Propagation of sound pulses in a metal with an open Fermi surface

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The effect of open electron orbits on the propagation of a modulated acoustic wave in a metal is investigated. A bulk sample and a thin metallic layer are considered under conditions of multichannel surface reflection of the carriers. The appearance of a system of secondary sound pulses is predicted. The magnitudes of the pulses and the delay (advance) time are found as functions of the applied magnetic field.

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1. Magnetoacoustic effects in a metal at $k^{-1} \ll r \ll l$ are sensitive to the geometry of its Fermi surface, i.e., their character is determined by the shape of the electron trajectories in the magnetic field. If these trajectories differ from the usual closed curves, the dependence of the sound wave damping on the magnetic field $\mathbf{H} \perp \mathbf{k}$ no longer reduces to the well-known Pippard oscillations,[1,2] and has a more complicated form.

Open (in the x-y plane perpendicular to H) electron orbits can result from multiple specular reflections of charges from the boundary of the sample, and, for electrons which do not interact with the boundary, in metals with an open Fermi surface (FS). Here **k** is the acoustic wave vector, and r and l are the characteristic Larmor radius and the free path length of the electron.

The features of the damping of a sinusoidal acoustic wave

$$U(x, t) = e^{i\xi}, \quad \xi = k(x - st)$$

under these conditions have been thoroughly studied,[3,4,5] and can be applied to the study of bulk and surface properties of conductors. However, in the high-frequency region ($\omega = ks \gtrsim 10^8\ \mathrm{s}^{-1}$) the pulse method of experiment turns out to be more advantageous (see, for example Ref. 6). The spatial extent of the acoustic pulse $\delta\xi/k$ is a new parameter, with the dimensions of length, and if it is smaller than the free path length of the electrons, the conditions for its interaction with the sound wave can change materially. In this connection, it is of interest to consider the passage, through a metal, of elastic waves in the form

$$U(\xi) = u(\xi) e^{i\xi}$$

with smooth envelopes localized in the interval $1 \ll \delta\xi \lesssim kr$.

Because of the great difference between the Fermi and the sound velocities $v_F \gg s$ the effective interaction of the electron with the deformation of the crystal takes place virtually simultaneously at several different depths x_i (according to the number of points $\mathbf{k} \cdot \mathbf{v} = ks$ at which it moves in the wavefront). Therefore, the sound pulse located at x perturbs the electron distribution function at all the points $x + x_i - x_{i'}$. As will be shown below, this leads to the appearance of a system of "satellites" of the fundamental signal; the structure of this system is determined by the geometry of the electron trajectories.

Thus, an open orbit in an unbounded metal contains an infinite set of points x_i, and the number of produced secondary signals is limited only by the ratio l/r.

In a thin conducting layer (thickness $d < l$) transitions between open and closed sections of the FS become possible in the case of reflection of the carriers by the boundaries of the metal. Thus groups of "glancing" electrons at both boundaries of the layer turn out to be connected even at $r \ll d$ and form a secondary pulse at a distance $\lesssim d$ from the fundamental one. Consequently, a "precursor" of the acoustic signal travels through the sample with a velocity close to the Fermi velocity.

2. We shall start out from the dynamical equations of elasticity theory

$$\rho(U_{tt}'' - s_0^2 U_{xx}'') = f\{U(x, t)\}, \tag{1}$$

the self similar solution of which $U(x, t) = U(\xi)$ is assumed to have the form $U(\xi) = u(\xi)e^{i\xi}$. Here ρ is the density of the crystal, s_0 is the unperturbed sound velocity (for definiteness, we take the sound wave to be longitudinal), and the force which the electrons exert on the lattice is[7]

$$f = \frac{\partial}{\partial \xi} \frac{2eHk}{ch^3} \int dp_z \int d\tau \Lambda(p_z, \tau) \chi(p_z, \tau, \xi). \tag{2}$$

For its calculation, it is necessary to solve for electrons in a metal a kinetic equation that can be represented in this case in the form

$$\left(\frac{\partial}{\partial \tau} + kV \frac{\partial}{\partial \xi} + \nu\right) \chi(p_z, \tau, \xi) = -k^2 s \Lambda(p_z, \tau) U''(\xi). \tag{3}$$

Here $(\partial F_0/\partial \varepsilon)\chi$ is the nonequilibrium addition to the electron distribution function, τ is the phase of the Larmor precession of the electrons and has the dimensionality of time, ν is their relaxation frequency, $V \equiv v_x - s$, and Λ is the xx component of the deformation potential.[1)] The complete solution of (3) for the electrons that do not collide with the boundaries of the sample is, as is well-known,

$$\chi(\tau, \xi) = -k^2 s \int_{-\infty}^{\tau} d\tau' e^{\nu(\tau' - \tau)} \Lambda(\tau') U'' \left(\xi - \int_{\tau'}^{\tau} kV\, d\tau''\right) \tag{4}$$

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and is essentially determined by the shape of the electron trajectory.

We shall be interested principally in the role of open trajectories. However, in a metal there are always closed electron orbits and as a rule make the principal contribution to the sound attenuation; therefore, we shall consider initially the case of a closed convex FS. In the most characteristic limit for the magneto-acoustic effects

$$ks<\nu\ll\Omega, \quad kr\gg 1 \tag{5}$$

we can neglect the phase shift of the sound field during the Larmor period $T=2\pi/\Omega$, and reduce the expression (4) to the form

$$\chi(\tau,\xi)\approx\frac{-k^2s}{1-e^{-\nu T}}\int_{\tau-T}^{\tau}d\tau' e^{\nu(\tau'-\tau)}\Lambda(\tau')U'''\left(\xi-\int_{\tau'}^{\tau}kv_x d\tau''\right)$$
$$\approx-\frac{k^2s}{\nu T}\sum_{i=1}^{2}\left(\frac{2\pi}{ik\dot v_x(\tau_{ei})}\right)^{1/2}e^{\nu(\tau_{ei}-\tau)}\Lambda(\tau_{ei})U''\left(\xi-\int_{\tau_{ei}}^{\tau}kv_x d\tau'\right). \tag{6}$$

The integral with respect to τ' is taken by the stationary-phase method $[v_x(\tau_{ei})=0]$ since $kr\gg1$, and $u(\xi)$ is a smooth function in comparison with $e^{i\xi}$. These same assumptions allow us to calculate the integrals over τ and p_z without difficulty (in the terms that oscillate with variation of p_z) and also in Eq. (2). It is easy to see that, because of the Larmor precession of the electrons, the result turns out to be dependent not only on ξ but also on the value of the argument, shifted by

$$\pm 2kr=\pm kD_{extr}=\pm\left(\int_{\tau_{e1}}^{\tau_{e2}}kv_x d\tau\right)\Bigg|_{p_z=0}. \tag{7}$$

Thus, in the passage of a sound pulse $U_0(\xi)=u_0(\xi)e^{i\xi}$ through a medium the conduction electrons exert the force

$$f(\xi)=\rho ks\Gamma_H\left[U_0'''(\xi)+\beta_0\sum_{\pm}e^{\mp i\pi/4}U_0'''(\xi\pm 2kr)\right] \tag{8}$$

on the crystal lattice. Here Γ_H is the monotonic part of the damping of a sinusoidal sound wave:

$$\Gamma_H\approx\frac{4\pi eHk}{ch^3\rho\nu T}\int dp_z\sum_{i=1}^{2}\frac{\Lambda^2(p_z,\tau_{ei})}{|\dot v_x(p_z,\tau_{ei})|};$$
$$\beta_0=\frac{4\pi eH}{ch^3\rho\nu T\Gamma_H}\left(\frac{2\pi k}{|D_{p_z}''(0)|}\right)^{1/2}\left|\prod_{i=1}^{2}\frac{\Lambda(0,\tau_{ei})}{(\dot v_x(0,\tau_{ei}))^{1/2}}\right| \tag{9}$$

is a nondimensional coefficient of the order of $(kr)^{-1/2}\ll 1$.

Knowing the form of the right-hand side of the dynamic equation, we can find with its help the envelope of the resulting signal $u(\xi)$ its propagation velocity, and the damping decrement γ. Thus, for an even "bare" envelope $u_0(\xi)$ we get, in first order in β_0,[2)]

$$u(\xi)=\frac{\Gamma_H}{\gamma}\left[u_0(\xi)+\beta_0\cos\left(2kr-\frac{\pi}{4}\right)\sum_{\pm}u_0(\xi\pm 2kr)\right],$$
$$s-s_0=0, \tag{10}$$
$$\frac{\gamma}{\Gamma_H}\approx 1+2\beta_0\cos\left(2kr-\frac{\pi}{4}\right)\frac{u_0(2kr)}{u_0(0)}.$$

Formula (10) describes the appearance of leading and lagging (by $\Delta t=2r/s_0$) "satellites" of the fundamental signal, which duplicate its shape in a scale that oscillates with change in the magnetic field. This has been observed experimentally.[6] For sound signals with a Gaussian envelope ($u_0(\xi)\propto\exp[-(\xi/\delta\xi)^2]$) this result was obtained in Ref. 8 by a Fourier-transform method.

3. For a given orientation of H let there now be in the metal open electron trajectories and let the direction of drift $\langle\mathbf{v}\rangle$ make with the direction of sound propagation (the x axis) an angle θ such that there exists at least one pair of points $v_x(\tau_{ej})=s$, that repeat periodically and are spaced cG/eH apart (G is the period of the open FS).

It is clear from the above that the energy of the initial sound pulse, which was located at depth x, should be distributed among regions

$$x\pm n\Delta, \quad \Delta=(cG/eH)\cos\theta, \quad n=0,1,2,3,\ldots,$$

lying in the limits of the electron mean free path.

This conclusion is not critical for a specific shape of the FS, and we only assume that the axis of the openness, together with the direction of H, forms a reflection plane in the reciprocal lattice of the crystal.

At depths exceeding l, expression (4) is valid as before for the distribution function of the electrons drifting in the chosen direction. This expression, upon satisfaction of the inequality (5), can be transformed into

$$\chi_+(\tau,\xi)=-k^2s\sum_{n=0}^{\infty}\int_{\tau-\tilde T}^{\tau}d\tau'\Lambda(\tau')\exp\{\nu(\tau'-\tau-n\tilde T)\}U''\left(\xi-nk\Delta-\int_{\tau'}^{\tau}kv_x d\tau''\right),$$

where $\tilde T$ is the period of motion of the electron along the open trajectory $v_x(\tau+\tilde T)=v_x(\tau)$, $\Lambda(\tau+\tilde T)=\Lambda(\tau)$. The usual stationary-phase method yields

$$\chi_+(\tau,\xi)\sim-k^2s\sum_j\frac{\Lambda(\tau_{ej})(2\pi)^{1/2}}{[ik\dot v_x(\tau_{ej})]^{1/2}}$$
$$\times\sum_{n=0}^{\infty}\exp\{\nu(\tau_{ej}-\tau-n\tilde T)\}U''\left(\xi-nk\Delta-\int_{\tau_{ej}}^{\tau}kv_x d\tau'\right). \tag{11}$$

In the calculation of the force f_+ [according to formula (2)], it must be kept in mind that the quantity Δ is constant over the entire layer of the open section of the FS, while the integrals $\int_{\tau_{ej}}^{\tau_{ej'}}kv_x d\tau$ (at $j\neq j'$) are functions of p_z, so that the corresponding "cross" terms, after integration with respect to p_z, turn out to be small, to the extent that $(k\Delta)^{-1/2}\ll 1$. Neglecting them, we obtain

$$f_+\approx-\frac{4\pi eHk^2s}{ch^3}\int_{(\mathrm{open})}dp_z\sum_j\frac{\Lambda^2(\tau_{ej})}{|\dot v_x(\tau_{ej})|}\sum_{n=0}^{\infty}e^{-n\nu\tilde T}U'''(\xi-nk\Delta). \tag{12}$$

Allowance for the boundary of the sample ($x=0$) shows that at a depth x, the number of terms in the sum over n is equal to $E(x/\Delta)$, which is the integer part of the ratio x/Δ. However, at depths $x>l$ this fact can be disregarded since terms with $n\gtrsim l/\Delta$ are exponentially small.

By virtue of the assumed symmetry of the FS, exactly the same form is possessed by the contribution f_-, which is dependent on $\xi+nk\Delta$, of electrons drifting out of the interior of the metal. In lower-symmetry cases, the coefficients of terms with $\xi + nk\Delta$ and $\xi - nk\Delta$ can of course be different.

As pointed out above, there are always closed electron orbits in real metals, in addition to the open orbits. The role of the corresponding asynchronous (i.e., depending on the shifted values of the argument $\xi \pm 2kr$) terms was explained in Sec. 2, and we exclude them from consideration here, assuming, for example, that the central section of the FS is open.[3)] However, as is seen from a comparison of formulas (12) and (10), the synchronous term in the force is generally speaking determined by these closed sections of the FS.

Summing up the above, we obtain the following expression for the force:

$$f(\xi) \approx -\rho k s \Gamma_{H1} \left[U_0'''(\xi) + \beta_1 \sum_{\substack{n=-\infty \\ (n\neq 0)}}^{+\infty} e^{-|n|\delta} U_0'''(\xi - nk\Delta) \right],$$

$$\Gamma_{H1} = \Gamma_H + \frac{8\pi e H k}{\rho c h^3} \int_{(\text{open})} dp_z \sum_j \Lambda^2(\tau_{ej}) / |\dot{v}_x(\tau_{ej})|, \qquad (13)$$

$$\beta_1 = \frac{4\pi e H k}{\rho c h^3 \Gamma_{H1}} \int_{(\text{open})} dp_z \sum_j \Lambda^2(\tau_{ej}) / |\dot{v}_x(\tau_{ej})|.$$

Here the integration is carried out over the layer of the open FS sections containing the points[4)] $v_x(\tau_{ej}) = 0$, and in Γ_H [see formula (9)] over values of p_z corresponding to closed orbits. If the structure of the electron spectrum does not contain additional small parameters,[5)] the quantity $\beta_1 \sim \nu T \ll 1$ ($\delta = \langle \nu \bar{T} \rangle$ is the averaged probability of scattering of the electron within the period of motion over an open trajectory.

In analogs with Sec. 2, we get with the aid of the dynamical equation (1) the envelope of the resulting signal (in first order in β_1):

$$u(\xi) = \frac{\Gamma_{H1}}{\gamma} \left[u_0(\xi) + \beta_1 \sum_{\substack{n=-\infty \\ (n\neq 0)}}^{+\infty} e^{-|n|\delta} \cos(nk\Delta) u_0(\xi - nk\Delta) \right], \qquad (14)$$

where, for even $u_0(\xi)$,

$$\frac{\gamma}{\Gamma_{H1}} = 1 + 2\beta_1 \sum_{n=1}^{\infty} e^{-n\delta} \cos(nk\Delta) \frac{u_0(nk\Delta)}{u_0(0)}. \qquad (15)$$

Thus, in the presence of open electron trajectories, the profile of the envelope of the wave in the interior of the sample, the quantity $x > l$, represents the complete set of small [in the measure $\beta_1 |\cos(nk\Delta)|$], (but slowly decreasing with the number $|n|$) "satellites" located on both sides of the fundamental signal and spaced $k\Delta = (kcG/eH)\cos\theta$ apart up to $\xi \sim kl$ (see Fig. 1). The maximum observed time lag (lead) is $\delta t \sim l/s_0$ and can reach $10^{-6} - 10^{-5}$ sec in pure metals.

We note that this result satisfies a special "correspondence principle": for a sinusoidal wave [$u_0(\xi)$ = const], by carrying out the summation in (15), we obtain

$$\frac{\gamma}{\Gamma_{H1}} \approx 1 + 2\beta_1 e^{-\delta} \frac{\cos(k\Delta) - e^{-\delta}}{1 - 2e^{-\delta}\cos(k\Delta) + e^{-2\delta}}, \qquad (16)$$

i.e., "resonance oscillations" similar to those predicted in Ref. 3.

4. In conclusion, we analyze the passage of sound pulses through a metallic layer, the surfaces of which $x = 0$, d reflect electrons almost specularly. We shall be interested in the case of multi-channel reflection, when an electron reflected by a boundary of the metal with conservation of the tangential quasimomentum (p_y, p_z) can go over with a certain probability W into another cavity of the non-singly-connected FS. In a magnetic field parallel to the layer such processes modity substantially the "glancing" trajectories of the electrons, and this leads to a new type of magnetoacoustic oscillations.[5]

Consider, for example, a FS whose central section is shown schematically[6)] in Fig. 2a. Thanks to the presence of two reflection channels corresponding to the open and closed parts of the FS, electrons, glancing along one of the boundaries of the layer, become able from time to time to go over to another layer in an arbitrarily strong magnetic field ($r \ll d$). Such

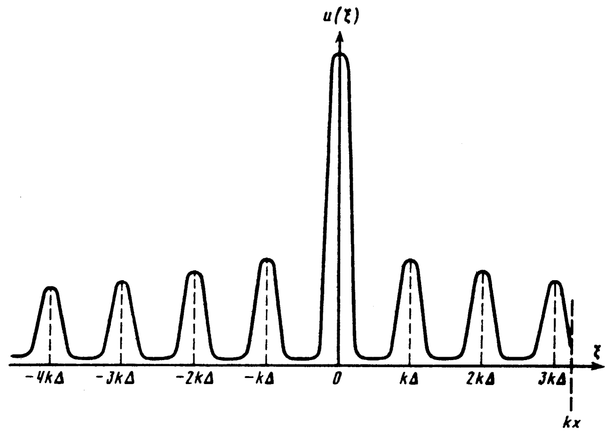

FIG. 1. Envelope of the wave in the interior of a metal with open FS. If the width of the initial signal is $\delta\xi \ll k\Delta$, then the secondary pulses (near $\xi = nk\Delta$) duplicate its shape in the scale $\beta_1 \exp(-|n|\delta)|\cos(nk\Delta)|$. At the depth x, the number of "precursors" ($n > 0$) does not exceed $E(x/\Delta)$.

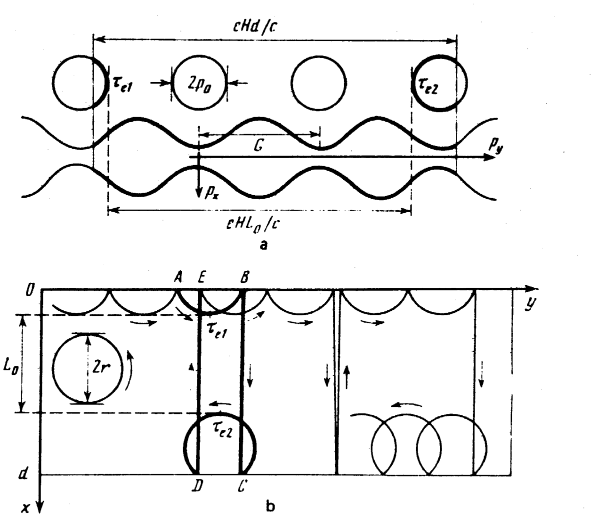

FIG. 2. a) G—period of the open sections of the FS and $2p_0 = 2eHr/c$ is the maximum diameter of the closed ones. For definiteness, $2p_0 < G/2$. b) Corresponding trajectories of electrons in the layer (the sections of the open trajectories are shown by straight lines), $\tau_{e1,2}$ are the turning points of the electrons colliding with the boundary of the metal.

electrons, interacting with the acoustic signal localized near $x=0$, transport part of their energy to the surface $x=d$ and form there a second burst of the acoustic field.

It is seen from Fig. 2 that the transition from an open orbit to a glancing one is possible in both surface layers at values of the magnetic field

$$n-2p/G<Z<n+2p/G, \quad Z=eHd/cG, \quad n=1,\ 2,\ 3,\ \dots \qquad (17)$$

Here the distance between turning points of the electrons, over which the signal is transported, is

$$L=\int_{\tau_{e1}}^{\tau_{e2}} v_x\,d\tau=\frac{d}{Z}\left[E\left(Z+\frac{1}{2}\right)-\frac{2p}{G}\right]. \qquad (18)$$

We point out the basic features of the calculation. The problem of finding the 4-component distribution function χ with the help of two pairs of boundary conditions turns out to be rather cumbersome. However, it can be simplified in the case

$$1>W\gg q,\ d/l, \qquad (19)$$

which is more favorable for the observation of the expected effect, when W greatly exceeds both the fraction q of the diffusely reflected electrons and the probability of their bulk scattering between two reflections by the boundary ($\lesssim d/l$). Within the free path time the electron manages to pass repeatedly through all four possible parts of the trajectory and, as exact calculations confirm, the transition probability W does not enter at all in the result. Setting now $W\approx 1$, we see that the time T of traversing the figure $ABCDE$ (Fig. 2b) turns out to be the effective period of the motion, which allows us to represent χ in the form

$$\chi=\frac{-k^2s}{1-e^{-\nu\bar{T}}}\int_{(ABCDE)} d\tau' e^{-\nu(\tau-\tau')}\Lambda(\tau')U''\left(\xi-\int_{\tau'}^{\tau} kv_x\,d\tau''\right). \qquad (20)$$

This integral differs in practice from zero only in those regions of space where the points τ_{e1} and τ_{e2} lie:

$$\chi\approx-\frac{k^2s}{\nu\bar{T}}\left(\frac{2\pi}{|\dot{v}_{xe}|k}\right)^{1/2}\Lambda_e\left\{\theta(x)e^{-i\pi/4}U''\left(\xi-\int_{\tau_{e2}}^{\tau} kv_x\,d\tau'\right)\right.$$
$$\left.+\theta(x+L)e^{i\pi/4}U''\left(\xi-\int_{\tau_{e1}} kv_x\,d\tau'\right)\right\}. \qquad (21)$$

where the function $\theta(x)=1$ at

$$\max\{0;\ d-D-L\}\leqslant x\leqslant\min\{D;\ d-L\}$$

and is equal to zero for other x.

The calculation of the force f is similar to that described in Secs. 2 and 3. The electrons on the considered orbits have asynchronous terms with arguments $\xi\pm kL_0$ while the synchronous part (which depends on ξ) of the force is determined as before by the closed section of the FS, since the probability of scattering by the bulk electrons during the period νT of their motion turns out to be much smaller than $\nu\bar{T}$:

$$T/\bar{T}\approx T|\langle v_x\rangle|/2d\sim r/d\ll 1$$

($\langle v_x\rangle$ is the drift velocity of the electrons).

As a result, we get for the envelope of the electron force

$$f_0=fe^{-i\xi}/i\rho ks\Gamma_H \qquad (22)$$

the result

$$f_0\approx u(\xi)+\begin{cases}\beta u(\xi+kL_0), & \text{a) } 0\leqslant x\leqslant d-L_0,\\ 0, & \text{b) } d-L_0<x<L_0,\\ \beta^* u(\xi-kL_0), & \text{c) } L_0\leqslant x\leqslant d,\end{cases} \qquad (23)$$

where

$$\beta=\beta_2\exp(ikL_0+i\pi/4);$$
$$\beta_2\approx\frac{4\pi eH}{ch^3\rho\nu\bar{T}\Gamma_H}\left(\frac{2\pi k}{|D_{p_z}''|_0}\right)^{1/2}\frac{\Lambda_{e0}^2}{|\dot{v}_{xe0}|}\sim\left(\frac{G}{p_0kdZ}\right)^{1/2},$$

and the index 0 denotes the central section of the FS. The regions of existence of the asynchronous terms are shown for the case

$$0\leqslant E(Z+1/2)-Z\leqslant 2p_0/G,$$

when they are directly adjacent to the layer boundary. We shall not write out the asynchronous terms with arguments $\xi\pm 2kr$ (see Sec. 2) since we shall be interested in the transport of the sound pulses to a large distance $L_0\gg 2r$.

Let us consider this process. In the zeroth order in $\beta_2\ll 1$, the solution of the dynamic equation

$$\frac{k(s^2-s_0^2)}{is\Gamma_H}u=\frac{\gamma}{\Gamma_H}u=f_0 \qquad (24)$$

is the function $u(i\xi)$, $\xi=k(x-s_0t)\equiv kx-\omega t$ which satisfies the boundary condition $u(0,t)=u_0(-\omega t)$, i.e., $u_0(\xi)$.

In first approximation in β_2, the solution $u(x,t)$ is self-similar, strictly speaking, only within the regions indicated in (22), and the values of the coefficient γ in them turn out to be different:

$$\text{a) } \gamma=\gamma'=[1+\beta u_0(kL_0)/u_0(0)]\Gamma_H,$$
$$\text{b) } \gamma=\Gamma_H,$$
$$\text{c) } \gamma=\gamma''=[1+\beta^* u_0(-kL_0)/u_0(0)]\Gamma_H.$$

The functions $u(L_0,t)$ and $u(0,t-L_0/s_0)$ [in regions b) and a), respectively] are not identical, but are connected by a relation which follows from formulas (23) and (24) and from the boundary condition

$$\gamma' u(L_0,\ t)-\gamma'' u(0,\ t-L_0/s_0)=\gamma' u(L_0,\ t)-\gamma'' u_0(kL_0-\omega t)$$
$$=\Gamma_H[\beta^* u_0(-\omega t)-\beta u_0(2kL_0-\omega t)].$$

Now shifting t by $(d-L_0)/s_0$ and neglecting quantities $\sim\beta_2^2$, we find the envelope of the signal at the "exit" ($x=d$)

$$u(d,\ t)\approx\mathrm{Re}[(\gamma''/\gamma')u_0(kd-\omega t)$$
$$+\beta^* u_0(kd-\omega t-kL_0)-\beta u_0(kd-\omega t+kL_0)]. \qquad (25)$$

Thus, in addition to the main signal (at the instant $t=d/s_0$) a "precursor" and a "follower" should be observed at the second boundary of the sample at the instants $t=(d\mp L_0)/s_0$. They are of the same shape as the main signal but are decreased by $\beta_2|\cos(kL_0+\pi/4)|$. The lead (lag) interval is

$$\frac{L_0}{s_0}=\frac{d}{s_0Z}\left[E\left(Z+\frac{1}{2}\right)-\frac{2p_0}{G}\right]. \qquad (26)$$

and approaches in a strong magnetic field, when $Z\equiv eHd/cG\gg 1$, the time at which the main signal passes through the sample.

In conclusion, we calculate the damping of the harmonic wave [$u_0(\xi)=\text{const}$] passing through a metallic layer under the considered conditions:

$$\gamma d=(\gamma'+\gamma'')(d-L_0)+\Gamma_H(2L_0-d)=\Gamma_H[d+(d-L_0)\cdot 2\beta_2\cos(kL_0+\pi/4)],$$

which can be represented in the form

$$\gamma(Z)/\Gamma_H \approx 1+2\beta_2 R \cos(kdR+\psi), \qquad (27)$$
$$R(Z) = 1-Z^{-1}[E(Z+{}^1/_2)-2p_0/G].$$

Thus, in addition to the Pippard oscillations, there are specific damping oscillations that are modulated by the sawtooth function $R(Z)$.

5. For experimental observation of a series of N precursors (or followers, Sec. 3), the receiving apparatus should resolve sound pulses separated by the interval $\delta t \lesssim l/Ns_0$. Therefore the metals with open FS, in which large carrier mean free paths are possible, are most convenient for the measurements, for example, Ga used in Ref. 6 and also Cd, Cu, Ag and Re (see Ref. 9). Realization of the size effect described in Sec. 4 requires of course, high quality of the surface of the sample.

A common feature of the particular cases considered is the transport of the sound pulses in the metal with velocities of the order of the Fermi velocity through distances that are characteristic of a system of electron orbits in an external magnetic field. Chief interest here is attached to trajectories that are infinite in the direction of propagation of the wave **k** (Sec. 3), or which possess anomalously large spatial periods, as in Sec. 4.[7)]

The first case can be realized also for a closed FS if the angle between **k** and **H** differs from a right angle. Here, too, an equidistant system of "satellites" of the fundamental signal arises at distances that are multiples of the extremal period of drift of the electrons Δ_{xe}. At $k\Delta_{xe} \gg 1$, their maximum relative value is

$$u(n\Delta_{xe})/u(0) \sim (k\Delta_{xe})^{-1/2} \exp(-|n|\Delta_{xe}/l).$$

In a thin layer of metal, in the case of certain orientations of the crystallographic axes, it may turn out that the different channels of surface reflection correspond simply to repetitions of the closed cavity of the FS in the broadened band scheme. Then (see Ref. 5) the shift D_x of the turning points of the electron after such a transition is the same for the entire layer p_z. This circumstance eliminates the smallness of the coefficient β, due to interference of the different $D_x(p_z)$, and can lead at $2r \gtrsim d$ to the formation of secondary pulses $u(\xi \pm kD_x)$ comparable in magnitude with the fundamental pulse.

These effects illustrate the new possibilities of pulsed magneto-acoustic experiments.

The calculation method used by us makes it possible to analyze the propagation of a smoothly modulated acoustic wave in a metal without specifying the shape of its envelope.

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[1)] In the general case, the right-hand side of the kinetic equation is equal to

$$\Lambda_{ik}\dot{U}_{ik} - e\mathbf{v}(\mathbf{E}+[\mathbf{v}\times\mathbf{H}]/c),$$

where $\dot{U}_{ik}$ is the time derivative of the distortion tensor; the total electric field **E** is calculated by means of Maxwell's equations. However, in the case of a longitudinal strain, the corresponding renormalization of Λ_{xx} is unimportant.

[2)] In the principal approximation in $\nu T \ll 1$, the renormalization of the signal velocity

$$(s-s_0)/s_0 \approx \beta_0 \sin(2kr-\pi/4)[u_0(2kr)-u_0(-2kr)]/u_0(0)$$

vanishes for even $u_0(\xi)$.

[3)] If there are nevertheless closed sections with extremal diameter D, these terms can be made sufficiently small by varying the parameter β_0.

[4)] The number of such points in one period of the motion can generally speaking, vary with p_z. However, this fact does not change the order of magnitude of β_1 and all the more of Γ_{H1}.

[5)] For example, these can be the smallness of the corrugation of the open FS, a significant difference in the values of Λ for different regions of the FS, etc.

[6)] For example, the FS of rhenium has such a shape.[9] For us, however, only the constancy of the sign of v_x on the open sections and the presence of a reflection plane $\{p_x, p_z\}$ are important, the latter for simplification of the formulas. The other details of the structure of the FS are unimportant here.

[7)] Such trajectories can arise also in a bulk sample under conditions of magnetic breakdown between the open and close cross sections of the FS.

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