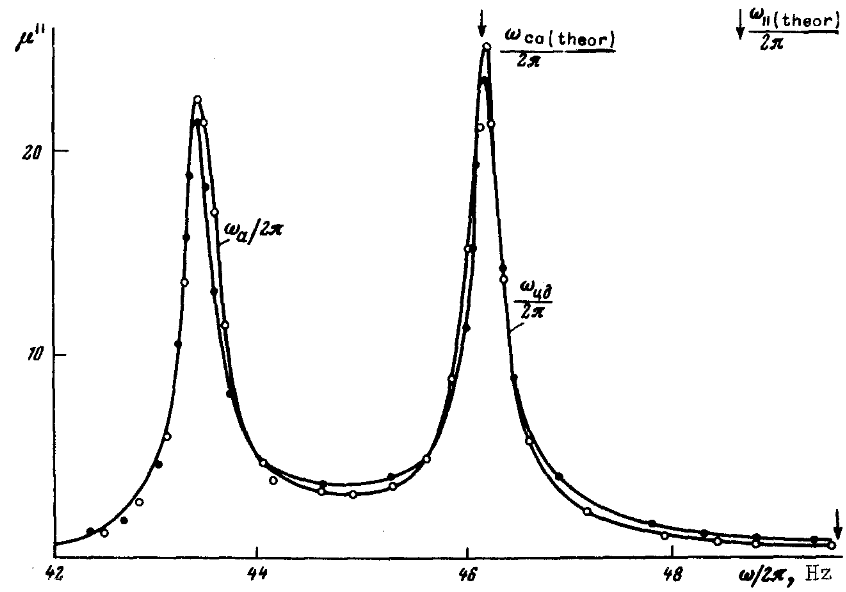

Magnetic spectra of magnetoplumbite crystal; ● – $h \perp H^{!}_{sat}$, o – $h \parallel H^{!}_{sat}$.

(H_a = 13.75 kOe [5]).

The high-frequency peak represents resonant absorption connected with the precession of the cylindrical-domain magnetization. The arrow in the figure indicates the theoretical frequency $\omega_{c.d.}$ calculated from [3], where H_a = 15 kOe and M = 320 G. The agreement with the experimental frequency is very good. Finahly, this statement is also favored by the fact that the optimal conditions for the excitation of the most intense and narrowest maximum $\omega_{c.d}(\theta = 87^{\circ})$ coincide with the optimal conditions for the nucleation of a cylindrical domain structure.

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TEMPERATURE DEPENDENCE OF THE CONDUCTIVITY OF A THIN PLATE

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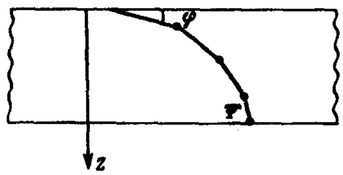
The dependence of the resistance of a plate of arbitrary thickness on the thickness and temperature is found, with allowance for electron diffusion (due to collisions with phonons) and for the arbitrary character of the electron reflection from the surfaces of the plate.

There have been many experimental studies of the resistance of thin metallic samples as functions of the temperature and dimensions (see [1], and also the review [2]). A consistent theoretical study of this question reduces to an analysis of the electron diffusion due to collisions with phonons between the walls of the sample, the reflection from which depends strongly on the angle of incidence on the wall. Nevertheless, the presently available theoretical papers deal with cases in which the sample is so thin that there is no diffusion [3-7].

We report here the first study of a general case of a plate of arbitrary thickness, and show that the form of the dependence of the resistance on the temperature varies significantly



with the sample thickness and with the character of the reflection from the surface (formulas (1), (4 - 6)). We explain the results by using simple qualitative arguments.

The electron executes Brownian motion in angle space (see, e.g., [8, 9]). The angle at which the electron moves after the n-th collision with a phonon is therefore $\phi_n \sim \phi + \alpha\sqrt{n}$, where ϕ is the initial angle (see the figure) and $\alpha = T/\theta$ is the characteristic angle of rotation after a single collision (T is the temperature and θ is the Debye temperature). Since it will be shown that only small angles matter, it suffices to take into account only the departure of the electron as a result of diffusion from this angle region.

The calculation of the mean free path is different for electrons whose angles relative to the surface do not exceed $\chi = d/\ell_T$ along their entire path, and for electrons for which $\phi > \chi$. Here d is the thickness of the plate and ℓ_T is the length of the "step" of the electron in the diffusion, i.e., the path between two successive electron-phonon collisions: $\ell_T = \ell_\theta \alpha^{-3}$, $\ell_\theta = hv_F/\theta$, h is Planck's constant, and v_F is the Fermi velocity.

In the first case (ϕ, $\phi_n < \chi$) the electron path between collisions n and n + 1 with the phonons in the z direction (see the figure) is $z_n \sim \ell_T \phi_n$. The number of collisions in one turn (i.e., on the path of the electron from one surface to the other) can be obtained from the fact that the sum of the values of z_n in one "turn is a quantity nearly equal to d, namely, for the (k + 1)-st "turn" we have

$$\sum_{n=n_k}^{n_{k+1}} z_n \sim \ell_T (\phi + \alpha\sqrt{n_k})(n_{k+1} - n_k) \sim d.$$

This enables us to determine the function n_k, i.e., the number of "steps" after k "turns":

$$n_k \sim \left[\frac{\phi}{\chi k} + \left(\frac{\alpha}{\chi k}\right)^{2/3}\right]^{-1}$$

The effective electron path, on which it contributes to the current density, terminates when the probability of diffuse reflection becomes of the order of unity, i.e.,

$$\sum_{k=1}^{k_{tot}} \{1 - q(\phi_{n_k})\} \sim 1$$

($q(\psi)$ is the coefficient of specular reflection). This enables us to determine the total number of turns along this path, meaning also the total number n_{tot} of "steps" on this path.

In the second case ($\phi > \chi$) one "step" spans several "turns." The electron flight angle does not change in one "step," and thus $\ell_T \phi_n / d$ collisions take place in one "step" (n is the number of the "step"). The total number of electron-phonon collisions over the effective mean free path is therefore determined from the equation

$$\sum_{n=0}^{n_{tot}} (\ell_T \phi_n / d)(1 - q(\phi_n)) \sim 1 .$$

The value of n_{tot} determines the effective path length $\sim \ell_T n_{tot}(\phi)$ of an electron leaving the surface at an angle ϕ. Of course, in any case the maximum possible path length cannot exceed the usual mean free path in the bulky metal, governed by either the phonons or by the impurities: $\ell_{bulk}^{-1} \sim \ell_{ep}^{-1} + \ell_i^{-1} \sim (\ell_T/\alpha^2)^{-1} + \ell_i^{-1}$ (ℓ_i is the free path due to impurities). The total effective electron mean free path is therefore of the order of $\ell_{eff}^{-1}(\phi) \sim (\ell_T n_{tot}(\phi))^{-1} + \ell_{bulk}^{-1}$, and the effective electric conductivity is given by

$$\sigma \sim \frac{Ne^2}{p_F} \int_0^{\pi/2} \ell_{eff}(\phi)\, d\phi ,$$

where N is the charge density, p_F is the Fermi momentum, and e is the electron charge.

For $d < \ell_{ep}$ we obtain, in order of magnitude

$$\sigma \sim \frac{Ne^2 d}{p_F} \left\{ \ln \frac{1}{\Phi + \Psi} + \frac{1}{\chi\, \Phi^{-1}[1 + (\Phi/\alpha)^2]^{-1} + \kappa \Psi^{-1}} \right\}, \qquad (1)$$

where Φ and Ψ are determined from the relations

$$[1+(\Phi/a)^2]\,\Phi(1-q(\Phi)) = \chi, \quad \chi = d/\ell_T;$$

$$\Psi(1-q(\Psi)) = \kappa, \quad \kappa = \frac{d}{\ell_i} + \frac{d}{\ell_{ep}}, \quad a = \frac{T}{\theta}. \tag{2}$$

The obtained formula shows, in particular, that the temperature at which the main role is assumed by scattering by phonons corresponds to $\ell_T < \ell_i < \ell_{ep}$ and, generally speaking, has a complicated dependence on the thickness, being given by the equation

$$\ell_T\,\Phi[1+(\Phi/a)^2] \sim \ell_i\,\Psi. \tag{3}$$

In limiting cases, formula (1) can be readily simplified. Thus, if scattering by impurities can be neglected, then we have at $\ell_{ep} > d > \ell_T T/\theta$

$$\sigma \sim \frac{Ne^2 d}{p_F}\left\{\ln\frac{1}{\tilde{\Phi}} + \frac{\tilde{\Phi}^3}{\eta}\right\},$$

$$\tilde{\Phi}^3(1-q(\tilde{\Phi})) = \eta, \quad \eta = \frac{d}{\ell_{ep}}. \tag{4}$$

At $d < \ell_T T/\theta$ in the absence of scattering by impurities we can have either formula (4) (when $\Phi > a$) or the formula

$$\sigma \sim \frac{Ne^2 d}{p_F}\left\{\ln\frac{1}{\Phi^*} + \frac{\Phi^*}{\chi}\right\},$$

$$\Phi^*(1-q(\Phi^*)) = \chi \tag{5}$$

when $\Phi < a$.

We see therefore that the dependence of the resistance on the temperature and on the thickness is governed strongly by the character of the electrons from the surface, and makes it possible to determine the function $q(\Psi)$ from the experimental data. The description of the reflection from the surface by such a function is certainly justified in the case of extremely thin samples, when the integral term in the general linear boundary condition can be neglected. In the general case, the introduction of $q(\Psi)$ has a qualitative character.

At temperatures such that the role of the diffusely-reflected electrons is negligible, formulas (4) and (5) yield in the simplest cases :

1) If $q(\psi) = \begin{cases} 1, & \psi < \phi_0 \\ 0, & \psi > \phi_0 \end{cases}$, then

$$\sigma \sim T^{-3} \quad \text{for} \quad \phi_0 < \frac{T}{\theta} < \left(\frac{\ell_0}{d}\phi_0\right)^{1/3},$$

$$\sigma \sim T^{-5} \quad \text{for} \quad \frac{T}{\theta} < \min\left\{\phi_0,\ \left(\phi_0^3\frac{\ell_0}{e}\right)^{1/5}\right\}. \tag{6a}$$

2) if $1 - q(\Psi) \sim \Psi/\Phi_0$ ($\Psi < \Phi_0$, Φ_0 is a characteristic angle separating the regions of essentially specular scattering from that of essentially diffuse scattering), then

$$\sigma \sim T^{-3/2} \quad \text{for} \quad \frac{T}{\theta} < \min\left\{\left(\phi_0\frac{d}{\ell_0}\right)^{-1},\ \left(\phi_0\frac{\ell_0}{d}\right)^{1/3}\right\}, \tag{6b}$$

$$\sigma \sim T^{-5/4} \quad \text{for} \quad \left(\phi_0 \frac{d}{\ell_\theta}\right)^{-1} < \frac{T}{\theta} < \left(\phi_0^3 \frac{\ell_\theta}{d}\right)^{1/5}.$$

Comparison with experiment [1] apparently favors a q(Ψ) dependence of the "step" type (formula (6a)).

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INVESTIGATION OF THE RADIATION OF STRONG-CURRENT PULSED DISCHARGES IN METAL VAPOR IN THE VACUUM ULTRAVIOLET REGION

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A specially developed photoemission-scintillation analyzer is used to investigate the radiation of strong-current pulsed discharges in Li and In at quantum energies 10 - 100 eV. The measured spectra reveal maxima connected with recombination of the ions LiIII and in II, inIII, and inIV.

In a number of problems connected with the investigation of a dense radiating discharge plasma, great interest attaches to measurements of the distribution of the radiation up to quantum energies $h\nu \sim 100$ eV. The use of diffraction spectrometers in measurements of this kind is made difficult both by their low efficiency (the minimum total number of registered quanta is 10^{14} $cm^{-2}sr^{-1}$ phen photography on film is used [1]) and by the contamination of the grating by the discharge products. We have measured the radiation distribution of strong-current pulsed discharges in Li and In vapor at quantum energies 10 - 100 eV. The parameters of these discharges (I_{max} = 220 kA, $\tau_{\frac{1}{2}}$ = 70 μsec, $N \approx 10^{18}$ cm^{-3}, energy input E = 50 J/cm^3) are given in [2].

The method employed is based on measurement and analysis of the integral spectrum $I = I_0\int f(h\nu,\ \theta,\ \varepsilon)d\theta d\varepsilon$ of the electrons knocked out from the photocathode by direct radiation from the plasma. To this end, a planar grid analyzer transmitted to the detector, from the total photoelectron flux I_0, the electrons with energy $\varepsilon \geq eU_R$, where U_R is the retarding potential of the grid. The electrons passing through the grid were accelerated to 15 kV and registered by a scintillator + light pipe + photomultiplier system (Lincke detector [3]). The photomultiplier signal was recorded with an oscilloscope. The chosen time resolution was $\Delta\tau$ = 1 μsec, and the efficiency of the instrument was 10^{11} quanta-$cm^{-2}sec^{-1}$. A detailed description of the analyzer is given in [4].

The experimental results, $I/I_0 = f(U_R)$, are shown in Fig. 1. They were reduced by using the available data on the energy and angular distributions of the photoelectrons. The latter, as follows from [5], is independent of the quantum energy, at a constant angle of incidence of the quanta, up to an energy on the order of several keV. The energy distribution of the photoelectrons usually contains two smeared maxima [5, 6]. The first is in the 0 - 5 eV region and is due to electrons leaving the photocathode after multiple scattering. The second maximum is