Dislocation slowing down by electrons in metals in strong magnetic fields

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(Submitted January 24, 1974)
Zh. Eksp. Teor. Fiz. 67, 240–249 (July 1974)

We investigate the variation, due to the restructuring of the electron energy spectrum in a quantized magnetic field, of the electron component of dislocation friction parallel to the magnetic field. At extremely low temperatures ($T\to 0$ °K) and in sufficiently pure metals (electron mean free-path time $\tau\to\infty$) the dependence of the friction force on the dislocation velocity has a threshold. For finite values of τ the dependence is nonlinear and complicated, and may even be nonmonotonic at large values of τ.

1. INTRODUCTION

Interest in the effect of electrons on dislocation mobility in metals was stimulated by the experimentally-observed influence of the superconducting transition on sound absorption and on plasticity[1]. A theoretical investigation of electron friction of dislocations in superconductors has shown that it is quite sensitive to the structure of the energy spectrum of the electrons. In addition to the superconducting transition, one of the factors that lead to an appreciable restructuring of the electron spectrum of a metal is a strong magnetic field with $\omega_H\tau \gg 1$ (ω_H is the cyclotron frequency and τ is the electron mean free path). Consequently, placing a metal in a strong magnetic field should lead to an appreciable change of the electronic component of the dislocation friction force, and at low temperatures, when this component plays a noticeable role in the balance of the forces that determine the dislocation mobility, it should lead also to a change in those mechanical properties of the metal which are connected with the dislocation motion.

When calculating the friction forces it is customary to confine oneself to allowance for the deformation interaction between the conduction electrons and the elastic field of the dislocation. For a dislocation moving with constant velocity V, the deformation potential $U(r, t)$ can be represented in the form of a packet of plane waves with dispersion $\omega_q = q\cdot V$:

$$U(\mathbf{r},t)=\sum_{\mathbf{q}} U(\mathbf{q})\exp\{-i\omega_q t\}, \qquad (1)$$

where $U(q) = \lambda_{ik} u_{ik}^{q} e^{iqr}$, u_{ik}^{q} is the Fourier component of the elastic-strain tensor, λ_{ik} is the tensor of the strain-potential constants. In the linear approximation, the action of each of these waves on the electrons can be regarded independently. It turns out also that the main contribution to the drag force is made by waves with $q \sim r_0^{-1} \sim a^{-1}$, where a is the crystal-lattice parameter and r_0 is the dimension of the dislocation core. These circumstances make the problem of electron drag of dislocations equivalent in many respects to the problem of absorption of high frequency sound at $ql \gg 1$ (l is the electron mean free path).

A strong magnetic field leads, as is well known, to different quantum oscillations of thermodynamic and kinetic characteristics of the metal, including the sound-absorption coefficient[2]. In some cases, there is also an appreciable change of the monotonic, non-oscillating part of the absorption coefficient, for example an increase of this part by a factor $\omega_H\tau$[2,3]. It is natural to expect the magnetic field to exert an analogous influence on the electron friction force on the dislocation. To be sure, in the case of a dislocation the oscillation picture is strongly smoothed out in comparison with the case of a monochromatic wave, owing to the additional integration with respect to the component of the wave vector q.

Calculation of the friction force in a magnetic field can be carried out, as in the absence of a field, by two methods[1]. The first (kinetic) is based on a solution of the classical kinetic equation for electrons with a deformation potential (1) as the external perturbation, followed by calculation of the dissipation function. This method was already used earlier, for a study of this problem, by Kravchenko[4], who calculated the drag force of a straight-line dislocation oriented parallel to the magnetic-field direction. At low dislocation velocities $q_m\tau V \ll 1$ ($q_m \sim 1/a$ is the maximum value of the modulus of the wave vector q) the drag force F turned out to be larger by approximately $\omega_H\tau$ times than its value in the absence of a field

$$F={}^1/_2\omega_H\tau B_0 V; \qquad (2)$$

here $B_0 = m^2b^2\lambda^2 q_m/(2\pi\hbar)^3$ is the drag coefficient in the absence of the field[1]. This result is analogous to that obtained by V. Gurevich for the sound-absorption coefficient in the classical case[3].

At higher velocities, $q_m\tau V \gg 1$, the friction force obtained by the kinetic method turned out to be a smooth function that increases slowly with increasing velocity V:

$$F=\omega_H B_0 q_m^{-1}\ln(q_m\tau V). \qquad (3)$$

The applicability of the classical kinetic equation and of formulas (2) and (3) obtained on its basis has, naturally, definite limitations. The classical description of the behavior of electrons in a magnetic field is in general not valid at sufficiently low temperature $T \ll \hbar\omega_H$. At the same time, the electronic components of the dislocation friction force assumes the principal role precisely in the region of low temperatures, and it is this region which is therefore of fundamental interest in its study.

In addition, the problem in question becomes essentially a quantum problem at sufficiently high values of the velocity V and at electron free path time τ in the following two cases: first, when the dislocation can cause electronic transitions between different Landau levels ($q_m V > \omega_H$), and second when the dislocation has time to cross the electron orbit within the time between the collisions that take the electron out of a definite quantum state ($\tau V > r_H$, where r_H is the Larmor radius). Formulas (2) and (3) which are derived classically, therefore do not admit of a transition to the limit as $\tau \to \infty$, and the function $F(V)$ does not contain the oscillating part that should appear at $V > q_m^{-1}\omega_H$ (see Sec. 2).

A second method of calculating the friction force, based on regarding the interaction of the electron with the elastic waves as a quantum-mechanical electron-phonon collision process, is free of the foregoing limitations. The applicability of this method is justified by the decisive role played by waves with $q \gg l^{-1}$. The calculation of the drag force by this method reduces to counting the number of the electronic transitions with energy absorption. This method was used earlier to calculate the dislocation friction force in a normal metal in the absence of a field and in a superconductor[1]. It will be used in the present paper to calculate the dislocation friction force in a metal situated in a strong magnetic field. This method yields for the friction force an expression that is valid in a wide range of variation of the parameters H, τ, and V, in which, in particular, it is possible to analyze the limiting transitions $\tau \to \infty$ and $H \to 0$.

The purpose of the present paper is to establish certain most essential features of the influence of magnetic quantization on the mobility of dislocations in metals. We confine ourselves, following Kravchenko[4], to consideration of the simplest variant of the geometry of the problem—a straight-line dislocation moving with constant velocity V and oriented parallel to the direction of the magnetic field H. We assume the field H to be strong enough, and the temperature and concentration of the impurities low enough, to ensure satisfaction of the conditions

$$\hbar\omega_H \gg T, \quad \omega_H\tau \gg 1. \tag{4}$$

It will be shown below that in a strong magnetic field a very important role will be played by collisions, because not only the absolute value of the friction force F, but also the qualitative form of its dependence on the velocity V depends strongly on the electron mean free path time τ. We have therefore deemed it advantageous to divide the problem into two parts. In Sec. 2 we consider the limiting case $\tau \to \infty$, after which, in Sec. 3, we analyze the role of collisions.

2. FRICTION FORCE IN THE ABSENCE OF COLLISIONS

We use rectangular coordinates with the z axis parallel to the direction of the magnetic field and the dislocation line, and the x axis along the velocity direction **V**, and assume also a standard vector-potential gauge[5]. The quantum-mechanical state of the electron is determined by the oscillator number n, the wave number k_z along the field direction, and the wave number k_x in a plane perpendicular to the field, while the energy levels are described by the formula

$$\varepsilon_{n,k_z} = \hbar\omega_H(n+1/2) + \hbar^2 k_z^2/2m. \tag{5}$$

Since the electron spectrum in a magnetic field is degenerate in k_x, while the deformation potential (1) of a dislocation oriented parallel to the field does not depend on the coordinate z, electronic transitions with energy absorption are possible only between levels with different values of the number n under the condition $\omega_q \geq (n' - n)\omega_H$. Consequently, the function $F(V)$ should have a threshold in the absence of collisions, namely $F(V) = 0$ at $V < V_c$, where V_c is the threshold velocity defined by the relation

$$V_c = q_m^{-1}\omega_H. \tag{6}$$

The general expression obtained for the drag force F by counting the total number of electronic transitions with energy absorption that are induced by the dislocation field is

$$F = \frac{4\pi}{V}\sum_{a,a'}\sum_{\mathbf{q}}\omega_{\mathbf{q}}|U_{a'a}(\mathbf{q})|^2[f(\varepsilon_a) - f(\varepsilon_{a'})]\delta(\varepsilon_{a'} - \varepsilon_a - \hbar\omega_{\mathbf{q}}); \tag{7}$$

here the subscript a labels the set of quantum numbers n, k_z, and k_x, while $U_{a'a}(\mathbf{q})$ is the matrix element of the Fourier component of the deformation potential (1) calculated with the unperturbed wave functions of the electron and of the magnetic field[5], and $f(\epsilon)$ is the Fermi function. Assuming the wave functions to be normalized to a unit volume, we automatically obtain with the aid of (7) the drag force per unit dislocation length.

Using the known properties of Hermite polynomials[6], we obtain

$$|U_{a'a}(\mathbf{q})|^2 = |\lambda_{ik}u_{ik}^{\mathbf{q}}|^2 e^{-x^2} n!n'! \left[\sum_{k=0}^{n} \frac{(-1)^k x^{n-k} x^{n'-k}\chi(n'-k)}{k!(n-k)!(n'-k)!}\right]^2 \delta_{k_z',k_z}\delta_{k_x',k_x+q_x}, \tag{8}$$

where $x = ql_H$, $l_H = (\hbar c/2eH)^{1/2}$ is the magnetic length, and $\chi(n)$ is a unit step function ($\chi(n) = 1$ at $n > 0$ and $\chi(n) = 0$ at $n < 0$). It is convenient to express the sum in (8) in terms of associated Legendre polynomials $L_n^{n'-n}(x^2)$. We shall need later on only the matrix elements with $n \leq n'$, which take the form

$$|U_{a'a}(\mathbf{q})|^2 = |\lambda_{ik}u_{ik}^{\mathbf{q}}|^2 e^{-x^2}(x^2)^{n'-n}\frac{n!}{n'!}[L_n^{n'-n}(x^2)]^2 \delta_{k_z',k_z}\delta_{k_x',k_x+q_x}. \tag{9}$$

The subsequent calculations call for knowledge of the specific form of the deformation potential. By way of example, we consider a screw dislocation, for which

$$\lambda_{ik}u_{ik}^{\mathbf{q}} = ib(\lambda_{xz}q_y - \lambda_{yz}q_x)/q^2, \tag{10}$$

where b is the value of Burger's vector[1)]. Substituting (10) and (9) in (7) and changing from summation with respect to k_z, k_x, and q to integration, we obtain after simple transformations

$$F = \frac{2B_0 V\hbar}{\pi m q_m l_H}\sum_{n,n'}\frac{n!}{n'!}\int dk_z[f(\varepsilon_{n,k_z}) - f(\varepsilon_{n',k_z})]$$
$$\times\int_0^{q_m l_H} dx\, e^{-x^2}x^{2(n'-n)}[L_n^{n'-n}(x^2)]^2 \tag{11}$$
$$\times\int_{-\pi/2}^{\pi/2} d\varphi\cos\varphi(\beta_1\sin^2\varphi + \beta_2\cos^2\varphi)\,\delta\left(n' - n - \frac{Vx}{l_H\omega_H}\cos\varphi\right),$$

where

$$\beta_1 = \frac{4\lambda_{xz}^2}{\lambda_{xz}^2 + 3\lambda_{yz}^2}, \quad \beta_2 = \frac{4\lambda_{yz}^2}{\lambda_{xz}^2 + 3\lambda_{yz}^2},$$

and the coefficient B_0 contains the following combination of the components of the tensor λ_{ik}: $\lambda^2 = (\lambda_{xz}^2 + 3\lambda_{yz}^2)/4$.

In the calculation that follow we shall retain only the first nonvanishing terms in the small parameter $\hbar\omega_H/\epsilon_F$ (ϵ_F is the Fermi energy), and put $T \to 0$ (at finite temperature, the electrons are scattered by phonons, so that we cannot put $\tau \to \infty$). In this approximation, after integrating with respect to k_z and φ, we obtain

$$F=\chi(V-V_c)\frac{8\omega_H B_0}{\pi q_m}\left(\frac{l_H\omega_H}{V}\right)^3\sum_{n=0}^{n_m-k_m}\sum_{k=1}^{k_m}\frac{n!k^4}{(n+k)!}$$

$$\times\left(\frac{\varepsilon_F}{\hbar\omega_H}-\frac{1}{2}-n\right)^{-1/2}\int\limits_{l_H\omega_H k/V}^{l_H q_m} dx e^{-x^2}x^{2k-3}[L_n^k(x^2)]^2 \tag{12}$$

$$\times\left[\frac{\beta_1}{k^2}\left(\frac{V^2}{l_H^2\omega_H^2}x^2-k^2\right)^{1/2}+\beta_2\left(\frac{V^2}{l_H^2\omega_H^2}x^2-k^2\right)^{-1/2}\right];$$

Here $n_m = [\epsilon_F/\hbar\omega_H - 1/2]$, and $k_m = [V/V_c]$ are the integer parts of the quantities in the brackets; $\chi(z) = 1$ at $z > 0$ and $\chi(z) = 0$ at $z < 0$.

It is easily seen that the main contribution to the summation over n in (12) is made by the terms with $n \gg 1$. We can therefore use the asumptotic values of the Laguerre polynomials $L_n^k(x^2)$, and the Stirling formula for $n!$ Substituting these asymptotic forms in (11) and noting, in the integration with respect to x, that one can neglect the contribution of the terms in which the integrand contains oscillating factors, we obtain

$$F=\chi(V-V_c)\frac{8B_0V}{3\pi^2}\left(\frac{V_c}{V}\right)^4\sum_{n=1}^{n_m-k_m}\left[n\left(\frac{\varepsilon_F}{\hbar\omega_H}-\frac{1}{2}-n\right)\right]^{-1/2}$$

$$\times\sum_{k=1}^{k_m}\left[\frac{\beta_1}{2}\left(\frac{V^2}{V_c^2}-k^2\right)^{3/2}+\beta_2\left(\frac{V^2}{V_c^2}-k^2\right)^{1/2}\left(\frac{V^2}{V_c^2}+\frac{1}{2}k^2\right)\right].$$

In real cases we have $n_m \gg k_m$, so that the summation with respect to n can be easily performed. We then obtain for the force F the following final expression:

$$F=\chi(V-V_c)\frac{8B_0V}{3\pi}\left(\frac{V_c}{V}\right)^4\sum_{k=1}^{k_m}\left[\frac{\beta_1}{2}\left(\frac{V^2}{V_c^2}-k^2\right)^{3/2}\right.$$
$$\left.+\beta_2\left(\frac{V^2}{V_c^2}-k^2\right)^{1/2}\left(\frac{V^2}{V_c^2}+\frac{1}{2}k^2\right)\right]. \tag{13}$$

We note the following most essential singularities of the function $F(V)$ as $\tau \to \infty$. As expected, in this case the function $F(V)$ has a threshold; the friction force appears only at dislocation velocities $V > V_c = q_m^{-1}\omega_H$, when transitions between neighboring Landau levels, in which the oscillator number n changes by unity, become possible. The character of the singularity of the function $F(V)$ near the threshold velocity V_c is obvious:

$$F(V)=\chi(V-V_c)\frac{4\sqrt{2}\beta_2}{\pi}B_0V_c\left(\frac{V}{V_c}-1\right)^{1/2},\quad \frac{V}{V_c}-1\ll 1. \tag{14}$$

The behavior of the function $F(V)$ in the velocity interval $(V_c, 2V_c)$ is described by the formula

$$F(V)=\frac{8}{3\pi}B_0V\left(\frac{V_c}{V}\right)^4\left[\frac{1}{2}\beta_1\left(\frac{V^2}{V_c^2}-1\right)^{3/2}+\beta_2\left(\frac{V^2}{V_c^2}-1\right)^{1/2}\left(\frac{V^2}{V_c^2}+\frac{1}{2}\right)\right],$$
$$V_c<V<2V_c. \tag{15}$$

With further increase of the dislocation velocity, electronic transitions with change of quantum number n by 2, 3, etc. become possible, so that the number of terms in (13) increases and the friction force acquires periodic increments. The schematic form of the function $F(V)$ as $\tau \to \infty$ is shown in Fig. 1.

At very large velocities $V \gg V_c$, the friction force approaches asymptotically its value in the absence of a field:

$$F(V)=B_0V[1-O(V_c/V)]. \tag{16}$$

This result is perfectly natural, since at large velocities $V \gg V_c$ the characteristic energy $\hbar q_m V$ transferred to the electrons from the dislocation extends over an increasing number of Landau levels and the dislocation ceases to "feel" the discreteness of the electron spectrum. We note also that the inequality $V \gg V_c$ can be realized at any finite value of the velocity V in a sufficiently weak magnetic field. Thus, formula (16) permits a natural limiting transition to the case $H = 0$ and yields the friction force in the absence of a field: $F_0 = B_0V$.

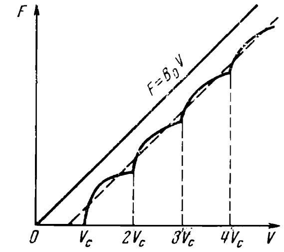

FIG. 1. Dependence of the dislocation friction force in a magnetic field on the dislocation velocities of $\tau \to \infty$. Straight line—friction force in the absence of the field.

3. INFLUENCE OF COLLISIONS

The presence of scattering leads to a smearing of the electron levels by an amount $\delta\epsilon \sim \hbar/\tau$. In the case of interest to us, that of a strong magnetic field ($\omega_H\tau \gg 1$), this smearing is much smaller than the distance $\hbar\omega_H$ between the Landau levels. In spite of this, the influence of the collisions on the friction force $F(V)$ and the character of its dependence on the dislocation velocity V can be appreciable. The validity of the statement made above can be illustrated by the following three samples. First, it is obvious that the smearing of the Landau levels should lead to a vanishing of the threshold effect in the dependence of $F(V)$, and by the same token alter qualitatively the course of this dependence near values of V that are multiples of V_c. Second, at dislocation velocities satisfying the inequality $\hbar q_m V < \delta\epsilon$, the dislocation-induced electronic transitions will take place practically in the continuous spectrum, and the friction force F should then be in any case not smaller than in the absence of a field; we recall that as $\tau \to \infty$ the friction force at these velocities is rigorously equal to zero.

For a detailed account of the collisions in the considered problems it is necessary, generally speaking, to use the quantum-kinetic equation. However, semiquantitative estimates for the friction force $F(V)$, which are of practical interest, can be obtained also within the framework of the quantum-mechanical perturbation-theory method employed in Sec. 2, if account is taken of the uncertainty, due to the collisions, in the energy conservation law when an electron is scattered by deformation-potential waves. We shall take this uncertainty into account by following the work of Gurevich, Skobov, and Firsov[7], replacing the δ-functions in formulas (7) or (11) by a Lorentz function of width $\delta\epsilon = \hbar/\tau$:

$$\delta(\varepsilon_{a'}-\varepsilon_a-\hbar\omega_q)\to\frac{1}{\pi}\frac{\hbar/\tau}{(\varepsilon_{a'}-\varepsilon_a-\hbar\omega_q)^2+\hbar^2/\tau^2}. \tag{17}$$

We confine ourselves henceforth to the velocity interval $V < V_c$, in which the influence of the collisions is most significant. At higher collision velocities, while the collisions do lead to a vanishing of threshold effects near velocity values that are multiples of V_c, everywhere else they produce insignificant changes both in the value and in the form of the function $F(V)$.

Substituting (17) in (11) and noting that at $V < V_c$ the contribution of the terms with $n \neq n'$ can be neglected, we obtain

$$F=\frac{2\hbar\omega_H\tau B_0}{\pi^2 m q_m l_H}\sum_{n=0}^{\infty}\int_{-\infty}^{\infty}dk_z\int_{0}^{l_H q_m}dx\int_{-\pi/2}^{\pi/2}d\varphi$$
$$\times\cos\varphi\,(\beta_1\sin^2\varphi+\beta_2\cos^2\varphi)\frac{e^{-x^2}[L_n(x^2)]^2}{1+(\tau^2V^2/l_H^2)x^2\cos^2\varphi}\qquad(18)$$
$$\times\left[f(\varepsilon_{n,k_z})-f\left(\varepsilon_{n,k_z}+\frac{\hbar V}{l_H}x\cos\varphi\right)\right],$$

where $L_n(x^2) \equiv L_n^0(x^2)$.

At the temperatures $T \ll \hbar\omega_H$ of interest to us, we can put $T = 0$ in the function $f(\epsilon)$, as before (the resultant error is obviously small to the extent that $T/\hbar\omega_H$ is small). It must be remembered, however, that in crystals that are sufficiently pure of impurities, the temperature influences also the free-path time τ. If this influence is appreciable, then it is necessary to take into account the $\tau(T)$ relation, but since τ is introduced phenomenologically, the allowance for this relation is not reflected in the subsequent calculations.

Replacing $f(\epsilon)$ in (18) by a step function and retaining, in the integration with respect to k_z and φ, the first terms that do not vanish in the small parameters $\hbar\omega_H/\epsilon_F$ and $V/q_m l_H V_c$, we obtain

$$F=\frac{4B_0\omega_H\tau V}{\pi q_m l_H}\left(\frac{l_H}{\tau V}\right)^4\sum_{n=0}^{n_m-1}\left(\frac{\varepsilon_F}{\hbar\omega_H}-n-\frac{1}{2}\right)^{-1/2}$$
$$\times\int_0^{l_H q_m}dx\frac{e^{-x^2}[L_n(x^2)]^2}{x^3}\left\{\beta_1\left[1-\left(\frac{\tau^2V^2}{l_H^2}x^2+1\right)^{1/2}+\frac{\tau^2V^2}{2l_H^2}x^2\right]\right.\qquad(19)$$
$$\left.+\beta_2\left[\left(\frac{\tau^2V^2}{l_H^2}x^2+1\right)^{-1/2}+\frac{\tau^2V^2}{2l_H^2}x^2-1\right]\right\}.$$

The subsequent calculations (integration with respect to x and summation over n) raise no particular difficulties, but lead to extremely cumbersome results. We therefore proceed directly to discuss the limiting cases in which the expression for the force F acquires a relatively compact form; at the same time, an analysis of these cases provides a qualitative idea of the character of the dependence of F on V and on τ.

The decisive role in our problem is played by the value of the parameter $r_H^{-1}\tau V$, where $r_H = 2l_H(\epsilon_F/\hbar\omega_H)^{1/2}$ is the Larmor radius of the electron. This parameter is none other than the ratio of the lifetime τ of the electron in a definite quantum state to the time of crossing of the electron orbit and the moving dislocation r_H/V (i.e., the characteristic time of interaction of the dislocation with the electron in the given state). We consider first the case when during the time r_H/V the electron is taken many times out of the given quantum state by the collisions, i.e., $r_H^{-1}\tau V \ll 1$. This case incorporates in turn three more particular cases defined by the inequalities $q_m\tau V \ll 1$; $q_m\tau V \gg 1$, but $l_H^{-1}\tau V \ll 1$ and $l_H^{-1}\tau V \gg 1$. Allowance for these inequalities in (19) leads to the following expressions for the force F:

$$F=\frac{\beta_1+3\beta_2}{4\pi}\tau\omega_H B_0 V,\quad q_m\tau V\ll 1;\qquad(20)$$

$$F=\frac{2(\beta_1+2\beta_2)\omega_H B_0}{3\pi q_m}\left(1-\frac{3}{2}\frac{\beta_1+\beta_2}{\beta_1+2\beta_2}\frac{1}{q_m\tau V}\right),$$
$$q_m\tau V\gg 1,\quad l_H^{-1}\tau V\ll 1;\qquad(21)$$

$$F=\frac{2(\beta_1+2\beta_2)\omega_H B_0}{3\pi q_m}\left[1-\frac{3}{2}\frac{\beta_1+\beta_2}{\beta_1+2\beta_2}\frac{1}{q_m\tau V}\right.$$
$$\left.+\frac{3(\beta_1+3\beta_2)}{\beta_1+2\beta_2}\left(1-\frac{1}{\pi}\right)\frac{\tau V}{2r_H}\ln\left(\frac{2r_H}{\tau V}\right)\right],\quad l_H^{-1}\tau V\gg 1.\qquad(22)$$

Comparing formulas (20)—(22) with (2) and (3), we see that the quantum-mechanical method of calculating F(V) leads in the case $r_H^{-1}\tau V \ll 1$ to a result that agrees qualitatively with the result obtained on the basis of the classical kinetic equation, namely the friction force at small velocities is linear in the velocity with a coefficient on the order of $\tau\omega_H B_0$, and at larger velocities it increases slowly with increasing velocity. The problem becomes essentially a quantum problem in the opposite limiting case $r_H^{-1}\tau V \gg 1$, i.e., when the collisions do not manage to take the electron out of the specified quantum state during the time of interaction of the dislocation. In this case we have

$$F=\frac{(\beta_1+\beta_2)\omega_H B_0}{\pi q_m}\frac{2r_H}{\tau V}\ln\left(\frac{\tau V}{2r_H}\right).\qquad(23)$$

This expression admits of a transition to the limit as $\tau \to \infty$; it is obvious that as $\tau \to \infty$ we have $F \to 0$, as should be the case in accordance with (13) at $V < V_c$.

The schematic form of the F(V) dependence at different values of τ is shown in Fig. 2. The qualitative cause of the transition from the monotonic F(V) dependence at $\tau = \tau_1$ to the nonmonotonic one at $\tau = \tau_3$ in the velocity interval $V < V_c$ consists in the following. At sufficiently low values of τ and V, when the characteristic energy $\hbar q_m V$ transferred to the electron in a single interaction is much lower than the smearing $\hbar/\tau$ of the Landau level, the friction force is proportional to $\hbar q_m V$ and to the number of approximately equivalent interaction acts $\omega_H\tau$, namely, $F \propto \hbar q_m V\omega_H\tau$. With further increase of the velocity or of the free-path time, the situation is realized wherein $\hbar q_m V \gg \hbar/\tau$, but still $\tau V \ll r_H$ (the latter inequality means that after $\omega_H\tau$ revolutions of the electron in the magnetic field the dislocation does not have time to cross its orbit); in this case the characteristic energy transferred to the electron in a single interaction act is of the order of $\hbar/\tau$, and the friction force is $F \propto (\hbar/\tau)\omega_H\tau$, i.e., F becomes insensitive to either a growth of V or a growth of τ. Finally, at $\tau V \gg r_H$ the energy transfer is as before on the order of $\hbar/\tau$, and the number of approximately equivalent interaction acts becomes equal to $r_H\omega_H/V$ (the number of electron revolutions in the time that the dislocation crosses its orbit); it is obvious that in this case $F \propto \hbar\omega_H r_H/\tau V$, i.e., it decreases with decreasing V and τ. Since the interval of variation of the velocities was limited by us to the value V_2, only the first of the cases indicated above can be realized at $\tau \ll r_H V_c^{-1}$, and at $\tau \gg r_H V_c^{-1}$ the third case can also be realized.

Thus, the foregoing results show that the magnetic quantization leads to the appearance of two essential singularities in the function F(V): first, an oscillating increment appears in this function at $V > V_c$; second, the function F(V) becomes nonmonotonic at sufficiently large electron mean free path times $\tau > r_H V_c^{-1}$. These

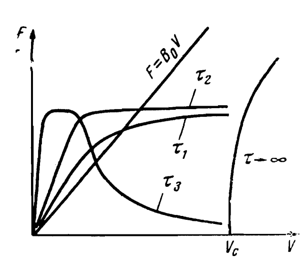

FIG. 2. Plot of F(V) at different values of τ: $\tau_1 < \tau_2 < \tau_3$, $\tau_1, \tau_2 \ll r_H V_c^{-1}$, $\tau_3 \gg r_H V_c^{-1}$.

two singularities of the electronic component of the dislocation friction force should become manifest in the study of the low-temperature plasticity of metals in strong magnetic fields. Simple estimates show that the conditions for their realization are within the capabilities of modern experimental facilities ($T \sim 1^\circ K$, $H \sim 10^5$ Oe, $\tau \sim 10^{-8}$–10^{-9} sec, $V \sim 10^4$–10^5 cm/sec).

The authors are grateful to M. I. Kaganov, V. Ya. Kravchenko, A. M. Kosevich, and V. G. Peschanskiĭ for interest in the work and for useful discussion of the results.

[1)]The components of the tensor λ_{ik}, which enter in this expression, should take into account the screening of the bare deformation potential by the conduction electrons [2,7]. In a strong magnetic field near the dislocation core, bound states appear [8,9], so that the character of the screen, and hence of λ_{ik}, depends in general on the magnitude of the field (the need for taking bound states into account was pointed out to the authors by A. M. Kosevich). The analysis of this dependence, however, is beyond the scope of the present article and should be the subject of a separate investigation.

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Translated by J. G. Adashko
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