FIG. 3. Curve 1) the ratio $T_C^{id}(\lambda_0)/T_C^0$ of the Curie temperatures of ideal quasi-two-dimensional and isotropic ferromagnetic crystals according to Lines[10]; 2) the ratio $T_C^{am}(\lambda_0)/T_C^0$ of the Curie temperatures of the amorphous isotropic and ideal isotropic ferromagnets; 3) the ratio $\bar{T}_C(\lambda_0)/T_C^0$ of the Curie temperatures of the amorphous isotropic FM (in the approximation of an average $\bar{I}$) and the ideal isotropic FM; 4) relative increase $T_C^{am}(\lambda_0)/T_C^{id}(\lambda_0)$ of the Curie temperature on amorphization of a quasi-two-dimensional FM.

ferromagnets with anisotropically distributed exchange couplings, such as quasi-two-dimensional magnets. We have shown, e.g., that the amorphization of a quasi-two-dimensional FM leads to a substantial increase of T_C, especially for small values of the interplanar exchange. In real substances we must, of course, also take into account the decrease of the quantities K_0 and J_0 themselves, which leads to a certain lowering of the curve 2 in Fig. 3. For example, for $\lambda_0 = 1$ the ratio T_C^{am}/T_C^0 will be less than unity. However, the results are not qualitatively changed. A similar remark can also be made concerning the magnetization of an amorphous FM. It is clear, of course, that the amorphization of a quasi-one-dimensional FM is not qualitatively different from the case considered. As already stated, the situation is considerably more complicated in the case of amorphization of quasi-two-dimensional anti-ferromagnets. However, we may expect that for small values of $|\lambda_0|$ the amorphization of systems with $J_0 > 0$ and $K_0 < 0$ can lead to ferromagnetism with a relatively high T_C. In view of the fact that quasi-low-dimensional magnets constitute a broad class of magnets, we may suppose that their amorphization could turn out to be an important way of obtaining new magnetically ordered substances.

[1]A. I. Gubanov, Fiz. Tverd. Tela(Leningrad) 2, 502 (1960) [Sov. Phys. Solid State 2, 468 (1960)].
[2]K. Handrich, Zh. Eksp. Teor. Fiz. 64, 1383 (1973) [Sov. Phys. JETP 37, 703 (1973)].
[3]C. G. Montgomery, J. I. Krugler and R. M. Stubbs, Phys. Rev. Lett. 25, 669 (1970).
[4]E-Ni Foo and S. M. Bose, Phys. Rev. B9, 347 (1974).
[5]J. E. Gubernatis and P. L. Taylor, Phys. Rev. B9, 3828 (1974).
[6]K. Tamura and H. Endo, Phys. Lett. 29A, 52 (1969).
[7]A. W. Simpson and D. R. Brambley, Phys. Stat. Sol. 49B, 685 (1972).
[8]J. J. Rhyne, S. J. Pickart and H. A. Alperin, Phys. Rev. Lett. 29, 1562 (1972).
[9]K. A. Sablina, G. A. Petrakovskiĭ, E. N. Agartanova and V. P. Piskorskiĭ, Pis'ma Zh. Eksp. Teor. Fiz. 24, 357 (1976) [JETP Lett. 24, 324 (1976)].
[10]M. E. Lines, Phys. Rev. 133, A841 (1964).
[11]S. V. Tyablikov, Metody kvantovoĭ teorii magnetizma (Methods in the Quantum Theory of Magnetism), Nauka, M., 1965 (English translation published by Plenum Press, N. Y., 1967).
[12]B. Velický, S. Kirkpatrick and H. Ehrenreich, Phys. Rev. 175, 747 (1968).
[13]P. Soven, Phys. Rev. 156, 809 (1967).
[14]R. J. Elliott, J. A. Krumhansl and P. L. Leath, Rev. Mod. Phys. 46, 465 (1974).
[15]N. Mermin and H. Wagner, Phys. Rev. Lett. 17, 1133 (1966).

Translated by P. J. Shepherd.

Contribution to the theory of electromechanical forces in metals

M. I. Kaganov and V. B. Fiks

Physical Problems Institute, USSR Academy of Sciences
(Submitted March 23, 1977)
Zh. Eksp. Teor. Fiz. 73, 753–760 (August 1977)

A mechanism is considered for production of electromechanical forces caused by electric current and concentrated near crystallite interfaces. The order of magnitude of the forces is ascertained for simple models of the interface and of the dispersion law.

An electric field **E** applied to a metal produces an electron current and motion of lattice defects: atoms, dislocations, inclusions, grain boundaries. It does not disturb the mechanical equilibrium of an ideal crystal lattice. This is a dynamic equilibrium: the forces of the external field, acting on the ions of the "skeleton" of the metal, are exactly compensated by the forces produced on scattering of the electrons in the lattice (the electronic wind[11]). The observable manifestations of the force of the electronic wind depend on the mech-

anism of transfer and redistribution of the momentum given up to the lattice by the electronic subsystem, and they always are connected with a disturbance of the local equilibrium. In all defective regions of the crystal, the equilibrium is certainly disturbed, and electromechanical forces necessarily appear. The disturbance of the lattice structure occurs with special intensity in the process of plastic deformation, and the electromechanical forces are concentrated where structural inhomogeneities have arisen. It is perhaps for this reason that electromechanical forces play a comparatively important role in plastic deformation.[2-4] The purpose of the present article is to point out several mechanisms of production of the electromechanical forces and to estimate their magnitude.

We consider the simplest situation: passage of current parallel to an intercrystallite boundary that serves to produce additional scattering of electrons. Assuming that the dimensions of the crystallites are considerably larger than the length l of the free path of the electrons, we suppose that each crystallite occupies a half-space. The plane $z = 0$ serves as the boundary. We shall attach the index $s = 1$ to all quantities related to the left crystallite ($z < 0$) and the index $s = 2$ to all related to the right. To avoid preoccupation with effects caused by unimportant anisotropy of each crystallite, we shall suppose that the direction of the current (of density j) coincides with principal directions of both crystallites. In this case the electric field intensity $\mathbf{E}$ is parallel to j in both crystallites. We take the x axis along $\mathbf{E}$ and j ($E_x = E$, $j_x = j$). With such geometry, the electromechanical force density $\vec{\mathscr{F}}$ is also directed along the x axis, $\mathscr{F}_x = \mathscr{F}$. According to Ref. 5, calculation of $\mathscr{F}$ requires knowledge of the dependence on the coordinate z of the nonequilibrium correction $f_s(\mathbf{p}, z)$ to the Fermi distribution function $F(\varepsilon)$ of the electrons:

$$\mathscr{F}_s(z) = \frac{2}{(2\pi\hbar)^3} \int \Lambda_{xz}^{(s)}(\mathbf{p}) \frac{\partial f_s(\mathbf{p}, z)}{\partial z} d^3p,$$

$\mathbf{p}$ is the quasimomentum, and $\varepsilon = \varepsilon_s(\mathbf{p})$ is the energy of a conduction electron in the s-th crystallite. The tensors $\Lambda_{ik}^{(s)}(\mathbf{p})$ can be determined only by a model calculation of the electronic energy spectrum of the metal. For order-of-magnitude estimates, we shall use the expression

$$\Lambda_{ik}^{(s)} = \tilde{m}_s v_i^{(s)} v_k^{(s)},$$

$\mathbf{v}_s(\mathbf{p}) = \partial \varepsilon_s(\mathbf{p})/\partial \mathbf{p}$ is the electron velocity, and the parameter $\tilde{m}_s \approx m_0$, where m_0 is the mass of a free electron.

In an invariant description (in a coordinate system attached to the crystallographic axes), the dispersion laws $\varepsilon_s(\mathbf{p})$ and the tensors $\Lambda_{ik}^s(\mathbf{p})$ are of course the same in both crystals (independent of s); but in the coordinate system that we have chosen, they differ from one to another (see below). We shall omit the index s when this can cause no misunderstanding. We shall furthermore write $v_x = u$, $v_z = w$. Thus

$$\mathscr{F}(z) = \frac{2}{(2\pi\hbar)^3} \int \tilde{m} u w \frac{\partial f}{\partial z} d^3p. \qquad (1)$$

Because of the neutrality of the metal, the total force exerted on it by the electric field must be zero. If

$$\int_{-\infty}^{\infty} \mathscr{F}(z) \, dz \neq 0,$$

then this means that the intercrystallite boundary is acted upon by a concentrated force $\mathscr{F}_{\text{surf}} \delta(z)$, such that

$$\mathscr{F}_{\text{surf}} = -\int_{-\infty}^{\infty} \mathscr{F}(z) \, dz. \qquad (2)$$

For $|z| \gg 1$, the role of the intercrystallite boundary is insignificant, and f is asymptotically independent of z:

$$f(\mathbf{p}, z) \to -\frac{\partial F}{\partial \varepsilon} eu\tau E, \qquad |z| \to \infty, \qquad (3)$$

τ is the free-passage time of an electron. The electric field intensity E, from the irrotationality condition (curl $\mathbf{E} = 0$), is independent of the coordinates.

On both sides of the boundary, the function f is a solution of Boltzmann's kinetic equation

$$w \frac{\partial f}{\partial z} + \frac{f}{\tau} = -\frac{\partial F}{\partial \varepsilon} euE, \qquad (4)$$

for which the boundary conditions at $z = \pm \infty$ are the equalities (3). The conditions on the boundary (at $z = 0$) are of course wholly determined by the structure of the intercrystallite layer, whose thickness is of the order of the interatomic distance. They cannot be formulated in general terms. But it is easy to imagine two limiting situations.

A. A diffuse boundary, scattering of the electrons on which is accompanied by complete loss of directed motion[1]:

$$f_1^{(-)}|_{z=0} = 0, \qquad f_2^{(+)}|_{z=0} = 0. \qquad (5)$$

We have written $f_s^{(\pm)}$ for $f_s(z, u_s \gtrless 0)$.

B. A specular boundary, on which the tangential projection $\mathbf{p}_t$ of the quasimomentum of an electron is conserved. An electron wave passing through the boundary is refracted and reflected. The possibility of assigning electron states with a value of the quasimomentum parallel to the boundary means that the bicrystal is periodic in planes parallel to the boundary (Fig. 1). Since the
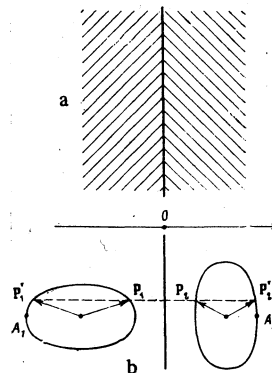

FIG. 1. a, Structure of a specular boundary. b, Isoenergetic surfaces $\varepsilon = \varepsilon_s(\mathbf{p})$. An electron with quasimomentum $\mathbf{p}_1$ has, on reflection from the boundary, quasimomentum $\mathbf{p}_1'$; and on passage across the boundary, quasi-momentum $\mathbf{p}_2'$ ($\mathbf{p}_1' \to \mathbf{p}_1'$, $\mathbf{p}_2'$). Analogously, $\mathbf{p}_2 \to \mathbf{p}_2'$, $\mathbf{p}_1'$. A_s are limiting points of the Fermi surface, corresponding to the asymptotic behavior of the electromechanical force.

energy as well as p_t is conserved in passage of the electron across the boundary, a relation can be established between the quasimomenta and velocities of an electron in the two crystallites. Figure 1b shows how to accomplish this. We remark that in the case of an anisotropic Fermi surface, there must always be electrons that undergo total internal reflection at the boundary of the crystallites. The boundary conditions for the distribution functions $f_s^{(\pm)}$ are consequences of the law of conservation of the number of particles on passage across the plane $z = 0$, on which

$$|w_1| f_1^{(-)} = R_1 w_1 f_1^{(+)} + (1 - R_2) |w_2| f_2^{(-)},$$
$$w_2 f_2^{(+)} = R_2 |w_2| f_2^{(-)} + (1 - R_1) w_1 f_1^{(+)}. \tag{6}$$

The coefficients R_s $(0 \leqslant R_s \leqslant 1)$ describe properties of the boundary; their physical meaning is clear from equations (6), which are a natural generalization of the Fuchs condition for specular reflection. When $R_1 = R_2 = 1$, the electrons of both half-spaces are specularly reflected, without crossing the plane $z = 0$ (two separate half-spaces with specular boundaries). If the two crystallites are identical, then $R_1 = R_2 = 0$ (then, of course, $v_1 = v_2$, $\tau_1 = \tau_2$, and in equation (1) there is no term at all containing $\partial f / \partial z$). For electrons that undergo total internal reflection, $R \equiv 0$.

Calculation of the coefficients R_s in the general case is a complicated problem, not yet solved (as far as we know). We shall hereafter treat the R_s as phenomenological parameters; but at the end of the article, we shall calculate the R_s by assuming that on both sides of the boundary the approximation of effective anisotropic masses is valid.

A. Diffuse boundary

The solution of equation (4) with the boundary conditions (3) and (5) has the form

$$
\left.
\begin{aligned}
f_1^{(+)} &= -\frac{\partial F}{\partial \varepsilon} e \tau_1 u_1 E \\
f_1^{(-)} &= -\frac{\partial F}{\partial \varepsilon} e \tau_1 u_1 E \left(1 - \exp\left\{\frac{z}{\tau_1 |w_1|}\right\}\right)
\end{aligned}
\right\} \quad z < 0,
$$

$$
\left.
\begin{aligned}
f_2^{(+)} &= -\frac{\partial F}{\partial \varepsilon} e \tau_2 u_2 E \left(1 - \exp\left\{-\frac{z}{\tau_2 w_2}\right\}\right) \\
f_2^{(-)} &= -\frac{\partial F}{\partial \varepsilon} e \tau_2 u_2 E.
\end{aligned}
\right\} \quad z > 0. \tag{7}
$$

On substituting the expressions (7) in formula (1), we have

$$\mathscr{F}_1 = -\frac{2eE}{(2\pi\hbar)^3} \int_{w_1<0} \tilde{m}_1 u_1^2 \frac{\partial F}{\partial \varepsilon} \exp\left\{\frac{z}{\tau_1 |w_1|}\right\} d^3p, \quad z < 0,$$

$$\mathscr{F}_2 = -\frac{2eE}{(2\pi\hbar)^3} \int_{w_2>0} \tilde{m}_2 u_2^2 \frac{\partial F}{\partial \varepsilon} \exp\left\{-\frac{z}{\tau_2 w_2}\right\} d^3p, \quad z > 0. \tag{8}$$

In the general case (for an arbitrary dispersion law), it is of course impossible to calculate the integrals that enter in (8). But it is possible to make an estimate without specializing the dispersion law. For example, it follows from (8) that

$$\mathscr{F}(-0) = {}^1\!/_2 \gamma_1 eEn, \quad \mathscr{F}(+0) = {}^1\!/_2 \gamma_2 eEn, \tag{9}$$

where n is the electron density, and where the γ_s are

dimensionless parameters $(\gamma_s = \langle \tilde{m}_s / m_{x_s s} \rangle$, the mean (over the band) ratio of $\tilde{m}_s$ to the effective mass $m_{x_s s}$ $= (\partial u_s / \partial p_x)^{-1})$.

The surface force $\mathscr{F}_{\text{surf}}$ in this case is nonzero. According to (2) and (7)

$$\mathscr{F}_{\text{surf}} = -\frac{2eE}{(2\pi\hbar)^3}\left\{\int_{w_1>0} \tau_1 \tilde{m}_1 u_1^2 w_1 \left(-\frac{\partial F}{\partial \varepsilon}\right) d^3p\right.$$
$$\left. + \int_{w_2>0} \tau_2 \tilde{m}_2 u_2^2 w_2 \left(-\frac{\partial F}{\partial \varepsilon} \int d^3p\right)\right). \tag{10}$$

We shall hereafter use the simplest anisotropic dispersion law, which in the invariant description has the form

$$\varepsilon(\mathbf{p}) = p_\perp^2 / 2m_\perp + p_\parallel^2 / 2m_\parallel, \tag{11}$$

and we shall suppose that in the first crystallite

$$\varepsilon_1(\mathbf{p}) = \frac{p_x^2}{2m_\parallel} + \frac{p_y^2 + p_z^2}{2m_\perp}, \quad u_1 = \frac{p_x}{m_\parallel}, \quad w_1 = \frac{p_z}{m_\perp}, \tag{11a}$$

whereas in the second crystallite

$$\varepsilon_2(\mathbf{p}) = \frac{p_x^2 + p_y^2}{2m_\perp} + \frac{p_z^2}{2m_\parallel}, \quad u_2 = \frac{p_x}{m_\perp}, \quad w_2 = \frac{p_z}{m_\parallel}. \tag{11b}$$

Then, supposing that $-\partial F / \partial \varepsilon = \delta(\varepsilon - \varepsilon_F)$, we get from (10)

$$\mathscr{F}_{\text{surf}} = -\frac{3}{16} eEn\left\{\frac{\tilde{m}_1}{m_\parallel} l_1 + \frac{\tilde{m}_2}{m_\perp} l_2\right\}, \tag{12}$$

where $l_{1(2)} = \tau_{1(2)} (2\varepsilon_F / m_{\perp (\parallel)})^{1/2}$ is the free-path length in the first (second) crystallite. In the calculation it was assumed that τ_s and $\tilde{m}_s$ are independent of $\mathbf{p}$.

The structure of the electromechanical force corresponding to the expressions (8), (9), and (12) is represented in Fig. 2. An important characteristic of the electromechanical force is its asymptotic behavior for $|z| \gg l$. According to (8), the asymptotic behavior of $\mathscr{F}(z)$ is determined by those sections on the Fermi surface where $|w|$ attains maxima, i.e., sections near a limiting point (the points A_1 and A_2 in Fig. 1). At a limiting point, as a rule, u_s vanishes. As a result, the electromechanical force attenuates faster than an exponential function. By using the expressions (11a) and (11b), one easily obtains from (8)

$$\mathscr{F}_s(z) \sim \frac{3}{2} neE \frac{\tilde{m}_s}{m_s} \left(\frac{l_s}{z}\right)^2 \exp\left(-\frac{|z|}{l_s}\right), \tag{13}$$
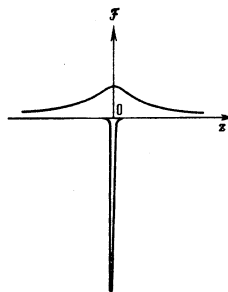$$m_1 = m_\parallel, \quad m_2 = m_\perp.$$

FIG. 2. Electromechanical force density with a diffuse boundary.

395   Sov. Phys. JETP 46(2), Aug. 1977
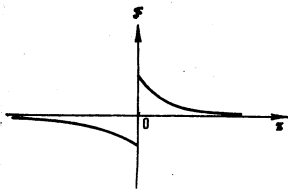
M. I. Kaganov and V. B. Fiks   395

FIG. 3. Electromechanical force density with a specular boundary.

B. Specular boundary

From the boundary condition (6) it follows that the force concentrated on the boundary of the crystallites in this case vanishes. This is in agreement with the "specularity" of the boundary, on which the electrons do not lose the component of momentum along the field. Thus around the intercrystallite boundary there arises, on passage of current, a peculiar force dipole (see below and Fig. 3).

The nonequilibrium part of the distribution function of electrons flying away from the boundary of the crystallites—the solution of equation (4) with the boundary condition (6)—has the form

$$f(z, w) = -\frac{\partial F}{\partial \varepsilon} e\tau u E \left(1 + \xi \exp\left\{-\frac{|z|}{\tau|w|}\right\}\right), \tag{14}$$

where

$$\xi_1 = R_1 - 1 + (1 - R_2)\frac{|w_2||u_2|\tau_2}{|w_1||u_1|\tau_1}$$

$$\xi_2 = R_2 - 1 + (1 - R_1)\frac{|w_1||u_1|\tau_1}{|w_2||u_2|\tau_2} . \tag{15}$$

The distribution function of electrons flying toward the boundary is no different from the distribution function in a homogeneous metal. We recall that the values of v_1 and v_2 are not independent: in calculating the electromechanical force for $z > 0$, we must, by use of the laws of refraction of electron waves, determine $v_1 = v_1(p)$, where p is the quasimomentum of an electron in the right crystallite; and conversely, in calculating $\mathcal{F}$ for $z < 0$, we must determine $v_2 = v_2(p)$ as a function of the quasimomentum of an electron in the left crystallite (see below).

If we do not specialize the dispersion law of the electrons, then ξ_1 and ξ_2 (like R_1 and R_2) must be considered phenomenological parameters describing the boundary. We shall estimate the value of the electromechanical force near the boundary of the crystallites (at $z = \pm 0$). From (14) and (1) we have

$$\mathcal{F}(-0) = \frac{1}{2}eEn\langle\xi_1\rangle, \quad \mathcal{F}(+0) = \frac{1}{2}eEn\langle\xi_2\rangle, \tag{16}$$

$$\langle\xi_i\rangle = \frac{4}{n(2\pi\hbar)^3}\int \xi_i \tilde{m}_i u_i^2\left(-\frac{\partial F}{\partial \varepsilon}\right)d^3p.$$

If the relative anisotropy of the crystallites is such that $|\langle\xi_1\rangle| \approx |\langle\xi_2\rangle| \approx 1$, then the values of $\mathcal{F}(\pm 0)$ coincide in order of magnitude with $\mathcal{F}(\pm 0)$ on a diffuse boundary. It must be borne in mind that, as a rule, $\langle\xi_1\rangle$ and $\langle\xi_2\rangle$ are of opposite signs and much smaller than unity (see below). The asymptotic behavior of $\mathcal{F}(z)$ (for $|z| \to \infty$) in this case also is determined by a limiting point of the Fermi surface (see above) and for the most part has a

dependence on $|z|$ in agreement with (13).

We shall carry out a calculation of the reflection coefficients R_1 and R_2 by using the simplest model (11), (11a), and (11b); for definiteness, we shall take $m_\parallel > m_\perp$. The wave function $\Psi(r)$ of an electron according to (11a) and (11b) satisfies the following Schrödinger equations:

$$\frac{\hbar^2}{2m_\parallel}\frac{\partial^2\Psi}{\partial x^2} + \frac{\hbar^2}{2m_\perp}\left(\frac{\partial^2\Psi}{\partial y^2} + \frac{\partial^2\Psi}{\partial z^2}\right) + \varepsilon\Psi = 0, \quad z < 0, \tag{17a}$$

$$\frac{\hbar^2}{2m_\perp}\left(\frac{\partial^2\Psi}{\partial x^2} + \frac{\partial^2\Psi}{\partial y^2}\right) + \frac{\hbar^2}{2m_\parallel}\frac{\partial^2\Psi}{\partial z^2} + \varepsilon\Psi = 0, \quad z > 0. \tag{17b}$$

The boundary conditions express the continuity of the wave function and of the probability current density:

$$\Psi(z = +0) = \Psi(z = -0), \quad \frac{1}{m_\perp}\frac{\partial\Psi}{\partial z}\bigg|_{z = -0} = \frac{1}{m_\parallel}\frac{\partial\Psi}{\partial z}\bigg|_{z = +0}. \tag{18}$$

We shall consider reflection and refraction, from the boundary of the crystallites, of an electron wave going from left to right, for which we shall seek $\Psi(r)$ in the following form:

$$\Psi(r) = \exp[i(rx + p_y y)/\hbar]\Psi(z), \tag{19}$$

where

$$\Psi(z) = \exp(iq_1 z/\hbar) + a\exp(-iq_1 z/\hbar), \quad z < 0$$
$$\Psi(z) = b\exp(iq_2 z/\hbar), \quad z > 0 . \tag{20}$$

We have written $p_x = r$, $p_z = q$.

The coefficients a and b are determined from the boundary conditions (18):

$$a = \frac{m_\parallel q_1 - m_\perp q_2}{m_\parallel q_1 + m_\perp q_2}, \quad b = \frac{2m_\parallel q_1}{m_\parallel q_1 + m_\perp q_2} . \tag{21}$$

But q_1 and q_2 are connected by the relation that expresses the laws of conservation of energy and of p_t:

$$q_1^2 - q_2^2 = (m_\parallel/m_\perp - 1)r^2. \tag{22}$$

Knowledge of the coefficient a enables us to determine the coefficient R_1:

$$R_1 = |a|^2 = \left|\frac{m_\parallel q_1 - m_\perp q_2}{m_\parallel q_1 + m_\perp q_2}\right|^2. \tag{23}$$

The coefficient R_1 can of course be determined also by use of the coefficient b.

In accordance with what was said above, electrons located inside the cone $q_1^2 - (m_\parallel/m_\perp - 1)r^2 = 0$ (for them $q_1^2 - (m_\parallel/m_\perp - 1)r^2 < 0$) undergo total internal reflection at the boundary of the crystallites. For them, $R_1 \equiv 1$. By considering incidence of an electron wave from the right on the boundary of the crystallites, we can show that R_2 is also determined by formula (23). In calculating characteristics related to the first crystallite, we must by means of (22) express q_2 in terms of q_1, while in calculating characteristics related to the second crystallite we must, conversely, express q_1 in terms of q_2.

We note that electrons of the second crystallite do not

undergo total internal reflection. According to (11), (15), (22), and (23), we have for $\tau_1 = \tau_2$

$$\xi_1 = -\frac{4m_{\parallel}m_{\perp}Q_1}{(m_{\parallel}+m_{\perp}Q_1)^2}(1-Q_1),$$

$$\xi_2 = \frac{4m_{\parallel}m_{\perp}Q_2}{(m_{\perp}+m_{\parallel}Q_2)^2}(Q_2-1); \qquad (24)$$

$$Q_{1,2}=\left[1\mp\left(\frac{m_{\parallel}}{m_{\perp}}-1\right)\frac{r^2}{q_{1,2}^2}\right]^{1/2}.$$

We recall that by hypothesis $m_{\parallel} > m_{\perp}$. Therefore $\xi_2 \neq 0$ for all r and q, whereas $\xi_1 \neq 0$ only for those values of the quasimomentum for which $(m_{\parallel}/m_{\perp}-1)\cdot(r^2/q^2)<1$ (see above). We note that ξ_1 and ξ_2 are of opposite signs, whence it follows that the electromechanical force density in the two crystallites also has opposite signs (Fig. 3). Calculation of $\langle \xi_1 \rangle$ and $\langle \xi_2 \rangle$ for arbitrary $m_{\parallel}/m_{\perp}$ (see (26)) is possible only by numerical integration. We shall restrict ourselves to calculation of the limiting values for $m_{\parallel}/m_{\perp} \gg 1$ and for $m_{\parallel} \approx m_{\perp}$. If $m_{\parallel} \gg m_{\perp}$, then

$$\langle\xi_1\rangle\approx-\tfrac{1}{15}(m_{\perp}/m_{\parallel})^4, \quad \langle\xi_2\rangle\approx m_{\perp}/m_{\parallel}, \qquad (25)$$

and for $m_{\parallel} \approx m_{\perp}$,

$$\langle\xi_1\rangle\approx-10^{-2}\left(\frac{\Delta m}{m}\right)^{1/2}, \quad \langle\xi_2\rangle\approx0,7\left(\frac{\Delta m}{m}\right)^{1/2}, \qquad (26)$$

$$\Delta m = m_{\parallel}-m_{\perp}, \; m = m_{\parallel} \approx m_{\perp}.$$

The estimate of $\langle \xi_s \rangle$ shows that the electromechanical force density, as was to be expected, is smaller in the case of a specular boundary than in the case of a diffuse one.

The sensitivity of $\mathscr{F}(z)$ to the electronic energy spectrum of the metal should be noted. The electromechanical forces naturally have their maximum value on a diffuse crystallite interface ($\mathscr{F}_{surf}$). It is convenient to express this force (see (13)) in terms of the current density. On omitting dimensionless factors (usually close to unity), we get

$$\mathscr{F}_{surf}\approx jp_F/e,$$

whence $\mathscr{F}_{surf}$ [dyn/cm^2] $\approx j$ [A/cm^2]. In some experiments the current density reaches $\approx 10^8$ A/cm^2; in such cases $\mathscr{F}_{surf} \approx 10^8$ dyn/cm^2.

[1] In the case of two different metal specimens with diffusely scattering surfaces, the boundary conditions would coincide with (5).

[1] V. B. Fiks, Ionnaya provodimost' v metallakh i poluprovodnikakh (Ionic Conduction in Metals and Semiconductors), Nauka, 1969.

[2] O. A. Troitskiĭ, Pis'ma Zh. Eksp. Teor. Fiz. 10, 18 (1969) [JETP Lett. 10, 11 (1969)].

[3] V. I. Spitsyn, O. A. Troitskiĭ, V. G. Ryzhkov, and A. S. Kozyrev, Dokl. Akad. Nauk SSSR 231, 402 (1976).

[4] K. M. Klimov, G. D. Shnyrev, and I. I. Novikov, Dokl. Akad. Nauk SSSR 219, 323 (1974) [Sov. Phys. Dokl. 19, 787 (1975)].

[5] V. M. Kontorovich, Zh. Eksp. Teor. Fiz. 59, 2116 (1970) [Sov. Phys. JETP 32, 1146 (1971)].

Translated by W. F. Brown, Jr.

The electrotropic cholesteric-nematic transition and the storage effect based on it

V. N. Chirkov, D. F. Aliev, and A. Kh. Zeĭnally

Azerbaĭdzhan State University
(Submitted March 30, 1977)
Zh. Eksp. Teor. Fiz. 73, 761–766 (August 1977)

The threshold, contrast, and time characteristics of the cholesteric-nematic phase transition induced by an electric field have been studied in compensated mixtures of cholesteric liquid crystals with nematic-liquid-crystal admixtures (up to 10%). A storage effect has been observed which consists in the prolonged preservation of the homeotropic orientation of the molecules near the cholesteric-nematic transition temperature. Judging by the principal electro-optical parameters, the investigated system surpasses the well-known nematic-cholesteric system in which the storage effect is due to texture transformations under conditions of dynamic scattering of light.

The twistedness of the hypomolecular structure determines the unusually strong optical activity of the cholesteric liquid crystal (CLC). In this case the plane of polarization of polarized light passing through the CLC may rotate either clockwise or anticlockwise, depending on the structure of the individual molecules of the cholesteric compound. By mixing in definite concentration proportions dextro- and levo-rotatory CLC, we can compensate the optical activity.[1] The compensation occurs at a definite—for each mixture—temperature