Eur. Phys. J. B **45**, 325–330 (2005)
DOI: 10.1140/epjb/e2005-00185-4

THE EUROPEAN PHYSICAL JOURNAL B

High frequency properties of a quasi-two-dimensional conductive film

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Received 23 November 2004 / Received in final form 1st March 2005
Published online 6 July 2005 –

Abstract. The propagation of a monochromatic longitudinal acoustic wave along the low conductivity axis of a thin quasi-two-dimensional conductive film, in the absence of an external magnetic field, is studied theoretically. It is shown that under certain conditions the formation of both a standing ordinary wave (OAW) and an anomalous acoustic wave (AAW) is possible. The frequency dependence of the amplitudes of both waves is derived. For certain values of the characteristic parameters, the AAW in the film may be dominant. From the resonance conditions for the formation of standing OAW and AAW waves (especially the AAW), it is possible to obtain information about the electronic structure of the quasi-two-dimensional conductors, e.g. the corrugation parameter η or the relaxation properties of the charge carriers.

PACS. 72.10.-d Theory of electronic transport; scattering mechanisms – 72.30.+q High-frequency effects; plasma effects – 72.50.+b Acoustoelectric effects

1 Introduction

The search for new superconducting materials is closely related to the interest in low-dimensional conductors. Most of these conductors are layered structures of organic origin with a sharply pronounced anisotropy in the conductivity: their electrical conductivity in the layers is significantly higher than normal to the layers. Layered conductors of organic origin are attractive for experimenters owing to their peculiar behaviour in strong magnetic fields and to a number of phase transitions at comparatively low pressures.

The discovery of Shubnikov-de Haas oscillations of the magnetoresistance in tetracyano-tetracen halogens and in a large family of tetrathiafulvalene based ion-radical salts with charge transfer in magnetic fields of several tens of teslas suggests that these compounds have metallic type conductivity even across the layers and that their carrier mean-free path ℓ reaches several micrometers [1–3]. This suggests that it might be reasonable to apply to the electronic description of these systems the well-developed concept of charge carrying quasi-particles in metals.

The low-dimensional electron energy spectrum of such conductors is quasi-two-dimensional (quasi-2D) in character

$$\varepsilon(\vec{p}) = \sum_{n=0}^{\infty} \varepsilon_n(p_x, p_y) \cos\left(\frac{anp_z}{\hbar}\right), \qquad (1)$$

i.e. it depends weakly on the momentum component $p_z = \vec{p}\vec{n}$ along the normal $\vec{n}$ to the layers, while the functions $\varepsilon_n(p_x, p_y)$ decrease sharply with increasing n: $\varepsilon_{n+1} = \eta\varepsilon_n \ll \varepsilon_n$. Here $\hbar$ is Planck constant, a is the lattice period in the z-direction and η is the quasi-2D parameter of the spectrum. The corresponding Fermi surface is a weakly corrugated open cylinder.

Over the past decade, high-frequency phenomena in layered conductors under magnetic fields were extensively studied, both experimentally [4–7] and theoretically [8–18]. The theoretical studies of high-frequency acoustic phenomena in these conductors show a number of characteristic features which appear distinct from the corresponding properties of three-dimensional conductors. In this respect, it might be of interest to perform a theoretical analysis of the high-frequency acoustic phenomena since they are highly informative and can be used successfully for a detailed study of the electronic structure of layered conductors, in particular for obtaining the dispersion relation and the relaxation properties of the charge carriers.

To the best of our knowledge, only two theoretical papers were published concerning high-frequency acoustic phenomena in quasi-2D conductors in the absence of an external magnetic field [19,20]. In reference [19] it was shown that, apart from the acoustic wave which propagates at the sound velocity (ordinary acoustic wave – OAW), there exists also a non-exponentially attenuated anomalous acoustic wave (AAW) propagating with a velocity $\sim \eta v_0$. The ordinary acoustic wave is attenuated at distance $\ell_{at} = \omega^*/\eta\, v_0$. The electrons which

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participate in the attenuation have velocities that are almost normal to the wave vector $\vec{q}\cdot\vec{v} \approx 0$, i.e. these are the electrons which are in phase with the sound wave. The excitation of the AAW is due to the “dragging” of the acoustic field directly by the conduction electrons. Specifically, when a longitudinal acoustic wave propagates in a conductor, the lattice deformations displace the electronic subsystem from equilibrium and thereby generate a longitudinal electric field. The momentum which the electrons acquire from this longitudinal field is transmitted to the lattice at distances of the order of their mean free path ℓ. It means that there exists a mechanism “dragging” the acoustic field at distances of the order of ℓ. The electrons which take part in this “dragging” of the acoustic field in 3D-metals have velocities along the sound wave vector ($v = v_F$) and they do not contribute to the sound attenuation at all. In quasi-2D conductors, in which the Fermi surface is a weakly corrugated cylinder open in the direction of propagation of the acoustic wave, these are the electrons whose velocity component in the direction of propagation of the sound reaches the highest value, $v_z^{max} = \eta\, v_0$. The acoustic wave which “survives” at distances ℓ is known as the AAW.

In a bulk conductor, at small distances from the boundary surface, the amplitude of the AAW is insignificant compared to the OAW. The situation is quite different if one considers the propagation of a high-frequency acoustic wave in a thin quasi-2D conductive film whose thickness d is much smaller than the mean-free path of the electrons ℓ, i.e. $d \ll \eta\ell$. In this case, at certain frequencies of the external acoustic field, a resonance of the AAW can be observed, i.e. standing AAWs can form.

2 Formulation of the problem

Consider a monochromatic longitudinal acoustic wave propagating along the low conductivity axis (z-axis) of a quasi-2D conductive film ($0 \le z \le d$). The displacement $U(0,t) = U_0 \exp(-i\omega t)$ of the surface ($z = 0$) is given. In the following the behaviours of an ordinary acoustic wave, $U^{\rm OAW}(z)$, and of an anomalous acoustic wave, $U^{\rm AAW}(z)$, are theoretically analysed.

The complete system of equations describing the propagation of acoustic waves in conducting media consists of the Boltzmann kinetic equation for the charge carrier distribution function, $f(\vec{r},\vec{p},t) = f_0(\varepsilon) - \chi(\vec{r},\vec{p})\,\partial f_0/\partial\varepsilon$ (here $\chi(\vec{r},\vec{p})$ is the non-equilibrium additive term to the equilibrium distribution function), the Maxwell equations, and the equations of elasticity [21].

The relaxation processes in quasi-2D conductors generally become more involved than in three-dimensional conductors, especially when the scattering on phonons is taken into account. However, if we assume that the temperature is fairly low, so that the charge carriers scatter mainly on impurities, we can confine the discussion to the relaxation-time approximation and analyse the distinctive features introduced into the phenomenon, i.e. the quasi-2D nature of the electron energy spectrum.

The Boltzmann kinetic equation for charge carriers, written in the relaxation-time approximation, takes the form

$$\frac{\partial\chi}{\partial z} - i\frac{\omega^*}{v_z}\,\chi = e\,E_z - \frac{i\omega}{v_z}\,\Lambda_{zz}\frac{\partial U}{\partial z} \tag{2}$$

where $\omega^* = \omega + i\nu$, $\vec{v} = \{v_x, v_y, v_z\}$ is the velocity of the charge carriers and Λ_{zz} is the deformation potential tensor component. For brevity, we restrict the discussion to the first two terms in equation (1) and set:

$$\varepsilon(\vec{p}) = \frac{p_x^2 + p_y^2}{2m} - \eta\frac{\hbar}{a}v_0\cos\left(\frac{ap_z}{\hbar}\right), \qquad v_{0^2} = \frac{2\varepsilon_F}{m}. \tag{3}$$

In such a spectrum, the reduced deformation potential Λ_{zz} depends only on p_z:

$$\Lambda_{zz} = \eta L\cos\left(\frac{ap_z}{\hbar}\right), \qquad L \approx \varepsilon_{\rm F}. \tag{4}$$

The effective interaction of electrons in a quasi-2D conductor with an acoustic wave propagating along its low conductivity axis is weakened to the extent that η is small. In the purely two-dimensional case, $\eta \to 0$ and Λ_{zz} vanishes.

The longitudinal electric field $E_z(z)$ in conductors with a high charge-carrier density can be derived from the electrical neutrality conditions $\langle\chi\rangle = 0$, which is equivalent to the continuity condition for the current, i.e.

$$j_z(z) = e\langle v_z\chi\rangle = 0. \tag{5}$$

The angular brackets indicate standard integration over the Fermi surface

$$\langle\cdots\rangle = -\frac{2}{(2\pi\hbar)^3}\int(\cdots)\frac{\partial f_0}{\partial\varepsilon}d^3p. \tag{6}$$

The acoustic field $U(z)$ in the film can be described by (2), (5) and the equation of elasticity:

$$\frac{\partial^2 U}{\partial z^2} + \left(\frac{\omega}{s}\right)^2 U = \frac{1}{\varrho s^2}\Big\langle\Lambda_{zz}\frac{\partial\chi}{\partial z}\Big\rangle, \tag{7}$$

where ϱ is the density of the metal and s is the acoustic wave velocity.

The behaviour of the acoustic field in the film generally depends on the type of scattering of the charge carriers on the boundaries. Since the ratio $v_z^{max}/v_F = \eta$ is small, only “glancing” electrons are present, and the specular boundary condition can be taken as an approximation for the electron distribution function [22]:

$$\chi_-(z=0) = \chi_+(z=0), \qquad \chi_-(z=d) = \chi_+(z=d), \tag{8}$$

where χ_- and χ_+ are the non-equilibrium terms produced by conduction electrons incident on and reflected from the surface, respectively.

$$U^{\mathrm{OAW}}(z) = U(0)\,\frac{f^2 + 2\gamma\,[1 + i\,c - \sqrt{(1+i\,c)^2 - \alpha^2 f^2}]}{f^2 + \gamma\, f^2\alpha^2\,[(1+i\,c)^2 - \alpha^2 f^2]^{-1/2}}\,\frac{\cos[a\,f\,(1-2\frac{z}{d})/2]}{\cos[a\,f/2]}, \tag{20}$$

3 Calculations

After Fourier transforming equations (2), (5), and (7) with respect to z, one obtains the following algebraic solution for the acoustic field:

$$U^n = \frac{n\pi}{d}\,\frac{[(-1)^n\,U(d) - U(0)][1+G_n]}{-\,(\omega s)^2 + (n\pi d)^2\,[1+G_n]}, \tag{9}$$

where $U(0)$ and $U(d)$ are displacements at $z = 0$ and $z = d$, respectively. The function G_k is

$$G_k = \frac{1}{\varrho s^2}\frac{\omega^*}{\omega}\Big[\Big\langle\frac{\omega^2\Lambda_{zz}^2}{v_z^2}\,1(\omega^*/v_z)^2 - k^2\Big\rangle - k\frac{\omega}{\omega^*}\Big\langle\Lambda_{zz}(\omega^*/v_z)^2 - k^2\Big\rangle^2\Big\langle\frac{1}{(\omega^*/v_z)^2 - k^2}\Big\rangle^{-1}\Big], \tag{10}$$

and $k = n\pi/d$. Integrating over the Fermi surface by (6), one obtains

$$G_k = \frac{mL^2}{2\pi^2\hbar^2\varrho a}\,\frac{\omega\omega^*}{k^2 s^2 v_0^2}\left\{1 - \sqrt{1 - \left(\frac{k\eta v_0}{\omega^*}\right)^2}\right\}. \tag{11}$$

Using the inverse Fourier transformation, the acoustic field in the film can be written

$$U(z) = A\,U(0) + BU(d), \tag{12}$$

where A and B are the following sums:

$$A = \frac{1}{d}\sum_{n=-\infty}^{\infty}\frac{\frac{n\pi}{d}(1+G_n)\sin\frac{n\pi z}{d}}{-\left(\frac{\omega}{s}\right)^2 + (\frac{n\pi}{d})^2(1+G_n)},$$
$$B = -\frac{1}{d}\sum_{n=-\infty}^{\infty}\frac{(-1)^n\frac{n\pi}{d}(1+G_n)\sin\frac{n\pi z}{d}}{-\left(\frac{\omega}{s}\right)^2 + (\frac{n\pi}{d})^2\,(1+G_n)}. \tag{13}$$

Applying to Poisson formulae [23] the sums A and B can be represented in the following forms:

$$A = \frac{1}{\pi}\left\{2\sum_{r=1}^{\infty}J_1(r) - iJ_2(r)\right\};$$
$$B = -\frac{1}{\pi}\left\{\sum_{r=1}^{\infty}[J_3(r) + J_4(r)] + J_5(r)\right\}, \tag{14}$$

where

$$J_1(r) = \int_{-\infty}^{\infty}F(k)\sin kz\exp[2irkd]dk;$$
$$J_2(r) = \int_{-\infty}^{\infty}F(k)\exp[ikz]dk,$$
$$J_3(r) = \int_{-\infty}^{\infty}F(k)\sin kz\exp[i(2r-1)kd]dk;$$
$$J_4(r) = \int_{-\infty}^{\infty}F(k)\sin kz\exp[i(2r+1)kd]dk,$$
$$J_5(r) = \frac{1}{2i}\int_{-\infty}^{\infty}F(k)\,\{\exp[ik(d+z)] - \exp[ik(d-z)]\}\,dk, \tag{15}$$

and

$$F(k) = \frac{k[1+G_k]}{-\left(\frac{\omega}{s}\right)^2 + k^2[1+G_k]} \tag{16}$$

Substituting equation (4) in equation (9), and after some manipulations, one obtains the dispersion function

$$D(k) = \left(\frac{ks}{\omega}\right)^2\left[1+G_k\right] - 1, \tag{17}$$

which besides the zeros $k = \pm k_0$ with

$$k_0 = \frac{\omega}{s}\left[1-\gamma-(\gamma\alpha)^2 - i\gamma\frac{\nu}{\omega} + \gamma\sqrt{\left(1+\gamma\,\alpha^2 + i\frac{\nu}{\omega}\right)^2 - \alpha^2}\right], \tag{18}$$

also has branch points $k_1 = \pm\omega^*/(\eta v_0)$. Hence, the solution can be represented as a sum of an OAW (U^{OAW}) whose velocity of propagation is close to the sound velocity s, and an AAW (U^{AAW}) determined by electrons with velocity $v_z \simeq \eta v_0$. Here,

$$\gamma = \frac{mL^2}{\pi^3\hbar^2\varrho\, a\, v_0^2} \sim \frac{m}{M},$$

where m and M are the electron and ion masses, respectively.

Applying the theorem of Cauchy, each integral in (15) can be represented as

$$J_n = -J_n^{\Gamma} + 2i\pi\sum_j Res_j, \qquad (n = 1,2,3,4,5). \tag{19}$$

The second term in (18) is connected with the OAW. Its amplitude can be determined if the residues at the roots of the dispersion equation (17) are calculated,

see equation (20) above

where the following notations are used,

$$f = \frac{s}{\omega}\,k_0; \qquad a = q\,d; \qquad \alpha = \frac{\eta\,v_0}{s}; \qquad c = \frac{\nu}{\omega}. \tag{21}$$

The amplitude of the AAW is determined by the contour integral along the "edges" of the cut in the complex plane from the branch point $k_1 = \omega^*/(\eta v_0)$ of the dispersion function $D(k)$ to infinity. It can be written in the following form

$$U^{\mathrm{AAW}}(z) = U(0)\frac{\gamma}{i}\int_1^{\infty}(1+ic)f(x)\frac{\cos[\frac{a}{2}\frac{x}{\alpha}(1+ic)(1-2\frac{z}{d}]}{\cos[\frac{a}{2}\frac{x}{\alpha}(1+ic)]}, \tag{22}$$

where

$$f(x) = \frac{\sqrt{x^2-1}}{x\,[1-(1+i\,c)^2x^2/\alpha^2]^2}; \qquad x = k\eta v_0/\omega^*. \tag{23}$$

Equations (20) and (22) have been obtained for the symmetrical case, $U(0) = U(d)$. Analogous relations can be derived for asymmetric case as well, i.e. for $U(0) = -U(-d)$.

4 Discussion

The anomalous acoustic field in the case of a quasi-2D electron energy spectrum is controlled by electrons near the extremum v_z of the Fermi surface, reached at $p_z = \pi\hbar/2a$. The projection of the electron velocities onto the normal to the layers is then $(v_z)_{ext} = \eta\, v_0$ and it is the velocity of propagation of the anomalous quasi-wave. For these electrons, it follows from (23) that the value of x is close to unity, $x \sim 1$. Therefore the upper limit of the integral in equation (22) can be cut at some value that depends on the frequency of the external acoustic field, $x_1(\omega)$. For $z = d/2$ and frequencies $1\,\pi \leq \omega d/(\eta v_0) \leq 20\,\pi$ this limit is in the range $1.2 \geq x_1 \geq 1.03$. It is obvious that with increasing frequency, the value of x_1 decreases and so does the amplitude of the AAW. Presumably, the reason for this is that an increasing frequency produces larger field inhomogeneities and so the "synchronization" can be achieved with a smaller number of electrons.

Equations (20) and (22) show that under some conditions, standing OAWs as well as AAWs can be formed in the film. In the collisionless limit, $\nu/\omega \to 0$, the resonance frequencies at which standing AAWs occur are determined by the conditions

$$\frac{a}{\alpha} = \frac{\omega_n d}{\eta v_0} = (2n+1)\pi, \qquad (n = 0, 1, 2, 3, \cdots). \tag{24}$$

The amplitudes of both waves depend on the parameters $a = q\,d$, $\alpha = \eta v_0/s$, and $c = \nu/\omega$. In the above expressions two new parameters can be introduced, $\theta = a/\alpha$ and $b = d/(\eta\ell)$, so that the only parameter that depends on the frequency is now θ. In the following, the amplitudes of the OAW and AAW are analysed in function of the parameters α, θ and b.

For $\alpha \gg 1$ $(\eta\, v_0 \gg s)$, the OAW undergoes collisionless damping linear in the frequency. The attenuation decrement is then

$$\Im k_0 = \gamma q \sqrt{\alpha^2 - 1}, \tag{25}$$

which is smaller than in an isotropic conductor (at $\alpha \gg 1$ it contains the factor $\eta\, v_0$). As indicated above, this occurs because the deformation interaction is small in a quasi-2D conductor. The attenuation decrement decreases radically when $\eta\, v_0$ approaches s. If $\eta\, v_0$ and s coincide, i.e for $\alpha = 1$,

$$k_0 = q\left(1 - \gamma + i\,\gamma\,\sqrt{\nu/\omega}\right), \qquad (|1-\alpha| \ll \nu^2/\omega^2 \ll 1), \tag{26}$$

and the attenuation decrement $\Gamma = \Im k_0$ approaches zero. The amplitude of the OAW has its maximum in this case because the condition for formation of standing OAWs in the film are fulfilled. At $\alpha \neq 1$ these conditions are destroyed and the amplitude of the OAW decreases rapidly (Fig. 1a).

Figure 1b shows the dependence of the amplitude of the AAW on the parameter α. It is obvious that the amplitude of the AAW has its maximum value also at $\alpha = 1$, as in this case the electron velocity v_z^{max} is along the wave vector which means that the interaction of these electrons with the electric field accompanying the wave is most effective.

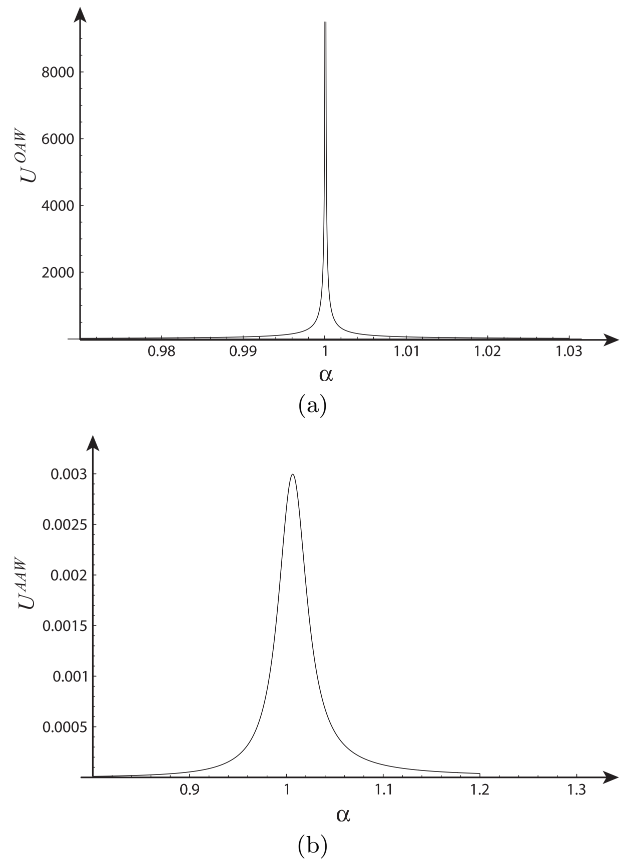


Fig. 1. The dependence of the amplitude of: a) OAW; b) AAW versus paramater $\alpha = \eta v_0/s\,(z = d/2, b = 0.05, g = 10^{-4}, \vartheta = \pi)$.

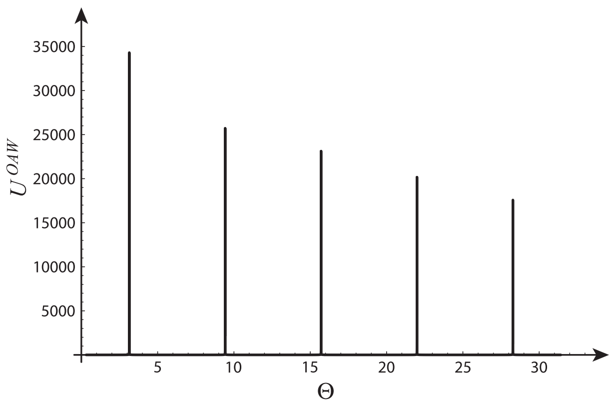
The frequency dependence of the amplitudes of the OAW and AAW at $\alpha = 1$ is shown in Figures 2a and 2b, respectively. In this case the frequencies determined by equation (24) are resonant for both waves but the amplitude of the OAW is much larger than that of AAW. With increasing α, the conditions for forming standing OAWs are violated but not the conditions for forming standing AAWs.

The amplitude of the OAW decreases with increasing frequency at $\alpha \gg 1$ since the attenuation decrement is proportional to ω (Eq. (25)). At $\alpha = 50$ and frequencies $\theta = \omega_n d/(\eta v_0) \geq 35\,\pi$ (i.e. $a = q\,d \geq 1750\,\pi$), the OAW is attenuated, and the AAW becomes dominant (Fig. 3a and Fig. 3b). For larger values of α the OAW is attenuated at even lower frequencies.

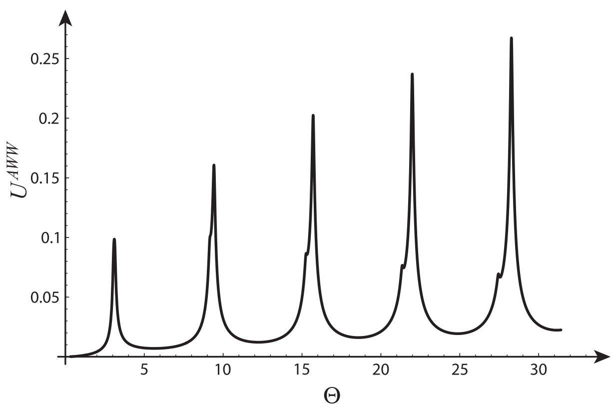
At $\alpha = 300$, the attenuation of the OAW occurs for frequencies $\theta \geq \pi$, i.e. $a \geq 300\,\pi$ (Fig. 4a and Fig. 4b).

When $\eta\, v_0 < s$, the OAW is attenuated by electron scattering. Then, the coefficient of attenuation does not depend on ω,

$$\Im k_0 = \gamma \frac{\nu}{s}\Big(\frac{1}{\sqrt{1-\alpha^2}} - 1\Big), \qquad (1 - \alpha \gg \nu^2/\omega^2). \tag{27}$$



(a)



(b)

Fig. 2. The frequency dependence of the amplitude of: a) OAW; b) AAW at $\alpha = 1\,(z = d/2, b = 0.05, g = 10^{-4})$.

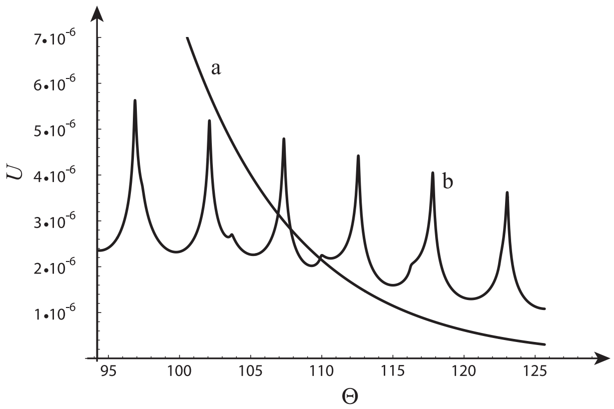


Fig. 3. The frequency dependence of the amplitude of: a) OAW; b) AAW at $\alpha = 50\,(z = d/2, b = 0.05, g = 10^{-4})$.

A purely two-dimensional electron gas ($\eta\, v_0/s \to 0$) would not interact with an acoustic wave propagating along the non-conducting axis and $k_0 = q$.

Comparing the nature of the attenuation of the ordinary and anomalous waves, one can easily determine that the asymptotic form of the acoustic field in a quasi-2D conductive film depends essentially on the value of α. There exists a "critical" value

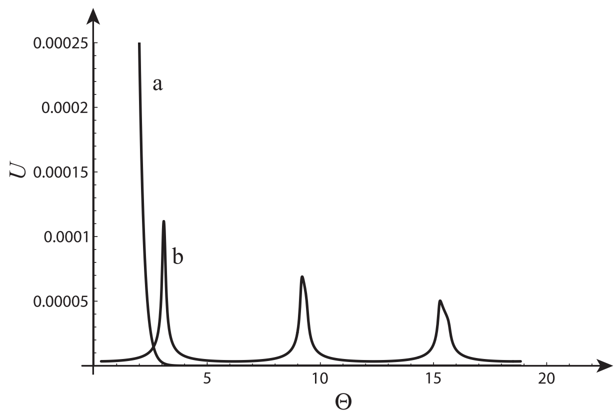
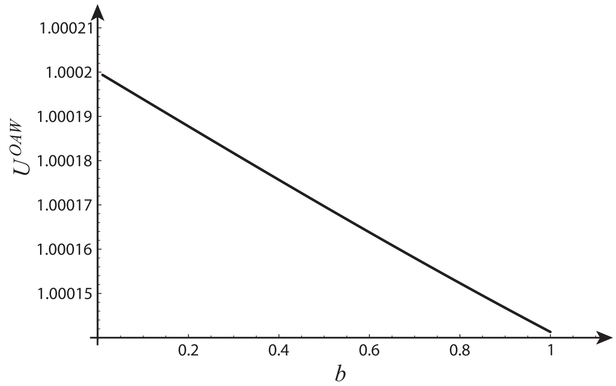
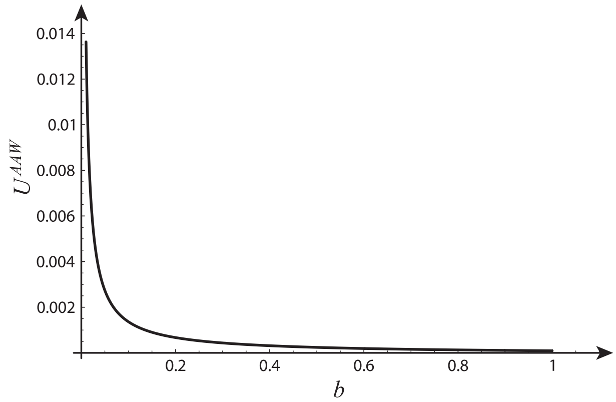


Fig. 4. The frequency dependence of the amplitude of: a) OAW; b) AAW at $\alpha = 300\,(z = d/2, b = 0.05, g = 10^{-4})$.



(a)



(b)

Fig. 5. The dependence of the amplitude of: a) OAW; b) AAW versus parameter $b = d/(nl)\,(z = d/2, \alpha = 2, g = 10^{-4}), \vartheta = \pi$.

$$\alpha_c = \sqrt{\frac{\nu}{\omega\gamma}}, \qquad (\alpha_c \gg 1) \tag{28}$$

below which the asymptotic form is like equation (20), but above which it is determined by the quasi-wave given by equation (22). These waves have substantially different velocities of propagation.

The corresponding dependence of both waves on the parameter $b = d/(\eta\ell)$ is shown in Figures 5a and 5b.

The amplitude of the OAW does not particularly depend on this parameter but the amplitude of the AAW shows a strong dependence on b. With increasing b it decreases rapidly and for $b \to 1$ (i.e. $d \to \eta\ell$) the condition for formation of standing AAWs is violated.

Formulas derivated above as well as the graphics obtained for different values of parameters show that the OAW in a quasi-2D conductive film in the absence of an external magnetic field is strongly attenuated, and the AAW is dominant, only for certain values of the characteristic parameters: for larger values of α and at high frequencies of the incident wave.

The investigation of AAWs is of interest because it opens the possibility of studying the electronic structure of quasi-2D conductors. For example, from the resonance conditions for anomalous standing waves, one can calculate the quasi-two-dimensionality parameter η, or one can obtain more information on the relaxation properties of the charge carriers.

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