Nonlocal effects during current flow in a quasi-1-*D* conductor with a charge-density wave

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It is shown theoretically that when a current I flows through part of a conductor with a charge-density wave an electric field and a voltage arise in a region in which no current is flowing. If I exceeds the threshold pinning current, domain walls—phase solitons—may be created at the boundaries of the region in which the current flows. These effects occur because a moving charge-density wave is deformed in the region with $I = 0$.

1. Considerable interest has recently been attracted to effects accompanying a current flow in a quasi-1-*D* conductor with a charge-density wave. In particular, it has been found that measurable quantities (such as the voltage and the current) depend on the positions of the measuring contacts and those which conduct the current.[1–3] Latyshev *et al.* (see the preceding paper) have recently observed the appearance of an electric field E outside the current-flow region (Fig. 1). They used filamentary samples with diameters up to several microns. A current I was passed through contacts 2 and 3, while the voltages V_{ik} (i, k specify the contacts) were measured between current contacts 2, 3 and outer contacts 1, 4. It was found that voltages V_{12} and V_{34} arise at a

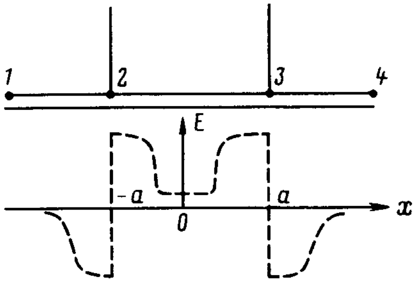

FIG. 1.

current above the threshold current for the pinning of a charge-density wave; these voltages are comparable in magnitude to V_{23} and opposite in sign.

Let us analyze the current flow in a conductor with a charge-density wave in the geometry of Fig. 1. We will show that the particular response of the phase of the charge-density wave to the electric field leads to results in qualitative agreement with the experimental data. Furthermore, we will predict some effects which can apparently be observed in future experiments (for example, the creation of phase solitons at contacts 2 and 3).

Since we will ignore the suppression of the energy gap Δ, it is sufficient to use only the equations for the phase (χ) of the charge-density wave and the current I. We use the results of the microscopic theory of Ref. 4. At low temperatures ($T \ll \Delta$), and at a Fermi-surface curvature small in comparison with Δ, the equations can be written

$$-\frac{D}{2}\chi'' + \lambda_1 \dot{\chi} = El \tag{1}$$

$$I = I_q + \sigma_N \dot{\chi}/l, \quad \left(\frac{\partial}{\partial t} - D_1 \frac{\partial^2}{\partial x^2}\right) I_q = \lambda \sigma_N \frac{\partial E}{\partial t}, \tag{2}$$

where $D = vl$ is the diffusion coefficient, l is the mean free path, $\lambda_1 = C_1 \dfrac{\Delta}{T}\exp(-\Delta/T)$ is the coefficient of friction of the charge-density wave,[1)] I_q is the quasiparticle current, $D_1 = C_2\left(\dfrac{T}{\Delta}\right)^{3/2} D$, $\lambda = C_3\left(\dfrac{T}{\Delta}\right)\exp\left(-\dfrac{T}{\Delta}\right)$, σ_N is the conductivity in the normal state (i.e., with $\Delta = 0$), and the coefficients C_1, C_2, and C_3 are numbers on the order of unity. Equation (2) was derived from the quasiparticle diffusion equation (within a numerical factor) by replacing the energy-dependent diffusion coefficient $D(\epsilon)$ by an energy-independent coefficient D_1. In the limit of a steady-state field E, or a field which is uniform over space, this equation is exact in the models of Ref. 3.

It should be kept in mind that Eqs. (1) and (2) contain the x components of the current density and the field averaged over the cross section of the conductor, so that we have $I(x) \equiv \langle j_x \rangle = I\theta(a - |x|)$, where $2a$ is the distance between contacts 2 and 3.

2. Let us assume that an alternating current $I(t) = I_\omega \exp(-i\omega t)$ is flowing through the conductor. The spatial profile of the phase of the charge-density wave can then be found easily from (1)–(2) and the continuity of the phase and its derivative at points 2 and 3:

$$\chi_\omega = -(I_\omega l/i\omega\sigma_N)\{[1 - e^{-\kappa a}\cosh(\kappa x)]\,\theta(a - |x|)$$
$$+ \frac{1}{2}(1 - e^{-2\kappa a})e^{(a - |x|)\kappa}\,\theta(|x| - a)\},$$

where $\kappa^2 = -i\omega D_2$ and $D_2 = \lambda D/2 + D_1 \approx D_1$. Using (1) and (2), we can find the distribution of the field $E_\omega(x)$ and calculate the voltages V_{23}, $V_{12} = V_{34}$, and V_{14}. We find

$$V_{23}=V_0\,[(D/2D_1)f+\lambda_1(1-f)],\quad V_{12}=\frac{V_0}{2}f[\lambda_1-D/2D_1],\quad V_{14}=\lambda_1 V_0\ ,$$

where

$$V_0=2aI_\omega\ /\sigma_N\ ,\qquad f=(1-e^{-2\kappa a})/(\kappa a).$$

At sufficiently low frequencies we have $f\approx 2$. Since λ_1 is small, we find $V_{23}=V_0(D/D_1)\approx -V_{ex}\equiv-(V_{12}+V_{34})$ in this case. The total voltage V_{14} is thus small in comparison with both V_{23} and the external voltage V_{ex}, which cancel each other out to a large extent. We might also note that the voltage V_{23} is quite different from the voltage which would prevail during the uniform slippage of a charge-density wave. The latter voltage is the same as V_{14}. The dashed curve in Fig. 1 is a sketch of the field distribution $E_\omega(x)$. The depth to which the field E_ω penetrates into the regions (1,2) and (3,4), $L_E(\omega)\approx\sqrt{D_1/\omega}$, decreases with increasing ω.

The results of this section apply at $V_0\ll\Delta$ and when pinning is ignored. In the high-temperature limit, $T\gg\Delta$, these effects are small, on the order of the parameter Δ/T.

3. When the current through contacts 2 and 3 is a large square pulse, the field distribution is initially as shown in Fig. 1. There is an increase in the distance L_E to which the field E penetrates and over which the velocity of the charge-density wave is nonzero ($\chi\neq 0$); the voltage V_{ex} also increases. The subsequent evolution of E and χ depends on the current I.

At a large current I or a large distance a such that the condition

$$I>I_c\sim\sigma_N\frac{\Delta}{a}\left(\frac{T}{\Delta}\right)^{3/2}\tag{3}$$

holds, an increase in the phase gradient $\partial\chi/\partial x$ leads to a suppression of the gap Δ at points 2 and 3. This occurs when $-V_{ex}\approx V_{23}\sim\Delta$, when the electric field penetrates a distance $(\sigma_N/I)\Delta\,(T/\Delta)^{3/2}$. In this case, phase-slippage centers arise in the system,[5] accompanied by nonlinear oscillations of Δ and χ which cannot be described by Eqs. (1) and (2).

If, on the other hand, the current I or the distance between the current contacts is small, so that condition (3) does not hold (but the condition $I\gg I_T$ does, where I_T is the threshold current for pinning), then perturbations of E and χ propagating away from contacts 2 and 3 begin to suppress each other after a time on the order of a^2/D_1, and the field D begins to decrease in magnitude, in proportion to $1/\sqrt{t}$, propagating ever further from contacts 2 and 3. The voltages reach a steady-state value $-V_{ex}\approx V_{23}=V_0(\Delta/T)^{3/2}$, $V_{14}=\lambda_1 V_0$. The decrease in the field E and its propagation in regions (1, 2) and (3, 4) continue until the distribution becomes fixed by pinning.

4. To take the pinning of the charge-density wave into account we adopt a very simple model, which can be defended rigorously in the case of a commensurate charge-density wave. Specifically, we add to the left side of Eq. (1) a term $E_0 l\sin n\chi$, where $E_0 l\sim\Delta\,(\Delta/\epsilon_F)^{n-2}$, and n is the order of the commensurability. Clearly, in the case of an alternating current we can ignore the pinning if $\omega\lambda_1\gg E_0 l$. What are the conse-

quences of incorporating pinning during the flow of a direct current I? If the current is weak, the phase is perturbed in region (2, 3), and it falls to zero as $x \rightarrow \pm \infty$ over a characteristic distance L ($L = \sqrt{D_1/\lambda E_0 l}$ is the size of the phase soliton; the increase in the length with decreasing temperature is due to screening[4]). In this case we find for $V_{23}^{23} = (V_0/\lambda)[1 - (1/2)f(a/L)]$. Under the condition $a \gg L$ we thus have $V_{23} = V_0/\lambda$, while at $a \ll L$ we have $V_{23} = (V_0/\lambda)(a/L)$; i.e., there is a change in the temperature dependence of the conductivity I/V_{23}. The voltage V_{14} is again related to the current, as in the homogeneous case, $V_{14} \approx V_0/\lambda$. With increasing I, the static deformation of the charge-density wave increases, and at a certain threshold current I_T there is no static solution for χ, which is unperturbed at infinity. At $I > I_T$ there is a static solution in the form of domain walls: a chain of phase solitons whose period depends on the difference $I - I_T$. At currents slightly above the critical current, the chain forms through the sequential creation of solitons at points 2 and 3 and their propagation away from these points. It is easy to derive an expression for $I_T = (2L/a)\lambda\sigma_N E_0$ for

$$a \ll L \text{ и } I_T = \lambda\sigma_N E_0 [1 + 24K^4(1/\sqrt{2})\,(L/a)^4]$$

for $a \gg L$. At high currents, $I \gg I_T$, the oscillations of the phase χ over space become progressively weaker, and the spatial dependence of χ approaches that of a static solution of the diffusion equation, $\chi' = \text{const}$.

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1)The coefficient λ_1 is small because the friction experienced by the charge-density wave results from the presence of excitations above the gap. If other possible friction mechanisms—in particular, states in the gap—were taken into account, λ_1 would increase.

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