[4] V. I. Talanov, JETP Letters 2, 218 (1965), transl. p. 138.
[5] S. A. Akhmanov, A. P. Sukhorukov, and R. V. Khokhlov, JETP 50, 1537 (1966), Soviet Phys. JETP 23, in press.
[6] N. F. Pilipetskii and A. R. Rustamov, JETP Letters 2, 88 (1965), transl. p. 55.
[7] Ya. B. Zel'dovich and Yu. P. Raizer, ibid. 3, 137 (1966), transl. p. 86.
[8] E. Garmire, R. Y. Chiao, and C. A. Townes, Phys. Rev. Lett. 16, 347 (1966).
[9] P. L. Kelley, ibid. 15, 1005 (1965).

1) When $\theta_{max} \ll 1$ the presence of the medium does not change r_f.

2) In considering the channel itself ($d\theta/dx = 0$) it is necessary to make Eqs. (1) and (2) more precise by writing in (1) $\delta\epsilon/2\epsilon_0$ in place of n_2E^2/n_0 and in (2) $\sqrt{\epsilon_0 + \delta\epsilon}$ in place of n_0 ($\epsilon = \epsilon_0 + \delta\epsilon$ is the dielectric constant). In the channel we obtain approximately: $\delta\epsilon(E^2) = \lambda_0^2/4\pi^2r_k^2$ and $P_{ch} = \sqrt{\epsilon_0 + \delta\epsilon}(cE^2/4\pi)\pi r_k^2(e - 1)$, where r_k is the effective radius of the channel. We eliminate the field from these equations by using the interpolation formula $\delta\epsilon = \epsilon_2E^2[1 + \epsilon_2E^2/\delta\epsilon_{max}]^{-1}$, which ensures "saturation" of the polarizability [7] ($\epsilon_2 = 2n_0n_2$, $\delta\epsilon_{max} \sim \epsilon_0$). We then obtain $P_{ch}/P_{cr} = (1 + \delta\epsilon)(1 - \delta\epsilon/\delta\epsilon_{max})^{-1}$, which together with the first of the equations gives the dependence of the radius on the power. When $P_{ch} \approx P_{cr}$ we have $r_k \sim \lambda_0[(P_{ch}/P_{cr}) - 1]^{-\frac{1}{2}}$, as in [3], and when P_{ch} exceeds P_{cr} by several times r_k is of the order of the wavelength and does not depend on P_{ch}.

3) Private communication from N. F. Pilipetskii.

OSCILLATIONS OF THE INTERMEDIATE STATE OF SUPERCONDUCTORS

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In the intermediate state, the volume of a superconductor constitutes a system of alternating layers of normal and superconducting phases. The regular arrangement of the layers corresponds to the minimum of energy, and therefore if the layers are bent by external forces, the system will ultimately return to the unperturbed state. This raises the question of whether this process will constitute weakly damped oscillations. This question is closely connected with the question of the motion of macroscopic filaments of magnetic flux, which was already discussed earlier [1,2].

In this paper we consider small oscillations of the intermediate state, the wavelength of which greatly exceeds the dimensions of the layers of the normal phase.

The motion of the interphase boundaries is accompanied by occurrence of an alternating electromagnetic field in the normal regions. This field is described by the equations

$$\text{curl } \vec{H} = \frac{4\pi}{c}\vec{j}, \text{ and curl } \vec{E} = \frac{i\omega}{c}\vec{H}, \tag{1}$$

where ω is the oscillation frequency.

It is necessary to satisfy on the interphase boundaries the conditions of continuity of the tangential component of the vector $\vec{E}$ and of the normal component of $\vec{H}$ in the coordinate frame fixed on the boundary [3]:

$$H_x = ik_z H_c \zeta; \quad E_y = -\frac{i\omega}{c} H_c \zeta; \quad E_z = 0 \quad (x = \pm a_n/2), \tag{2}$$

where ζ is the x component of the displacement of the boundary, k is the two-dimensional oscillation vector ($k_x = 0$), a_n is the thickness of the normal layer, and H_c is the critical magnetic field. The coordinate system is chosen such that the x axis is normal to the unperturbed separation boundary, the yz plane lies in the middle of the normal layer, and the z axis is directed along the magnetic field $\vec{H}_c$ in the absence of oscillations. The coordinate dependence of all the quantities is of the form $f(x)\cdot\exp(i\vec{k}\cdot\vec{r})$.

In order to obtain the complete system of equations, we must find, first, the connection between the current $\vec{j}$ and the electric field $\vec{E}$, and second, the boundary conditions for H_z. The latter is easily done by using the equality of the normal forces on both sides of the boundary

$$\frac{H_z^2}{8\pi} = \frac{H_c^2}{8\pi} + F, \quad x = \pm a_n/2, \tag{3}$$

where F is the force exerted on the boundary by the electrons. In the right side of (3) it is necessary to write out also a term connected with the surface tension. For long waves, however, this term is immaterial.

The current density $\vec{j}$ should be determined by solving the kinetic equation for the electron distribution function $f = f_0 + (\partial f/\partial E)\chi$ (f_0 is the equilibrium value):

$$v_x \frac{\partial \chi}{\partial x} + \Omega \frac{\partial \chi}{\partial \tau} + I\{\chi\} = -e\vec{v}\cdot\vec{E}, \tag{4}$$

where $\vec{v}$ is the electron velocity, Ω the cyclotron frequency in the field H_c, and τ the dimensionless time of revolution (see the paper of Lifshitz et al. [4]). $I\{\chi\}$ is the collision integral. We have neglected in (4) the terms containing the derivatives with respect to the time and with respect to y and z, assuming that the following conditions

$$ka_n \ll 1, \quad \omega \ll \nu,$$

where ν is the electron collision frequency, are satisfied.

We assume also that the temperature is low compared with the critical temperature of the superconducting transition in the absence of a magnetic field. In this case it is easy to find the boundary condition for the function χ for $x = \pm a_n/2$ (see [5]):

$$\chi(\vec{v}) + \chi(-\vec{v}) = 0. \tag{5}$$

We took account here of the fact that the change in the momentum of the electronic excitations upon reflection from the boundary is proportional to the distance from the Fermi surface, and is consequently small at low temperatures. For the same reason, we can put $F = 0$ in (3).

In the considered case of long waves, we can neglect in Eq. (4) the dependence of $\vec{E}$ on the coordinate x. Then this equation has a solution that does not depend on x and coincides with the solution corresponding to an infinite normal metal. Using the symmetry of the collision integral with respect to the transformation $\vec{v} \to -\vec{v}$, we can readily verify that this solution satisfies the boundary condition (5) automatically. We thus arrive at the conclusion that the dependence of the current on the electric field is determined by the static conductivity of the bulk metal in a magnetic field H_c, namely $\vec{j} = \hat{\sigma}\vec{E}$.

Using (1) and boundary conditions (2) and (3), we can readily obtain the following relations [1)]:

$$\frac{ck_z^2}{4\pi} H_c\zeta + j_y = 0; \quad E_x = \frac{4\pi i\omega}{c^2k_z^2} j_x; \quad E_y = -\frac{i\omega}{c} H_c\zeta; \quad E_z = 0, \tag{6}$$

which make it possible to find the function $\omega(\vec{k})$ if the conductivity $\hat{\sigma}$ is specified.

If $\Omega \ll \nu$, the conductivity in the absence of a magnetic field can be used. It is then easy to verify that the oscillations attenuate over a distance of the order of their wavelength.

In the opposite extreme case ($\Omega \gg \nu$), using the known values of $\hat{\sigma}$ [4], we find

$$\omega = \frac{cH_c}{4\pi Ne} k_z^2, \tag{7}$$

where e is the electron charge and N the difference between the number of electrons and holes.

It is assumed here that there are no open trajectories and $N \neq 0$. When $N = 0$ the oscillations are strongly damped.

Thus, weakly damped oscillations of the interphase boundaries exist in the intermediate state of a superconductor with unequal number of electrons and holes when $\Omega \gg \nu$. The spectrum (7) of these oscillations is very similar to the spectrum of helicoidal waves (see [6]). We emphasize, however, that the oscillations (7) can exist also when $kr \gg 1$ (r is the Larmor radius).

The presence of weakly damped oscillations of the layers was recently observed experimentally in indium [7]. However, the condition $T \ll T_c$ was not satisfied in these experiments, and therefore the experimentally observed values of ω are lower than those given by (7). The experimental data of [7] approach the theoretical values with decreasing temperature.

It is possible to solve in similar manner the previously considered problem [2] of the motion of a curved filament of the normal phase. We obtain in this case

$$V_y = \frac{c^2}{4\pi R} \frac{\sigma_{xy}}{\sigma_{xx}\sigma_{yy} - \sigma_{xy}\sigma_{yx}}, \quad V_x = \frac{c^2}{4\pi R} \frac{\sigma_{xx}}{\sigma_{xx}\sigma_{yy} - \sigma_{xy}\sigma_{yx}}, \tag{8}$$

where V_y and V_x are the components of the velocity of a line element of the filament perpendicular and parallel to the curvature plane, σ_{ik} the conductivity of the bulk normal metal in a field H_c directed along the z axis, and R the radius of curvature. If $\Omega \gg \nu$ and $N \neq 0$, then relation (8) yields $V_y = cH_c/4\pi NeR$ and $V_x \ll V_y$ (cf [2]). We note that the direction of

V_y depends on the sign of N. When N = 0 we always have $V_y \sim V_x$. This explains why the motion of the normal regions was observed in [8] in indium ($N \neq 0$) but not in tin (N = 0).

[1] J. Bardeen, Phys. Rev. Lett. 13, 747 (1964).

[2] P. G. De Gennes and P. Nozieres, Phys. Lett. 15, 216 (1965).

[3] L. D. Landau and E. M. Lifshitz, Elektrodinamika sploshnykh sred (Electrodynamics of Continuous Media), Fizmatgiz, 1959, p. 228

[4] I. M. Lifshitz, M. Ya. Azbel', and M. I. Kaganov, JETP 31, 63 (1956), Soviet Phys. JETP 4, 41 (1957).

[5] A. F. Andreev, JETP 47, 2222 (1964), Soviet Phys. JETP 20, 1490 (1965).

[6] E. A. Kaner and V. G. Skobov, JETP 46, 1106 (1964), Soviet Phys. JETP 19, 749 (1964).

[7] B. W. Maxfield and E. F. Johnson, Phys. Rev. Lett. 15, 677 (1965).

[8] W. De Sorbo, Philos. Mag. 11, 853 (1965).

1) For long waves $\zeta(x = a_n/2) \approx \zeta(x = -a_n/2)$.

PRESSURE DEPENDENCE OF ELECTRON EFFECTIVE MASS IN INDIUM ANTIMONIDE

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By changing the width of the forbidden band in a crystal, hydrostatic compression affects the value of the effective mass of the carriers. If the band is non-parabolic, as is the conduction band in InSb, the effective mass varies with the filling of the band, and the pressure can have diverse influences on the effective mass.

According to Kane's theory, the electron effective mass m_n at the bottom of the conduction band is given by the formula

$$\frac{m_0}{m_n} = 1 + \frac{\epsilon_p}{3}\left(\frac{2}{\epsilon_g} + \frac{1}{\epsilon_g + \Delta}\right) , \qquad (1)$$

where $\epsilon_g = \epsilon(\Gamma_{1c}) - \epsilon(\Gamma_{15v})$ is the width of the forbidden band at the Γ-point, Δ the energy of spin-orbit splitting of the valence band at the Γ-point, $\epsilon_p = (2m_0/\hbar^2)\mathcal{P}^2$, $\mathcal{P}$ the matrix element of the interaction between the states Γ_{1c} (conduction band) and Γ_{15v} (valence band), and m_0 the mass of the free electron.

Under the usual assumption that ϵ_p is not affected by small changes of the lattice period, the mass m_n is almost proportional, according to (1), to ϵ_g, which in InSb increases linearly with the pressure P with a coefficient $d\epsilon_g/dP = 1.6 \times 10^{-5}$ eV/atm.

We undertook an experimental study of the influence of hydrostatic pressure up to 16.5 katm on the effective mass m_n of the electrons in InSb at T = 96°K. To this end, we measured the thermal emf and the Hall effect in classically strong magnetic fields, when neither de-