Contribution to the theory of kinetic phenomena in Peierls dielectics

S. N. Artemenko and A. F. Volkov

Institute of Radio Engineering and Electronics, USSR Academy of Sciences
(Submitted 20 October 1980)
Zh. Eksp. Teor. Fiz. **80**, 2018–2030 (May 1980)

Equations are obtained for the Green's functions that describe the kinetics of a quasi-one-dimensional conductor with a charge-density wave [Peierls dielectric (PD)]. The equations are similar to those of the kinetics of superconductors and are valid when one-dimensional fluctuations can be disregarded (for example, as a result of three-dimensionality effects). The low-frequency conductivity of PD is calculated in two limiting cases: a) in the gapless state, at $\Delta \ll \nu \ll \eta$ (Δ is the order parameter, ν is the reciprocal momentum-scattering time, and η is the characteristic curvature of the Fermi surface of the quasi-one-dimensional conductor); b) in a pure PD ($\nu \ll \eta \ll \Delta$). The natural oscillations of the phase of the charge-density wave are investigated. It is shown that these oscillations exist in the form of a "soft" mode near the critical temperature and in the form of a "hard" mode at low temperatures.

PACS numbers: 72.15.Nj

It is known that when the temperature is lowered an instability of the Peierls type appears in quasi-one-dimensional conductors (metals and semimetals), and that a structural phase transition takes place near a certain temperature T_P. The initial lattice is then restructured and a superlattice appears in the crystal. In many quasi-one-dimensional conductors, the appearance of the superlattice is accompanied by formation of an electron charge-density wave (CDW) (see, e.g., Ref. 1). The instability sets in if the equality $\varepsilon(\mathbf{p}+\mathbf{Q})+\varepsilon(\mathbf{p})=0$ is satisfied in a certain region of variation of the electron momentum near the Fermi surface. Here $\mathbf{Q}$ is fixed and is the wave vector of the CDW. The energy $\varepsilon(\mathbf{p})$ is reckoned from the Fermi energy ε_F. As a result of the instability, a gap 2Δ appears in the excitation spectrum of the conductor, i.e., a transition to the dielectric state takes place. Such a conductor is called a Peierls dielectric (PD). By itself, the transition into the PD state recalls a superconducting transition of a metal. In the former case, however, a more important role is played by fluctuations, which lead to a certain smearing of the transition. If, however, the spectrum of the electrons (or phonons) is not purely one-dimensional then, just as in a superconducting transition in a metal, the role of the fluctuations is small, and the transition can then be described by using the self-consistent-field approximation.[2]

A most interesting circumstance is that, just as in a superconductor, the response of a PD to an electric field E is due not only to quasiparticles, but also to a "condensate," the role of which in a PD is played by a CDW.[3,4] The difference from a superconductor lies in the fact that the contribution made to the conductivity by this additional mechanism, called the Fröhlich mechanism, can be equal to zero if the CDW motion is blocked by pinning. The latter is due to lattice defects or to the effect of commensurability of the periods of the superlattice in the initial lattice. Convincing experimental proof has been obtained of the contribution of the CDW to the conductivity and of the influence of pinning on this effect[5,6] (see also the references therein).

The contribution of the CDW to the PD conductivity was theoretically analyzed in a number of papers.[4,7,8] It was concluded in Ref. 8 that to describe correctly the conductivity of a PD with allowance for the quasiparticles in the CDW it is necessary to go outside the framework of the two-liquid model and to analyze the kinetic equations. It was noted in Ref. 7 that the usual kinetic equation is not valid in this case, and the situation is analogous with that encountered in the development of the theory of nonequilibrium processes in superconductors. Two approaches were used, in the main, in the microscopic theory of kinetic phenomena in superconductors. In the first, the response of the superconductor to an external action was determined by expanding the temperature Green's functions in powers of the field E and of the order parameter Δ and by summing the corresponding diagrams,[9] followed by analytic continuation to the real-frequency axis. In the second approach, the Keldysh technique for the investigation of nonequilibrium processes, generalized to the case of superconductors,[10,11] was used. The last approach offers a number of advantages. First, by using the strong dependence of the Green's function $G(\varepsilon,\varepsilon',\mathbf{p},\mathbf{R})$ on the energy near the Fermi surface it was possible to obtain equations for Green's functions that are integrated with respect to the variable $\xi_p=v(p-p_F)$ and depend therefore on a smaller number of variables: $g=g(\varepsilon,\varepsilon',\theta,\varphi,\mathbf{R})$. These equations can be written in the form of a single matrix equation. Second, the obtained normalization relations simplify greatly the solution of the equations.[11] This approach made it possible to solve a number of interesting problems of the theory of nonequilibrium processes in superconductors (see, e.g., the review[12]).

The method of analytically continuing the responses was used by Gor'kov and Dolgov[7] to calculate the static conductivity of a PD near the transition temperature T_P. In the present paper we analyze a number of kinetic phenomena by another approach developed in superconductivity theory. We obtain general expressions for the matrix Green's functions $\check{g}$ that describe the PD. From these equations we obtain the linear response of the system to an external electric field $E(\omega, k)$. The static conductivity σ of the system is determined with allowance for electron scattering by impurities and for the motion of the CDW, with pinning neglected. In addition, we analyze the spectrum of the collective excitations of the PD and show that near T_P, under definite conditions, there exists a soft mode of such excitations, $\omega \sim kV$, while at low temperatures the collective excitations exist in the form of a hard mode, $\omega^2 \sim \omega_0^2 + k^2V^2$.

1. BASIC EQUATIONS

We use the model of Gor'kov and Dolgov as the basis.[7] We consider a quasi-one-dimensional metal whose spectrum is described by the formula

$$\varepsilon_p = p_\parallel^2/2m + \varphi(\mathbf{p}_\perp) - \varepsilon_F, \tag{1}$$

and assume that $|\varphi(\mathbf{p}_\perp)| \ll \varepsilon_F$, i.e., we consider a quasi-one-dimensional metal whose two flat Fermi surfaces $|p_\parallel^2 = \text{const}|$ are only weakly bent. This bending is described by the function $\varphi(\mathbf{p}_\perp)$. The three-dimensionality effects described by $\varphi(\mathbf{p}_\perp)$ make it possible, in particular, to disregard the one-dimensional fluctuations. In addition, if the characteristic energy of the transverse motion $\sim|\varphi(\mathbf{p}_\perp)|$ is high enough $[|\varphi(\mathbf{p}_\perp)| > \tau^{-1}$, where τ is the momentum relaxation time in scattering by the impurities], then the electron scattering can be treated by the usual cross technique, and the fan diagrams, allowance for which leads to Anderson localization in a disordered potential,[13] can be disregarded.

The Hamiltonian of the system, in the self-consistent field approximation and in the absence of external fields and impurities, is of the form

$$\hat{H} = \sum_p \varepsilon_p c_p^+ c_p + \sum_q \omega_q b_q^+ b_q + \sum_k (c_{k+Q_0}^+ c_{k-Q_0} \Delta + c_{k-Q_0}^+ c_{k+Q_0} \Delta^*), \tag{2}$$

where $\Delta = g\langle b_Q + b_{-Q}^+\rangle$ is the order parameter, $2\mathbf{Q}_0 = \mathbf{Q}$ is the CDW wave vector, g is the electron-phonon interaction constant, and ω_q is the phonon frequency. It is assumed that $|\mathbf{k}| \ll \mathbf{Q}_0$. The onset of a superlattice, i.e., the appearance of Δ, entangles the states on different parts of the Fermi surface: $\mathbf{p}_+ = \mathbf{k} + \mathbf{Q}_0$ and $\mathbf{p}_- = \mathbf{k} - \mathbf{Q}_0$. From the equations of motion of the operators c_p^+ and c_p it is easy to obtain equations for the matrix Green's functions. One of these functions, introduced by Keldysh

$$G_{\alpha\beta} = i^{-1}\langle c_{k+\alpha Q_0}(t) c_{k'+\beta Q_0}^+(t') - c_{k'+\beta Q_0}^+(t') c_{k+\alpha Q_0}(t)\rangle, \quad \alpha, \beta = \pm 1, \tag{3}$$

describes the kinetics of the system, and the two others (retarded and advanced) describe the spectrum of the system. The equations for the functions G and $G_{R(A)}$ are of the form

$$\{i\partial/\partial t - (\hat{\varepsilon}_k + \hat{\Delta})\}\hat{G} \equiv \hat{G}_0^{-1}\hat{G} = 0, \tag{4}$$

$$\hat{G}_0^{-1}\hat{G}^{R(A)} = \hat{1}\delta(t-t'), \tag{5}$$

$$\hat{\varepsilon}_k = \begin{pmatrix} \varepsilon(\mathbf{k}+\mathbf{Q}_0) & 0 \\ 0 & \varepsilon(\mathbf{k}-\mathbf{Q}_0) \end{pmatrix}, \quad \hat{\Delta} = \begin{pmatrix} 0 & \Delta \\ \Delta^* & 0 \end{pmatrix}.$$

In the stationary case, the Green's functions depend on the difference $t - t'$. From (5) we obtain for the Fourier component $\hat{G}_R(\varepsilon)$ the known expression[7,8]:

$$G^R(\varepsilon) = [(\tilde{\varepsilon}+i0)^2 - \xi^2 - |\Delta|^2]^{-1} \begin{pmatrix} \tilde{\varepsilon}+\xi & \Delta \\ \Delta^* & \tilde{\varepsilon}-\xi \end{pmatrix},$$

$$\tilde{\varepsilon} = \varepsilon - \eta(\mathbf{p}_\perp), \quad \eta(\mathbf{p}_\perp) = [\varepsilon(\mathbf{p}_+) + \varepsilon(\mathbf{p}_-)]/2, \quad \xi^2 = \eta^2 - \varepsilon(\mathbf{p}_+)\varepsilon(\mathbf{p}_-). \tag{6}$$

If the vector $\mathbf{Q}$ is parallel to the filaments ($\mathbf{Q} \parallel \mathbf{P}_\parallel$) we have

$$\eta = Q_0^2/2m - \varepsilon_F + \varphi(\mathbf{p}_\perp), \quad \xi = (Q_0/m)k = vk. \tag{7}$$

The PD spectrum is thus described by the formula

$$\varepsilon = \eta \pm (|\Delta|^2 + \xi^2)^{1/2}.$$

We consider now elastic scattering by the impurities. If $|\varphi(\mathbf{p}_\perp)| \gg \tau^{-1}$ we can use the usual cross technique. In addition, we disregard diagrams that are odd in the impurity potential and describe the pinning.[14] Neglect of pinning means that either the field E exceeds the threshold CDW collapse field, or that the frequency of the field E is high enough (see Ref. 7). We then have for the impurity self-energy part

$$\Sigma_{\alpha\beta}\left(\mathbf{k}+\frac{\mathbf{q}}{2}, \mathbf{k}-\frac{\mathbf{q}}{2}\right) = N_i \int d\mathbf{q}_1 \int \frac{d\mathbf{k}_1}{(2\pi)^3} \left| u\left(\mathbf{Q}_0(\alpha-\alpha') + \mathbf{k} - \mathbf{k}' + \frac{\mathbf{q}+\mathbf{q}_1}{2}\right)\right|^2 \times G_{\alpha\beta}(\mathbf{k}_1\mathbf{q}_1)\delta[\mathbf{Q}_0(\alpha-\alpha_1-\beta+\beta_1) + \mathbf{q} - \mathbf{q}_1], \tag{8}$$

where N_i is the impurity density and u is the potential of the interaction with the impurities. The wave vectors q and q_1 characterize the spatial inhomogeneity of the perturbations, so that q^{-1} is the characteristic scale of the inhomogeneity. Since we are considering perturbations that are smooth over interatomic distances, we can neglect q and q_1 compared with Q. It follows therefore from (8) that $\alpha - \alpha_1 - \beta + \beta_1 = 0$.

Neglecting in the interaction potential q, q_1 and $k_\parallel$, $k_{1\parallel}$ compared with Q, we can rewrite (8) in the form

$$\hat{\Sigma}(\mathbf{k}_\parallel, \mathbf{k}_\perp; \mathbf{R}) = \frac{1}{2\pi}\int d\xi \int \frac{d\mathbf{k}_\perp'}{S} \left\{ \nu_1(\mathbf{k}_\perp - \mathbf{k}_\perp')\hat{G}(\mathbf{k}_\parallel, \mathbf{k}_\perp'; \mathbf{R}) + \frac{\nu_2(\mathbf{k}_\perp - \mathbf{k}_\perp')}{2}[\hat{\sigma}_x\hat{G}(\mathbf{k}_\parallel, \mathbf{k}_\perp'; \mathbf{R})\hat{\sigma}_x + \hat{\sigma}_y\hat{G}(\mathbf{k}_\parallel, \mathbf{k}_\perp'; \mathbf{R})\hat{\sigma}_y]\right\}. \tag{9}$$

We have taken here a Fourier transform in the coordinate representation with respect to the summary coordinate $\mathbf{R} = (\mathbf{r} + \mathbf{r}')/2$, and S is the area of the intersection of the Brillouin zone with the plane $\mathbf{k}_\parallel = 0$. The collision frequencies ν_1 and ν_2 are determined by the forward and backward scattering of the particles (with a change $\mathbf{Q}$ in the value of the momentum). They are expressed in terms of the potential u with the aid of the equations

$$\nu_1(\mathbf{k}_\perp) = \frac{SN_i}{(2\pi)^2 v}|u(\mathbf{k}_\perp)|^2, \quad v = \frac{Q_0}{m}, \quad \nu_2(\mathbf{k}_\perp) = \frac{SN_i}{(2\pi)^2 v}|u(\mathbf{k}_\perp + \mathbf{Q})|^2. \tag{10}$$

Just as in the case of a superconductor, Σ does not depend of $\mathbf{k}_\parallel$, i.e., on ξ [accurate to terms $\sim\xi/\varepsilon_F \sim (\Delta, T)/\varepsilon_F$] and the functions $\hat{G}$ vary rapidly with changing ξ. Therefore, by eliminating ξ from the equations for $\hat{G}$, we can obtain equations for the Green's functions integgrated with respect to ξ:

$$\hat{g} = \frac{i}{\pi}\int d\xi \hat{G}(\xi, \mathbf{k}_\perp, \varepsilon, \varepsilon'). \tag{11}$$

Consider, for example, Eq. (4) for the functions $\hat{G}$. We write down this equation with allowance for the scattering of the particles by the impurities and for the presence of the electrostatic potential $\Phi(\mathbf{R},t)\equiv\Phi(x)$, as well as the conjugate equation

$$\hat{G}_0^{-1}(x)\hat{G}(x,x')-\Phi(x)\hat{G}=\hat{\Sigma}\hat{G},$$
$$\hat{G}(x,x')[\hat{G}_0^{-1}(x')]^*-\hat{G}\Phi(x')=\hat{G}\hat{\Sigma}. \tag{12}$$

We subtract these equations from each other, after multiplying the first by σ_z from the left and the second from the right. In the equation obtained in this manner we can integrate with respect to ξ. The result is an equation for the function $\hat{g}(\varepsilon,\varepsilon';\mathbf{R},\mathbf{k}_\perp)$:

$$\hat{\sigma}_z(\tilde{\varepsilon}-\hat{\Delta}(\mathbf{R}))\hat{g}-\hat{g}(\tilde{\varepsilon}'-\hat{\Delta}(\mathbf{R}))\hat{\sigma}_z+i(\hat{v}\nabla)\hat{g}$$
$$-\Phi(R)\hat{\sigma}_z\hat{g}+\hat{g}\hat{\sigma}_z\Phi(R)=\hat{\sigma}_z\hat{\Sigma}\hat{g}-\hat{g}\hat{\Sigma}\hat{\sigma}_z. \tag{13}$$

Here $\tilde{\varepsilon}'=\varepsilon'-\eta(\mathbf{k}_\perp)$

$$\hat{v}\nabla\hat{g}=\begin{pmatrix} v_+\nabla g_{++} & \frac{1}{2}(v_+-v_-)\nabla g_{+-} \\ \frac{1}{2}(v_+-v_-)\nabla g_{-+} & -v_-\nabla g_{--}\end{pmatrix},$$

$\mathbf{v}_\pm=\partial\varepsilon(\mathbf{p}_\pm)/\partial\mathbf{k}$ is the group velocity on the right- (left-) hand Fermi surface. The product of $\hat{\Delta}(\mathbf{R})\hat{g}$ and $\Phi(\mathbf{R})\hat{g}$ means the convolution

$$\hat{\Delta}(\mathbf{R})\hat{g}=\int\frac{d\omega}{2\pi}\hat{\Delta}_\omega(\mathbf{R})g(\varepsilon-\omega,\varepsilon'). \tag{14}$$

the right-hand side of (13) describes scattering by the impurities. The integral of the collisions with the phonons is of similar structure if the inelasticity of the scattering is neglected (see Ref. 7). A similar equation holds for the functions $\hat{g}^{R(A)}$. If the vector $\mathbf{Q}$ is parallel to the filaments ($\mathbf{Q}\parallel\mathbf{p}_\parallel$), then $\mathbf{v}_\pm=\mathbf{v}_\perp\pm\mathbf{v}_\parallel$, where $\mathbf{v}_\perp=\partial\varphi(\mathbf{k}_\perp)/\partial\mathbf{k}_\perp$, $\mathbf{v}_\parallel=v\mathbf{n}_\parallel$, and $\mathbf{n}_\parallel$ is a unit vector in the direction of the filaments. The gradient term can then be rewritten as follows:

$$\hat{v}\nabla g=\mathbf{v}\nabla_\parallel g+(\mathbf{v}_\perp\nabla_\perp)(\hat{\sigma}_z\hat{g}+\hat{g}\hat{\sigma}_z)/2. \tag{15}$$

We shall consider just this case $\mathbf{Q}\parallel\mathbf{p}_\parallel$. We write down the equation for the matrices $\hat{g}$ and $gR(A)$. It is more convenient, however, to introduce a new matrix of the Green's functions and of the order parameter

$$\hat{g}_n=\hat{\sigma}_z\hat{g},\quad \hat{g}_n^{R(A)}=\hat{\sigma}_z\hat{g}^{R(A)},\quad \hat{\Delta}_n=\hat{\sigma}_z\hat{\Delta}. \tag{16}$$

We introduce also the supermatrix gn in analogy with the procedure in superconductivity theory[11]:

$$\check{g}_n=\begin{pmatrix} \hat{g}_n^R & \hat{g}_n \\ 0 & \hat{g}^A\end{pmatrix}. \tag{17}$$

The equations for the Green's functions can then be written in the form of a single equation for the matrix g. This equation is easily obtained from (13) and from the analogous equations for $\hat{g}_{R(A)}$, with account taken of (15). It takes the form (we leave out the subscript n to simplify the notation)

$$\tilde{\varepsilon}\check{\sigma}_z\check{g}-\check{g}\check{\sigma}_z\tilde{\varepsilon}'+[\check{\Delta}-\Phi\check{\sigma}_z,\check{g}]_-+iv\nabla_\parallel\check{g}+\frac{i}{2}\mathbf{v}_\perp\nabla_\perp[\check{\sigma}_z,\check{g}]_+=[\check{\Sigma},\check{g}]_-; \tag{18}$$

$$\check{\Delta}=\begin{pmatrix}\hat{\Delta} & 0\\ 0 & \hat{\Delta}\end{pmatrix},\quad \hat{\Delta}=\begin{pmatrix}0 & \Delta\\ -\Delta^* & 0\end{pmatrix},\quad \check{\sigma}_z=\begin{pmatrix}\hat{\sigma}_z & 0\\ 0 & \hat{\sigma}_z\end{pmatrix},$$
$$\check{\Sigma}=\begin{pmatrix}\hat{\Sigma}^R & \hat{\Sigma}\\ 0 & \hat{\Sigma}^A\end{pmatrix},\quad [\check{a},\check{b}]_\pm=[\check{a}\check{b}\pm\check{b}\check{a}], \tag{19}$$
$$\hat{\Sigma}=-\frac{i}{2}\int\frac{d\mathbf{k}_\perp'}{S}\{v_1(\mathbf{k}_\perp-\mathbf{k}_\perp')\hat{\sigma}_z\hat{g}(\mathbf{k}_\perp')\hat{\sigma}_z-\frac{1}{2}v_2(\mathbf{k}_\perp-\mathbf{k}_\perp')$$
$$\times[\hat{\sigma}_x\hat{g}(\mathbf{k}_\perp')\hat{\sigma}_x+\hat{\sigma}_y\hat{g}(\mathbf{k}_\perp')\hat{\sigma}_y]\}.$$

If all the functions $\check{g}$ depend only on the longitudinal coordinate $R_\parallel$, then the arguments advanced by Larkin and Ovchinnikov[11] can be used to obtain the normalization relation

$$\int\frac{d\varepsilon_1}{2\pi}\check{g}(\varepsilon,\varepsilon_1;R_\parallel)\check{g}(\varepsilon_1,\varepsilon';R_\parallel)=2\pi\check{1}\delta(\varepsilon-\varepsilon'). \tag{20}$$

Equation (20) must be supplemented by the self-consistency condition. We write down the equation for the evolution of the operators b_Q and b^+_{-Q} using the interaction Hamiltonian

$$H_{ep}=g\sum_{\mathbf{k},\mathbf{q}}(c^+_{\mathbf{k}+\mathbf{q}}c_{\mathbf{k}-\mathbf{q}}b_{2\mathbf{q}}+\text{c.c.}).$$

Then

$$\left(i\frac{\partial}{\partial t}-\omega_Q\right)b_Q=g\sum_{\mathbf{k}}c^+_{\mathbf{k}-Q_0}c_{\mathbf{k}+Q_0},$$
$$-\left(i\frac{\partial}{\partial t}+\omega_Q\right)b^+_{-Q}=g\sum_{\mathbf{k}}c^+_{\mathbf{k}-Q_0}c_{\mathbf{k}+Q_0}.$$

From this we get

$$\left(1+\frac{1}{\omega_Q^2}\frac{\partial}{\partial t^2}\right)\Delta=\lambda\int\frac{d\mathbf{k}_\perp}{S}\int d\varepsilon g_{+-}(\varepsilon,t;R). \tag{21}$$

Here $\lambda=g^2S/2\pi\omega_Q v$. An equation for Δ^* can be obtained analogously.

We now write down equations suitable for the determination of the modulus and phase of the order parameter $\Delta=|\Delta|e_{i\chi}$

$$|\Delta|\cos\chi=-\frac{\lambda}{4(1-\omega^2/\omega_Q^2)}\int\frac{d\mathbf{k}_\perp}{S}\int d\varepsilon\,\mathrm{Sp}(i\hat{\sigma}_y\hat{g}), \tag{22}$$
$$i|\Delta|\sin\chi=\frac{\lambda}{4(1-\omega^2/\omega_Q^2)}\int\frac{d\mathbf{k}_\perp}{S}\int d\varepsilon\,\mathrm{Sp}(\hat{\sigma}_x\hat{g}). \tag{23}$$

The current is also expressed in terms of the function $\hat{g}$. For the current along the filaments we have

$$I=\frac{Nv^2}{8}\int\frac{d\mathbf{k}_\perp}{S}\int d\varepsilon\,\mathrm{Sp}\,\hat{g}, \tag{24}$$

where $N=2S/v$ is the electron density of states of the PD at $T\geqslant T_P$.

Equation (18) together with the normalization condition (20) and the self-consistency conditions (22) and (23) describes the kinetics of the PD in the self-consistent-field approximation with pinning disregarded. These equations will be used to solve a number of problems of PD kinetics.

2. LOW FREQUENCY CONDUCTIVITY OF THE SYSTEM

We calculate on the basis of the derived equations the low-frequency ($\omega\ll\nu,\Delta$) conductivity of a PD with allowance for the CDW motion. We consider two limiting cases: a) high impurity density and large curvature of the Fermi surface ($\Delta\ll\nu\ll\eta$). b) low impurity density and weak bending of the Fermi surface ($\nu\ll\eta\ll\Delta$). The first of the considered cases corresponds to a gapless PD.

a) To find the system conductivity σ in this case it is necessary to specify the concrete functions $\eta(\mathbf{k}_\perp)$ and $\nu(\mathbf{k}_\perp)$. We assume for simplicity that in the main region of variation of $\mathbf{k}_\perp$ the functions $\nu_{1,2}$ do not depend on $\mathbf{k}_\perp$. The right-hand side of (18) can then be rewritten in the form

$$-i\nu_1(\check\sigma_z\langle\check g\rangle\check\sigma_z\check g-\check g\check\sigma_z\langle\check g\rangle\sigma_z)/2+i\nu_2(\check\sigma_x\langle\check g\rangle\check\sigma_x\check g-\check g\check\sigma_x\langle\check g\rangle\check\sigma_x+\check\sigma_y\langle\check g\rangle\check\sigma_y\check g-\check g\check\sigma_y\langle\check g\rangle\check\sigma_y)/4, \tag{25}$$

$$\langle g\rangle=S^{-1}\int d\mathbf{k}_\perp g.$$

We find first the functions $\hat g_{R(A)}$, which we represent in the form $\hat g_{R(A)}=g_{R(A)}\sigma_z+f_{R(A)}i\sigma_y$. From (18) and (20) we obtain, taking (25) into account,

$$-2\varepsilon f^{R(A)}+(2\Delta-i\nu_1\langle f^{R(A)}\rangle)g^{R(A)}-i\nu_+f^{R(A)}\langle g^{R(A)}\rangle=0, \tag{26}$$

$$(g^{R(A)})^2-(f^{R(A)})^2=1; \tag{27}$$

where $\nu_+=\nu_1+\nu_2$.

Equations (26) and (27) differ from the corresponding equations for superconductors with magnetic impurities only in the presence of the function η. The role of the frequency τ_s^{-1} of collisions with spin flip is assumed in this case by the quantity $\nu_0=\nu_1+\nu_{2/2}$. Putting $\eta=0$, we can use the known results of superconductivity theory. In particular, at $\Delta\tau_s<1$ (correspondingly, $\Delta/\nu_0<1$), a gapless state takes plece. This analogy with superconductors was noted repeatedly (see, e.g., the review by Bulaevskiĭ[1]). Strictly speaking, however, one cannot put $\eta=0$ at $\nu\neq0$, for it is then necessary to take into account localization effects in the one-dimensional system. Allowance for the curvature of the Fermi surface also leads to a suppression of the critical temperature T_p. If $\eta>\nu$, then the influence of the curvature on the PD spectrum and on T_p predominates, and the influence of the impurities is negligile in this case (in accord with the ratio ν/η). In fact, let us determine $f_{R(A)}$ from (26) and (27) by perturbation theory, regarding Δ/ν as a small parameter. We obtain

$$f^R=\frac{\Delta}{(\varepsilon+i\nu_+/2)}\left[1+\frac{i\nu_1}{2}\langle(\varepsilon+i\nu_+/2)^{-1}\rangle\right]^{-1},$$

$$f^A=-f^{R*},\quad N_p(\varepsilon)=\frac{1}{2}\langle g^R-g^A\rangle=1+\frac{1}{4}\langle(f^R)^2+(f^A)^2\rangle. \tag{28}$$

It follows therefore that the PD state density N_p does not vanish at any value of ε. If we neglect the $\eta(k_\perp)$ dependence, then

$$f^{R(A)}=\pm\frac{\Delta}{\varepsilon\pm i\nu_0}.$$

From the self-consistency condition (22) we obtain for the transition temperature T_P

$$1=\frac{\lambda}{2}\int_0^{\varepsilon_c}d\varepsilon\frac{\varepsilon}{\varepsilon^2+\nu_0^2}\left(\operatorname{th}\frac{\varepsilon+\eta}{2T}+\operatorname{th}\frac{\varepsilon-\eta}{2T}\right),$$

where ε_c is an energy on the order of the width of the electron band. It is seen that at $\nu\ll\eta$ the $T_P(\nu)$ dependence is weak, since the main contribution to the integral is due to high energies. If $\nu\ll\eta\ll T_P$ we obtain from this formula

$$T_P=T_{P_0}(1-a\eta^2/T_{P_0}^2),\quad T_{P_0}=\varepsilon_c e^{-1/\lambda},\quad a\sim1.$$

We determine now the linear response of a gapless PD to a constant electric field E. Just as in superconductivity theory,[9,12] it is convenient to represent the perturbed Green's function $\hat g$ in the form of the sum

$$\delta\hat g=\hat g^{(r)}+\hat g^{(a)},\quad \hat g^{(r)}=\delta\hat g^R(\varepsilon,\varepsilon')\operatorname{th}\frac{\varepsilon'}{2T}-\delta\hat g^A(\varepsilon,\varepsilon')\operatorname{th}\frac{\varepsilon}{2T}. \tag{29}$$

Here $\hat g^{(r)}$ and $\hat g^{(a)}$ are respectively the regular and anomalous parts. The latter differs from zero in the nonequilibrium case and has complicated analytic properties. Since (18) contains the combination $\varepsilon-\Phi$, we seek the total Green's functions in the form

$$\check g(\varepsilon)=\check g_0(\varepsilon-\Phi)+\delta\check g,$$

$$g_0(\varepsilon)=(\hat g_0^R-\hat g_0^A)\operatorname{th}\frac{\varepsilon}{2T}, \tag{30}$$

$\delta\check g$ is defined by (29).

The function (30) must be substituted in the linearized equation (18) and written for the compoents $\delta\hat g^{R(A)}$ and $\hat g^{(a)}$. The latter take in this case the form $\hat g^{(a)}=g_1^{(a)}\hat 1+g_x^{(a)}\hat\sigma_x$ (the same holds for $\delta\hat g^{R(A)}$). It must be borne in mind that the problem is not entirely static, since the condensate current is proportional to $\partial\chi/\partial t$. The limit as $\omega\to0$ must therefore be taken with caution. For g_1^R and g_x^R we obtain the equations

$$-\omega g_1^R-2\Delta g_x^R+\frac{i\nu_1}{2}\langle f_+^R+f_-^R\rangle g_x^R=\frac{i\nu_1}{2}(f_+^R+f_-^R)\langle g_x^R\rangle-i\Delta\chi(f_+^R+f_-^R)-iv\frac{\partial g^R}{\partial\varepsilon}E,$$

$$\frac{i\nu_1}{2}\langle f_+^R-f_-^R\rangle g_1^R-2\varepsilon g_x^R-\frac{i\nu_+}{2}\langle g_+^R+g_-^R\rangle g_x^R=\frac{i\nu_1}{2}(g_+^R+g_-^R)\langle g_x^R\rangle-\frac{i\nu_-}{2}(f_+^R-f_-^R)\langle g_1^R\rangle-i\Delta\chi(g_+^R+g_-^R)-iv\frac{\partial f^R}{\partial\varepsilon}E, \tag{31}$$

where $f_\pm^R=f^R(\varepsilon\pm\omega/2)$ and $\nu_\pm=\nu_1\pm\nu_2$. Hence

$$\langle g_x^R\rangle=i\chi\langle f^R\rangle-\frac{iEv}{2\Delta}\left\langle\frac{\partial g^R}{\partial\varepsilon}\right\rangle,\quad g_1^R=\frac{1}{2}\frac{\partial g^R}{\partial\varepsilon}\frac{\partial\chi}{\partial t}. \tag{32}$$

$\hat g^{(a)}$ can be obtained similary

We now write the final expressions for the functions δg_1 and δg_x, which determine the current and the phase of the order-parameter

$$\langle\delta g_1\rangle=\left[vE+\frac{i\Delta}{2}\langle f^R+f^A\rangle\left(\frac{\partial\chi}{\partial t}-\frac{vE}{\nu_2}\right)\right]\left(2T\nu_2\operatorname{ch}^2\frac{\varepsilon}{2T}\right)^{-1}+\frac{1}{2}\frac{\partial\chi}{\partial t}\operatorname{th}\frac{\varepsilon}{2T}\frac{\partial\langle g^R-g^A\rangle}{\partial\varepsilon},$$

$$\langle g_x\rangle=i\chi\langle f^R-f^A\rangle\operatorname{th}\frac{\varepsilon}{2T}+\frac{iEv}{2\Delta}\frac{\partial\langle g^R-g^A\rangle}{\partial\varepsilon}\operatorname{th}\frac{\varepsilon}{2T}+\frac{\langle f^R+f^A\rangle}{4T\operatorname{ch}^2(\varepsilon/2T)}\left(\frac{\partial\chi}{\partial t}-\frac{vE}{\nu_2}\right). \tag{33}$$

Substituting $\langle g_1\rangle$ in (24), we determine the current

$$I=\sigma_N E(1-a)+\sigma_N\left(\frac{\nu_2}{v}\right)\frac{\partial\chi}{\partial t}\left(a-\frac{b}{2}\right), \tag{34}$$

where σ_N is the conductivity in the normal state

$$a=\int d\varepsilon\frac{i\Delta\langle f^R+f^A\rangle}{8T\nu_2\operatorname{ch}^2(\varepsilon/2T)},\quad b=-\int\frac{d\varepsilon}{2}\operatorname{th}\frac{\varepsilon}{2T}\frac{\partial\langle g^R-g^A\rangle}{\partial\varepsilon}.$$

If the CDW is slowed-down, then

$$\frac{\partial\chi}{\partial t}=0\text{ and }I=\sigma_N E(1-a).$$

But if the pinning is inessential and the CDW moves, then $\partial\chi/\partial t$ must be determined from the self-consistency condition (23), which yields

$$a\left(\frac{\partial\chi}{dt}-\frac{vE}{\nu_2}\right)+\frac{bv}{2\nu_2}E=0. \tag{35}$$

Determining the derivative $\partial\chi/\partial t$ from (35) and substituting it in (34), we get

$$j=\sigma_N E[1-b(1-b/4a)].$$

Let us estimate a and b. In the case $\nu \ll \eta \ll T$ considered by Gor'kov and Dolgov[7] we obtain

$$a \sim \Delta^2/\nu_2 T, \quad b \sim -\Delta^2/T^2.$$

At $\nu \ll \eta$ and $T \ll \eta$, putting for example $\eta(\mathbf{k}_\perp) = \eta_0 \cos(\mathbf{a}\cdot\mathbf{k}_\perp)$, which is typical of the tight-binding approximation of electrons on a filament, we obtain $a \sim \Delta^2/\nu_2\eta_0$ and $b \sim -\Delta^2/\eta_0^2$. In both cases $|b| \ll a$, i.e., in a gapless PD the CDW motion compensates for the decrease of the conductivity of the free electrons and leads to a negligible increase of the conductivity.

b) We determine now the conductivity of a quasi-one-dimensional conductor with a low impurity density ($\nu \ll \Delta$). We confine ourselves to the simplest case of weak bending of the Fermi surface, $\eta \ll \Delta$ and $\eta \ll T$. Then the equilibrium properties depend neither on ν nor on η. In particular, it follows from (26) and (27) that the functions $g^{R(A)}$ and $f^{R(A)}$ are of the same form as in a superconductor

$$g^{R(A)} = (\varepsilon/\Delta) f^{R(A)} = \varepsilon/\xi^{R(A)},$$
$$\xi^{R(A)} = \pm(\varepsilon^2-\Delta^2)^{1/2}\,\mathrm{sign}\,\varepsilon\cdot\theta(|\varepsilon|-\Delta) + i(\Delta^2-\varepsilon^2)^{1/2}\theta(\Delta-|\varepsilon|). \tag{36}$$

Proceeding as in the previous case, we evaluate the functions g_1 and g_x that determine the current and the phase of the CDW:

$$g_1 = \frac{1}{2}\frac{\partial\chi}{\partial t}\,\mathrm{th}\frac{\varepsilon}{2T}\frac{\partial(g^R-g^A)}{\partial\varepsilon} + \frac{\theta(|\varepsilon|-\Delta)\,(vE\xi^2+\nu_0\Delta^2\varepsilon\xi^{-1}\partial\chi/\partial t)}{2T\,\mathrm{ch}^2(\varepsilon/2T)\,(\xi^2\nu_2+\Delta^2\nu_0)}, \tag{37}$$

$$g_x = \frac{ivE}{2\Delta}\,\mathrm{th}\,\frac{\varepsilon}{2T}\frac{\partial(g^R-g^A)}{\partial\varepsilon} + \frac{\theta(|\varepsilon|-\Delta)}{2T\,\mathrm{ch}^2(\varepsilon/2T)}\frac{i\nu_0\Delta\varepsilon\,(vE-\nu_2\varepsilon\xi^{-1}\partial\chi/\partial t)}{\xi\,(\xi^2\nu_2+\Delta^2\nu_0)}. \tag{38}$$

Here $\xi = (\varepsilon^2 - \Delta^2)^{1/2}$. With the aid of g_1 we obtain the current

$$I = \sigma_N\begin{cases} E\left(1-\dfrac{\Delta}{2T}A\right) + \dfrac{\pi}{4}\dfrac{\Delta}{T}\dfrac{(\nu_0\nu_2)^{1/2}}{v}\dfrac{\partial\chi}{\partial t}, & \dfrac{\Delta}{T} \ll 1 \\ E\dfrac{2\sqrt{\pi}\,\nu_2 T}{\nu_0\Delta}e^{-\Delta/T} + \dfrac{\nu_2}{v}\dfrac{\partial\chi}{\partial t}, & \dfrac{\Delta}{T} \gg 1 \end{cases}, \tag{39}$$

$$A = 1 + \frac{\nu_0}{\nu_2}\int_1^\infty \frac{dx}{x^2+\nu_0/\nu_2-1}.$$

It is seen from (39) that in the case of an immobile CDW ($\partial\chi/\partial t = 0$) the PD conductivity is close to the conductivity σ_N of a normal metal near T_P and is exponentially small at low temperatures.

In the case of a moving CDW we must find $\partial\chi/\partial t$ from the self-consistency condition (23). We note that the term $\partial\chi/\partial t$ in g_x makes a diverging contribution to the integral. We calculate it therefore with logarithmic accuracy, cutting off the integration at $\varepsilon - \Delta \sim \eta$, where the growth of the state density stops. Expressing $\partial\chi/\partial t$ in terms of E with the aid of the self-consistency condition, we obtain

$$\frac{\partial\chi}{\partial t} = \frac{Ev}{\nu_2}\begin{cases}\pi(\nu_0/\nu_2)^{1/2}[\ln(\Delta/\eta)]^{-1}, & \Delta/T \ll 1 \\ T\Delta^{-1}[\ln(T/\eta)]^{-1}e^{\Delta/T}, & \Delta/T \gg 1\end{cases}. \tag{40}$$

Substituting this result in expression (39) for the current, we get

$$I = \sigma_N E\begin{cases}1 + \dfrac{\Delta}{2T}\left(\dfrac{\pi^2}{2}\dfrac{\nu_0}{\nu_2\ln(\Delta/\eta)} - A\right), & \Delta/T \ll 1 \\ \dfrac{T}{\Delta}e^{\Delta/T}[\ln(T/\eta)]^{-1}, & \Delta/T \gg 1\end{cases}. \tag{41}$$

At $\Delta \ll T$ the contribution of the CDW to the conductivity is thus less than the conductivity decrease by the gap formation (since the logarithm is large by assumption). At low temperature, the PD conductivity is large by virtue of the motion of the CDW, and can exceed substantially the conductivity above T_P.

3. COLLECTIVE EXCITATIONS

In this section we determine the linear response of a PD to an alternating electric field $E_\sim = -ik\Phi = E_\parallel \exp(-i\omega t + ikx)$ and determine the spectrum of the collective excitations connected with the perturbation of the phase χ of the CDW (i.e., with the oscillations of the phase of the order parameter Δ) and of the longitudinal electric field $E_\sim$ (and correspondingly of the potential). To solve this problem we must linearize (18). We then obtain linear equations for the matrix functions $\delta g(\omega, k)$, with right-hand sides proportional to the potential Φ and to the phase χ of the CDW. These equations can be solved in the general case of arbitrary ω and k, as well as arbitrary relations between the parameters $\nu_{1,2}$ η, and Δ. We confine ourselves for simplicity to the case of a pure PD ($\nu_{1,2} = 0$) and disregard the curvature of the Fermi surfaces ($\eta \ll \Delta$).

We must calculate the function $\delta\hat{g}$ that determines the current in the system and the phase of the CDW [see Eqs. (23) and (24)]. It is again convenient to represent $\delta\hat{g}$ as a sum of regular and anomalous parts, $\hat{g} = \hat{g}^{(r)} + \hat{g}^{(a)}$. The equations for $\hat{g}^{(a)}$ can be easily obtained from (18). As for the equations for $\delta\hat{g}^{R(A)}$, they are obtained in the same manner as for superconductors, and have a structure similar to that of the equation for $\hat{g}^{(a)}$. The difference is that the right-hand sides of these equations do not contain the additional factor

$$\mathrm{th}\,(\beta\varepsilon_-) - \mathrm{th}\,(\beta\varepsilon_+) \quad (\beta = 1/2T,\ \varepsilon_\pm = \varepsilon \pm \omega/2),$$

and the exponent $A(R)$ is replaced by the exponent $R(A)$ in the equations for $g^{R(A)}$. The solution of the matrix equations for the functions $\delta\hat{g}^{R(A)}$ and $g^{(a)}$ is simplified if these functions are represented in the form

$$\delta\hat{g}^{R(A)} = \hat{g}_-^{R(A)}\hat{\alpha}^{R(A)} - \hat{\alpha}^{R(A)}\hat{g}_+^{R(A)}, \tag{42}$$

$$g^{(a)} = [\mathrm{th}\,(\beta\varepsilon_-) - \mathrm{th}\,(\beta\varepsilon_+)](g_-^R\hat{\alpha} - \hat{\alpha}g_+^A).$$

Here $\hat{g}_\pm^R = g_\pm^R\sigma_z + f_\pm^R i\sigma_y$ are the unperturbed functions. In this notation[11] the linearized normalization relations are automatically satisfied. It turns out that the functions $\hat{\alpha}$ are of the form $\alpha_z\hat{\sigma}_z + \alpha_1\hat{1}$, where

$$\alpha_z = \frac{-kv\Phi + i\chi\Delta(F_+/G_-)\omega}{(\xi_{\varepsilon_+}^R + \xi_{\varepsilon_-}^A)^2 - (kv)^2}$$

$$\alpha_1 = -\frac{\Phi(\omega - 2\Delta F_+/G_-) + kv\Delta(F_+/G_-)\chi}{(\xi_{\varepsilon_+}^R + \xi_{\varepsilon_-}^A)^2 - (kv)^2}, \tag{43}$$

$F_+ = f_+^R + f_-^A$, $G_- = g_+^R - g_-^A$. The expressions for $\hat{\alpha}^R$ have exactly the same form, if the functions with exponent A are replaced in (43) by functions with exponent R (for $\hat{\alpha}^A$, correspondingly, the functions with exponent R are replaced by functions with exponent A).

With the aid of the functions α_z and $\alpha_z^{R(A)}$ it is possible to obtain expressions for the current $I(\omega, k)$ and an equation for the phase $\chi(\omega, k)$. Linearizing Eq. (23) and substituting in it expressions (42) and (43) we obtain after simple transformations an equation for the phase:

$$i\chi+i\chi^{(r)}+i\chi^{(a)}=0,$$

$$i\chi^{(r)}=\frac{\lambda_\bullet}{4\Delta}\int d\varepsilon[\operatorname{th}(\beta\varepsilon_+)+\operatorname{th}(\beta\varepsilon_-)](F_+^R\alpha_z^R-F_+^A\alpha_z^A), \tag{44}$$

$$i\chi^{(a)}=\frac{\lambda_\bullet}{2\Delta}\int d\varepsilon[\operatorname{th}(\beta\varepsilon_+)-\operatorname{th}(\beta\varepsilon_-)][F_+\alpha_z-(F_+^R\alpha_z^R+F_+^A\alpha_z^A)/2].$$

From (24) we get an expression for the current

$$I=I^{(r)}+I^{(a)},$$

$$I^{(r)}=\frac{Nv^2}{16}\int d\varepsilon[\operatorname{th}(\beta\varepsilon_+)+\operatorname{th}(\beta\varepsilon_-)](G_-^R\alpha_z^R-G_-^A\alpha_z^A), \tag{45}$$

$$I^{(a)}=\frac{Nv^2}{8}\int d\varepsilon[\operatorname{th}(\beta\varepsilon_+)-\operatorname{th}(\beta\varepsilon_-)][G_-\alpha_z-(G_-^R\alpha_z^R+G_-^A\alpha_z^A)/2].$$

Calculating the integrals in (44) and (45), we can find the response of the PD to an alternating electric field and the spectrum of the collective excitations. We shall perform the calculation using expressions (43) and assuming for simplicity that the frequencies are low enough: $\omega \ll \Delta$ and $kv \ll \Delta$. We note that the regular parts of $\chi^{(r)}$ and $I^{(r)}$ can be determined by closing the contour of integration with respect to ε in the upper or lower half-plane and calculating the residues at the poles of the tangents. The anomalous parts of $\chi^{(a)}$ and $I^{(a)}$ are due to the quasiparticles; at low temperatures $T \ll \Delta$ they contain a small factor $\exp(-\Delta/T)$. We consider two limiting cases.

a) *Low temperatures*: $T \ll \Delta$.

In this case the anomalous terms $I^{(a)}$ and $\chi^{(a)}$ contain the small factor $\exp(-\Delta/T)$ and they can be neglected. Calculating the regular parts by residues, we obtain (we assume that $\lambda\omega_Q^2 \ll \Delta^2$)

$$i\chi\left(\frac{2\omega^2}{\lambda\omega_Q^2}-\frac{k^2v^2}{2\Delta^2}\right)=\frac{kv\Phi}{\Delta^2}\left(1-\frac{k^2v^2}{6\Delta^2}\right), \tag{46}$$

$$I=-\frac{N\omega}{4v}\left[\frac{kv\Phi}{3\Delta^2}\left(1-\frac{k^2v^2}{5\Delta^2}\right)-i\chi\left(1-\frac{k^2v^2}{6\Delta^2}\right)\right]. \tag{47}$$

We obtain the collective-excitation spectrum from (46) and from the quasineutrality condition $kI=0$, i.e., $I=0$. If we disregard the electric field $-ik\Phi$, then we obtain from (46) the soft mode: $\omega=\lambda^{1/2}\omega_Q kv/2\Delta$. Allowance for the electric field changes the soft mode into a hard one:

$$\omega^2=\frac{3}{2}\lambda\omega_Q^2+\frac{\lambda\omega_Q^2(kv)^2}{20\Delta^2}. \tag{48}$$

The first term coincides with the formula given in Ref. 4, and the second term describes the spatial dispersion.

We write down also an expression for the conductivity $\sigma(\omega)$ of the system in the long-wave limit:

$$\sigma(\omega)=-\left(\frac{Nv^2}{12\Delta^2}\right)\left[\frac{3\lambda\omega_Q^2}{2i\omega}+i\omega\right].$$

b) *Temperatures close to the critical* T_P: $\Delta \ll T$.

In this case the contribution to the current is due mainly to the quasiparticles, i.e., to the terms $I^{(a)}$ in (45). Calculating $I^{(a)}$ in the long-wave limit ($kv \ll \omega$), we obtain, assuming $\omega \ll \Delta$:

$$I=\frac{Nv^2}{2}\omega\left[\frac{kv\Phi}{\omega^2}\left(1-\frac{\Delta}{T}\right)+i\chi\frac{\pi}{8}\frac{\Delta}{T}\right]. \tag{49}$$

The equation for the phase (44) yields

$$i\chi\frac{\omega^2}{\omega_0^2}=kv\Phi\frac{\pi}{2\Delta T},\quad \omega_0^2=\frac{4\Delta T}{\pi}\left(1+\frac{8\Delta T}{\lambda\pi\omega_Q^2}\right)^{-1}. \tag{50}$$

Substituting (50) in (49) we obtain the conductivity of the PD near T_P:

$$\sigma_\bullet=-\frac{Nv^2}{2i\omega}\left[1+\frac{\pi\Delta}{2T}\left(1+\frac{8\Delta T}{\lambda\pi\omega_Q^2}\right)^{-1}-\frac{\Delta}{T}\right]. \tag{51}$$

The equation (50) for the phase holds if the ratio ω/kv is not small. Calculating the regular terms for the case when they play the decisive role in Eq. (44) for χ, under the condition

$$\omega/kv \ll \Delta/T, \tag{52}$$

we obtain the condition for the phase

$$i\chi\left[\frac{2\omega^2}{\lambda\omega_Q^2}+\frac{7\zeta(3)}{\pi^2}\frac{\omega^2-k^2v^2}{T^2}\right]=\frac{14\zeta(3)kv\Phi}{\pi^2T^2}. \tag{53}$$

In this case of temperatures close to T_P, the term in the right-hand side of (53) due to the influence of the field E, turns out to be small. We obtain accordingly near T_P a soft mode

$$\omega^2=\frac{7\zeta(3)}{2\pi^2}\frac{\lambda\omega_Q^2}{T^2}k^2v^2, \tag{54}$$

which attenuates weakly under the condition [cf. (52)]

$$\lambda\omega_Q^2 \ll \Delta^2. \tag{55}$$

The situation with the collective excitations in PD near the critical temperature is similar to the superconducting case.[12] In both cases the characteristic field E, which causes the dissipative condensate-current-compensating quasiparticle current, turns out to be small by virtue of the large number of the quasiparticle. It leads therefore to small damping of the soft mode determined from the self-consistency condition.

We note in conclusion that the obtained equations allow us to analyze also nonlinear problems, and can also be generalized to include the case when pinning plays an important role.

The authors are most grateful to S.A. Brazovskiĭ for a discussion of the work and for helpful remarks.

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Translated by J. G. Adashko