversed, i.e., instability sets in. An instability of this type was investigated in detail for the three-dimensional case in Refs. 6 and 7. It is shown there that at $\Lambda > \Lambda_c$ the short-wave fluctuations of the carrier density upset the homogeneous state of the system. A thermodynamically more favored structure is one with a periodic dependence of the carrier density and of the lattice deformation on the coordinate. The threshold value of k at which instability begins to evolve is (in the three-dimensional problem!)

$$k_0 = r_D^{-1}(\Lambda/\Lambda_c - 1)^{-1/2},$$

where r_D is the "three-dimensional" Debye radius. At sufficiently high temperature and low electron density, k_0 can be small enough to land in the region where the macroscopic theory does not hold.

In contrast to the three-dimensional case, in our problem there is no need for a threshold condition on Λ. It is easy to estimate, however, that in any real case we have $Qr_0 \gtrsim 1$. For the smallest possible k_0 (at $T=0$)

we then obtain a value that does not depend on the charge surface density and is determined only by the parameters of the material:

$$k_0 \approx Q^2 r_0 = \pi\rho (c_l^2 - c_t^2)/\Lambda^2 m.$$

The ratio of this quantity and the reciprocal-lattice vector is 0.4 for Si and 0.2 for Ge. An increase of the screening radius r_0 only increases k_0, in contrast to the three-dimensional situation. Thus, the instability of the two-dimensional sound-charge waves in the monopolar case could be observed only in substances with anomalously small elastic moduli and with a large deformation potential. We shall not analyze here this hypothetical case.

In the region $\omega > kv$, $\omega t \gg 1$ the considered solution corresponds to a plasma wave. The two-dimensional plasmon is damped by the electron scattering and in accord with the Landau mechanism. Just as in the case of piezoelectric coupling,[1] additional damping is produced by the emission of volume sound waves. In the case of Fermi statistics this damping is obtained from the expression

$$\omega'' = \frac{(\varepsilon+1)\Lambda^2 c_t^2 k^3}{4\pi e^2 \rho c_l^2}\left(1 - \frac{k^2 v^2}{\omega^2}\right)^{3/2}. \tag{10}$$

The order of magnitude of ω''/ω can be estimated as $(k/Q)^2 c_t k/\omega \ll 1$. Thus, in the case of coupling via the deformation potential, the obtained damping is much less than in the analogous problem with piezoelectric coupling (see Ref. 1). We note that in the considered frequency region the sound wave emitted by the plasma is almost a shear wave: the longitudinal component of the displacement is smaller than the transverse one in a ratio $c_t k/\omega$. The sound wave propagation is practically perpendicular to the surface here.

2. SOUND-WAVE EMISSION BY TWO-DIMENSIONAL PLASMONS IN AN MIS STRUCTURE

In this section we consider a situation typical of experiments with silicon MIS structures.[4] An amorphous SiO_2 insulator film of thickness Δ is deposited on the (001) surface of silicon filling the half-space $z \leqslant 0$. Let the plasmon propagate in the [100] direction (the x axis). We shall now write down the obvious system of equations, and give only the form of the boundary conditions:

$$\begin{gathered}(T_{xz} - f_x)|_{z=-0} = T_{xz}|_{z=+0},\ T_{zz}|_{z=-0} = T_{zz}|_{z=+0},\\ \mathbf{u}|_{z=-0} = \mathbf{u}|_{z=+0},\ -\varepsilon_2 \partial\varphi/\partial z|_{z=+0} + \varepsilon_1 \partial\varphi/\partial z|_{z=-0} = 4\pi e \tilde{n}_s,\\ \varphi|_{z=-0} = \varphi|_{z=+0},\ T_{zz}|_{z=\Delta} = 0,\\ T_{xz}|_{z=\Delta} = 0,\ \varphi|_{z=\Delta} = 0,\end{gathered} \tag{11}$$

where ε_1 and ε_2 are the dielectric constants of the semiconductor and of the insulator, respectively.

The reasoning of the end of the preceding section remains in force in the present problem: the radiation consists almost completely of shear waves in a direction that makes a small angle with the normal to the surface, in view of the smallness of the parameter ck/ω, where c is a quantity on the order of the speed of sound in the insulator and in the semiconductor

The complete solution of this problem entails very cumbersome calculations. We confine ourselves therefore to a qualitative analysis of the results. The acoustic damping of the plasmon is described by an equation similar to (10), which contains a factor of the order of unity dependent on the densities and elastic moduli of the insulator and of the semiconductor. In addition, $\varepsilon + 1$ is replaced by $\varepsilon_1 + \varepsilon_2 \coth k\Delta$. Thus, in contrast to (10), at $k\Delta \ll 1$ the plasmon damping becomes proportional to k^2. We assume here and below satisfaction of the condition

$$\mathrm{cth}\, k\Delta \ll (\omega/ck)\, e^2 \rho c^2/\Lambda^2 k^2,$$

which is always satisfied in the present experiments with two-dimensional plasmons.

To obtain the amplitude of the sound wave emitted in the interior of the crystal, we solve the homogeneous problem in which the kinetic equation (3) includes the external electric field $\mathbf{E}\exp(ikx - i\omega t)$, $\mathbf{E} \parallel \mathbf{k}$, that acts on the plasma. Assuming that the dominant damping mechanism is the electron scattering,[1)] the sought amplitude A is given by

$$A = i\frac{E\Lambda k(\varepsilon_1 + \varepsilon_2\, \mathrm{cth}\, k\Delta)\,\omega_p^2/\omega^2}{4\pi e\omega(\rho_1 c_1 - i\rho_2 c_t\, \mathrm{tg}\, \omega_p\Delta/c_t)}\left[1 - \frac{\omega_p^2}{\omega^2}\left(1 - \frac{i}{\omega\tau}\right)\right]^{-1}, \tag{12}$$

where $\omega_p^2 = 4\pi e^2 N_s k/(\varepsilon_1 + \varepsilon_2 \coth k\Delta)m$, ρ_1 and ρ_2 are the respective densities of the semiconductor and of the insulator, c_t is the transverse sound velocity in the insulator, and c_1 is the velocity of the shear wave in the [100] direction of the cubic crystal.

The quantity A in (12) is the amplitude of the x component of that part of the displacement whose divergence is equal to zero. The corresponding wave propagation velocity is c_1. Outside the vicinities of the acoustic resonances Eq. (12) determines the principal part of the displacement in the radiated sound wave. In addition, we have the z component of the displacement and that part of the x component whose curl is equal to zero. Their contributions to the total displacement are proportional respectively to Ack/ω and $A(ck/\omega)^3$. We shall not carry out here the cumbersome investigation of the vicinities of the acoustic resonances, and note only that the radiated sound energy decreases noticeably when $\omega\Delta/c_t$ is close to $(n + \frac{1}{2})\pi$, where n is an integer.

From a comparison of (12) with the analogous expression for the amplitude of the emitted sound in the case of piezoelectric coupling (see Ref. 3) it is seen that the order of magnitude of the effect can be estimated by replacing the piezoelectric modulus β by the expression

$$\Lambda k(\varepsilon_1 + \varepsilon_2\, \mathrm{cth}\, k\Delta)/4\pi e.$$

If the electromagnetic wave acting on the plasma has an intensity of the order of 1 W/cm^2, the intensity of the hypersound wave is of the order of 10^{-8} W/cm^2. To obtain this estimate we used the parameters of a silicon MIS structure and assumed $N_s \sim 10^{12}$ cm^{-2}, $k \approx 2 \cdot 10^4$ cm^{-1}, $\Delta \approx 2 \cdot 10^{-5}$ cm, $\omega\tau \sim 10$.

The presented estimate of the sound-wave intensity is smaller by approximately four decades than in the case of piezoelectric coupling (see Ref. 3), but is still within the capabilities of contemporary experimental techniques. We noted that this low ultrasound intensity is still larger by four decades than that given by the electrostriction effect even for the best ferroelectrics. In the considered case of a silicon MIS structure, the difference amounts to 12 decades (the electrostriction constant is of the order of 10^{-14} cgs esu).

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1) This is precisely the case realized in the experiments reported in Refs. 4 and 5.

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Effect of pressure on the hyperfine magnetic fields at the nuclei of iron impurity atoms in antiferromagnetic chromium

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The pressure dependences of the hyperfine magnetic fields H at ^{57}Fe nuclei in alloys of chromium with alloys were measured by the Mössbauer gamma-spectroscopy method. Values $\Delta H/H\Delta p = (-1.2\pm0.5)\times10^{-2}$ and $(-2\pm0.5)\times10^{-2}$ kbar^{-1} were obtained for alloys with concentrations 1.5 and 3.5 at.% Fe, respectively. The values of $\Delta H/H\Delta p$ are anomalously large and agree approximately with the relative changes of the magnetic moments of the Cr atoms under pressure $\Delta\mu_{Cr}/\mu_{Cr}\Delta p$. Doubts are expressed concerning the model of H described in the paper of Herbert *et al*. [J. Phys. Chem. Solids **33**, 979 (1972)]. An alternate mechanism for the onset of H is proposed, based on concepts concerning the nature of the field at nuclei of nonmagnetic atoms dissolved in a magnetic host.

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Hyperfine magnetic fields at nuclei of impurity magnetic atoms in magnets are determined by the values of the magnetic moments of the impurity atoms, the host atoms, and the impurity-host interaction. Accordingly, the resultant hyperfine field H should consist of several contributions. The values and the signs of the individual contributions to H cannot be predicted beforehand, since the mechanism of the transport of the spin density from the atoms of the magnetic host to the nuclei of the impurity atoms is not yet clear. Important information on the role of the individual contributions to H is made by experiments on the influence of the pressure on the hyperfine interaction in magnets with different types of magnetic ordering. So far, however, the pressure dependences of the hyperfine field $H(p)$ at the impurity magnetic atoms in monatomic antiferromagnets (consisting of atoms of the same species) have not been investigated. The present study fills this gap in part. We have investigated the $H(p)$ dependences at ^{57}Fe nuclei in three diluted alloys of chromium with iron.

Metallic chromium can have an unusual variety of magnetic structures, depending on the temperature, species, and concentration of the impurities.[1] Below the Néel temperature $T_N = 311$ K, the magnetic moments of Cr form spin-density waves that can be represented in the form

$$\mathbf{P}(x) = P_0 \mathbf{n} \cos Qx, \tag{1}$$

where $\mathbf{P}(x)$ is the spin-polarization vector, $\mathbf{n}$ is a unit vector along the polarization vector, $\mathbf{Q}$ is the wave vector and is parallel to one of the $\langle 100\rangle$ directions in a bcc lattice, and z is the coordinate. The SDW are polarized transversely ($\mathbf{n}\cdot\mathbf{Q}=0$), in the temperature interval $T_N > T > T_{SF}$, where $T_{SF} \approx 122$ K, and longitudinaly ($\mathbf{n}\cdot\mathbf{Q}=q$) in the region $T < T_{SF}$. At $T < 70$ K we have $Q = 0.95162\pi/d$ (d is the lattice parameter) and in the interval 70 K $< T < T_N$ the value of Q increases to $0.9626\cdot 2\pi/d$; the SDW length is then of the order of $20d$ and is not commensurate with d.

Introduction of the Fe impurity influences strongly the magnetic properties of the chromium: T_N decreases by approximately 20 K per at.% Fe, T_{SF} decreases to zero at 1.5 at.% Fe, and Q becomes equal to $2\pi/d$ at a concentration higher than 2.5 at.% Fe. The Fe im-

