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Spin-Dependent Tunneling in Junctions Containing Metals with Charge Density Waves in a Magnetic Field

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Abstract—The dependences of the differential tunneling conductance G on the voltage V across a junction in an external magnetic field H are calculated for two types of junctions involving normal or superconducting metals with charge density waves (CDWs). Junctions of the first type are asymmetric CDW metal (CDWM)–insulator–ferromagnet junctions. The results of calculations for these junctions demonstrate that there occurs splitting between the components of the conductance $G(V)$ corresponding to the tunneling of electrons with spins aligned with the magnetic field H and opposite to it, as is the case with junctions containing a superconducting electrode instead of the CDWM electrode. Junctions of the second type are junctions between two normal or superconducting CDWM electrodes. For junctions with at least one normal CDWM electrode and $H \neq 0$, the conductance $G(V)$ also exhibits spin splitting. The form of the conductance $G(V)$ for tunnel junctions of both types depends on the phase of the order parameter of the charge density waves.

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1. INTRODUCTION

Since electrons have both charge and spin, there exists a possibility of controlling spin-polarized currents and, as a consequence, spin-encoded information flows. This problem has been already solved, and the branch of electronics dealing with similar phenomena is referred to as spin electronics or spintronics [1]. Among the various spintronics applications, the method proposed by Tedrow and Meservey stands out [2]: it was found that the spin polarization P of electrons is conserved in their tunneling through a ferromagnet–insulator–superconductor (FM–I–S) junction. The density of states of quasiparticles in Bardeen–Cooper–Schrieffer superconductors has a gapped character. This allows one to reveal and measure quantitatively the spin polarization of electrons if the tunnel junction is in an external magnetic field H (it should be noted that the relation between the polarization P thus determined and the spin polarization of electrons in the ferromagnet in the general case is ambiguous [3]). An important feature is that, in the Tedrow–Meservey configuration, the peaks of the tunneling conductance $G(V) = dJ/dV$ (where V is the bias voltages across the junction and J is the quasiparticle tunneling current) are split into spin-polarized components $G^{\pm}$ only in the case where at least one electrode is a normal metal. (In what follows, the upper and lower signs will correspond to the directions of the spin polarization for the "minority" and "majority" of tunneling electrons with respect to the direction of the magnetic field H. The mnemonic rule is that the energy levels of the minority of electrons lie higher than the energy levels of the majority of electrons.) When both electrodes are superconductors, no splitting occurs in the absence of spin flip, because the energy levels of quasiparticles at the edges of the superconducting energy gaps of both electrodes are shifted in the same direction by the same magnitude for each orientation of the spin.

Nonetheless, the effective Tedrow–Meservey method as applied to the determination of the spin polarization P encounters some difficulties. First and foremost, it should be noted that the spin–orbit interaction brings about a shift of the $G^{\pm}(V)$ peaks with a larger splitting (along the V axis) toward the peaks with a smaller splitting (and the opposite spin orientation) up to their complete merging [4]. It is known that the intensity of spin–orbit scattering $\hbar/\tau_{so}$ is proportional to Z^4, where $\hbar$ is Planck's constant, τ_{so} is the time of spin–orbit scattering, and Z is the atomic number of the electrode material [5]. As a result, the dimensional parameter b, which determines the deterioration of the spin splitting, is equal to $\hbar/3\tau_{so}\Delta \propto Z^4/\Delta$ [4], where Δ is the amplitude of the superconducting gap. Therefore, in order to observe the Tedrow–Meservey effect, it is necessary to use higher temperature superconducting materials or superconductors consisting of light elements. However, these requirements can appear contradictory.

Moreover, there exists another phenomenon that can render the spin spitting of the conductance $G(V)$ unresolved. The case in point is the main effect inherent in conventional Bardeen–Cooper–Schrieffer superconductors in a magnetic field, i.e., the orbital diamagnetic Meissner effect. Note that the paramagnetic Zeeman effect of the magnetic field (discussed in this paper) in conventional superconductors is a predecessor of the destruction of the Cooper pairs, which occurs at a magnetic field strength H higher than the so-called paramagnetic limit H_p [6]. For the vast majority of bulk superconductors, the paramagnetic limit H_p exceeds the upper critical magnetic field H_{c2} determined by the diamagnetic effect, although there are known exceptions, such as the α-(*BEDT–TTF*)$_2$KHg(SCN)$_4$ [7] and κ-(*BEDT–TTF*)$_2$Cu(NCS)$_2$ [8] compounds (the possible origins of this behavior were discussed in [9, 10]). Therefore, in order to suppress the Meissner diamagnetism, it is necessary to use superconducting electrodes in the form of thin films [2].

At the same time, there exists a class of solid-state objects, namely, insulators and metals with charge density waves (CDWs), which possess paramagnetic properties similar in many respects to the properties of superconductors [11] but exhibit gapped features in the energy spectrum due not to Cooper pairing but to electron–hole pairing. On this basis, we propose to use normal charge-density-wave metals (CDWMs) and charge-density-wave superconductors (CDWSs) as components for the tunnel detectors of ferromagnetic polarization, because these materials offer undeniable advantages over conventional superconductors.

In the first part of this paper, we analyze the tunneling currents in FM–I–CDWM junctions. In the second part of the paper, we consider the CDWM (CDWS)–I–CDWM (CDWS) symmetric junctions, for which the current–voltage characteristics differ significantly from those obtained for S–I–S junctions. In particular, the spin splitting of the conductance $G(V)$ in the former junctions [12] manifests itself due to the existence of nondielectrized (gapless) regions on the Fermi surface of the metals with charge density waves even at temperatures below the temperature T_d of the structural phase transition [9]. This stems from the fact that the current–voltage characteristics of the CDWM–I–CDWM junctions exhibit a combination of specific features of both asymmetric and symmetric tunnel junctions. Another characteristic feature of the junctions involving metals with charge density waves is the dependence of the tunneling conductance on the phases φ_i of the order parameter of charge density waves $\tilde{\Sigma}_i \equiv \Sigma_i e^{i\varphi_i}$, which is distinguished from the dependence of the tunneling conductance $G(V)$ for FM–I–S or S–I–S junctions. The difference lies in the fact that the so-called interband "normal" temperature Green's function $\mathcal{G}_c$ for a Peierls or excitonic insulator with charge density waves plays the same role as the anomalous Gor'kov Green's function $\mathcal{F}$ for superconductors [13]. Consequently, the quasiparticle current $J(V)$ is a specific analog of the Josephson tunneling current between superconductors.

In conclusion, we discuss particular materials for CDWM electrodes and the possibility of observing the predicted effects.

2. FORMULATION OF THE PROBLEM

2.1. Model of a Metal with Charge Density Waves

As the starting point, we use the Bilbro–McMillan Hamiltonian for a partially dielectrized superconducting metal with charge density waves [14]. According to this model (see also review [9]), the Fermi surface of this metal is separated into three sections, namely, two congruent sections ($i = 1, 2$) in which the spectrum of quasiparticles is degenerate (d) and one section with a nondegenerate (n) spectrum ($i = 3$). In the former case, the branches of the bare energy in the energy spectrum of quasiparticles $\xi_{1,2}(\mathbf{p})$ (which are reckoned from a common Fermi level) are related by the expression

$$\xi_1(\mathbf{p}) = -\xi_2(\mathbf{p}+\mathbf{Q}), \tag{1}$$

where $\mathbf{Q}$ is the CDW vector. Owing to the interaction between quasiparticles from different degenerate sections of the Fermi surface, there arises an interparticle correlation in the zero-sound channel. As a consequence, the dielectric CDW gap Σ appears in both these sections. If the above interaction is associated primarily with the Coulomb mechanism [15] and the branches $\xi_{1,2}(\mathbf{p})$ describe electrons and holes, respectively, the CDW dielectrization corresponds to the formation of an excitonic insulator (completely or partially dielectrized). A similar situation occurs in the case of an anisotropic quasi-one-dimensional spectrum with planar sections of the Fermi surface, for which the branches $\xi_{1,2}(\mathbf{p})$ can also be separated and the mutual attraction of quasiparticles is provided by the electron–phonon interaction. As a result, there arises a Peierls insulating state [16]. In both cases, the pairing involves quasiparticles with oppositely directed spins (singlet pairing). The third section of the Fermi surface remains undistorted and is described by the nondegenerate branch $\xi_3(\mathbf{p})$ of the spectrum. The separation of the Fermi surface into dielectrized and nondielectrized sections is described by the dielectrization parameter

$$\mu = N_{d0}(0)/N_0(0), \tag{2}$$

where $N_0(0) = N_{n0}(0) + N_{d0}(0)$ is the total initial density of states at the Fermi surface (above the temperature T_d; and $N_{d0}(0)$ and $N_{n0}(0)$ are the components of the total initial density of states, i.e., the densities of states at the degenerate and nondegenerate sections of the Fermi surface, respectively.

At temperatures below the critical temperature of the superconducting transition $T_c < T_d$, there appears a superconducting order parameter $\tilde{\Delta} \equiv \Delta e^{i\varphi}$ (which is the same over the entire Fermi surface). This means that, at temperatures below the critical temperature T_c, the degenerate sections of the Fermi surface are distorted as a result of the pairing of both types, which gives rise to a combined energy gap $D = \sqrt{\Sigma^2 + \Delta^2}$, whereas the electron spectrum in the nondegenerate section has a gap Δ [9, 14]. It should be emphasized that the quasiparticle current J under investigation does not depend on the phase ϕ of the superconducting order parameter, which determines the Josephson current flowing through a weak link between the superconductors.

The charge density wave can be either commensurate or incommensurate with the parameters of the crystal lattice. In the excitonic insulator model, the lattice distortion caused by the Coulomb interaction below the temperature T_d is commensurate. Moreover, the phase of the order parameter in an excitonic insulator always takes on the values $\varphi = 0$ or π [17]; hence, the order parameter $\tilde{\Sigma}$ is either a positive real quantity or a negative real quantity [9, 15]. However, incommensurate charge density waves in Peierls insulators can exhibit quite different dynamics, even though these waves in specific current measurements can be fixed with an arbitrarily frozen phase φ [16].

At $H \neq 0$, the quasiparticle currents flowing through junctions containing metals with charge density waves are functions of the following time Green's functions $G_{ij}(\omega)$ describing electrodes:

$$G_{11} = G_{22} \equiv G_d, \quad (3)$$

$$G_{12} = G_{21} \equiv G_c, \quad (4)$$

$$G_{33} \equiv G_n, \quad (5)$$

where the subscripts $i, j = 1, 2, 3$ number the Fermi surface sections (see above). The Green's functions for all other combinations of the subscripts i and j are equal to zero: $G_{ij} = 0$. The function G_c describes the electron–hole pairing.

As is known, the dependence of the current J on the voltage V is exponential at high voltages and ohmic at low voltages. We will restrict our subsequent analysis to the specific case of the ohmic regime, because the superconducting and dielectric gaps of the normal and superconducting CDW metals under investigation fall in the range 0.2–30.0 meV. However, the deviations from the Ohm law, which indicate a crossover to the tunneling regime, appear in the case where the change in the energy of an electron involved in tunneling eV becomes comparable to the conduction band width $W \geq$ 1 eV for a particular electrode (hereafter, $e > 0$ will be the elementary charge).

In analyzing the tunneling currents between metals with charge density waves, we use a standard tunneling Hamiltonian. Since the resistance for the quasiparticle current between the metals with charge density waves is ohmic, we can introduce a single parameter R, i.e., the resistance of the junction in the normal state. The resistance R is inversely proportional to the square of the tunneling matrix element averaged over the Fermi surfaces of both electrodes [9].

In an external magnetic field H (the spatial z axis is chosen to be aligned with the magnetic field H), the paramagnetic effect in the degenerate and nondegenerate quasiparticle states manifests itself in quite different ways. It should be noted once again that, in analyzing the spin splitting of the $G(V)$ peaks, we ignore the dynamic response of the metal with charge density waves. Of course, this does not imply that the temperature T_d is independent of the magnetic field H if we go beyond the scope of the approximation used. However, the analysis of the experimental data demonstrates that various characteristic nonlinear and oscillatory diamagnetic phenomena are observed [11] at a magnetic field strength H higher than the paramagnetic limit H_p^{CDWM} for normal and superconducting CDW metals [10], whereas the Zeeman splitting under investigation becomes noticeable at lower magnetic field strengths.

Now, we consider the paramagnetic properties of a metal with charge density waves in a nonzero magnetic field ($H < H_p^{\mathrm{CDWM}}$). The schematic diagram of the energy levels of quasiparticles is depicted in Fig. 1. The energy of electrons with the spin projection $s_z = +1/2$ onto the magnetic field H increases by $\mu_B^* H$, whereas the energy of electrons with the opposite spin direction $s_z = -1/2$ decreases by the same value. Here, $\mu_B^* = e\hbar/(2m^*c)$ is the effective Bohr magneton, c is the velocity of light, and m^* is the effective mass of charge carriers. In what follows, quasiparticles with the corresponding spin directions will be denoted by "+" and "–," respectively.

Quasiparticles that belong to the nondegenerate section of the Fermi surface in which the Fermi level at $T = 0$ separates the occupied and empty states behave in a conventional manner characteristic of quasiparticles in a normal metal. An increase in the magnetic field leads to an increase in the energy of the states from the "+" subband that, at $H = 0$, coincide in the energy with their "–" analogs (both subbands are equally filled). As a result, the quasiparticles pass from the "+" subband into the "–" subband, so the number of particles in the "+" subband decreases with an increase in the magnetic field H, whereas the number of occupied states in the "–" subband correspondingly increases. This spin polarization induced by the external magnetic field leads to a change in the chemical potential $\tilde{\mu}$ by a value of the order of $(\mu_B^* H/E_F)^2$, where E_F is the Fermi energy. Since we examine effects for which the quantity $\mu_B^* H$ is at least smaller than the order parameters Δ and Σ, the inequality $(\mu_B^* H/E_F)^2 \ll 1$ is satisfied and the change in the chemical potential $\tilde{\mu}$ can be disregarded altogether. It should be remembered that this is not the case with band Stoner ferromagnets, in which the chemical potential $\tilde{\mu}$ varies significantly in the ordered state [18].

The paramagnetic splitting of the quasiparticle states on the Fermi surface sections with a gap (the nondegenerate section at temperatures below the temperature T_c and the degenerate section at temperatures below the temperature T_d) can be analyzed in the same manner as for superconductors [2]. The similarity of the phenomena in these cases is explained by the fact that both electron–hole and Cooper pairs are spin-singlet objects and, hence, are destroyed by the Zeeman splitting [19]. As a consequence, the energies of quasiparticles from the "+" and "–" subbands of the corresponding Fermi surface sections with a gap are shifted in opposite directions at $H \neq 0$. Within this approach, we can perform only a qualitative analysis of the phenomenon, because all processes that are associated with the spin flip and result in a smearing of the ideal splitting are disregarded.

In the subsequent analysis, the tunneling in the magnetic field $H \neq 0$ will be described in terms of the aforementioned time Green's functions G_d, G_c, and G_n. As compared to the situation considered in our previous paper [9], the number of these Green's functions is doubled and six new functions dependent on the magnetic field H will be denoted with the use an additional subscript $s = \pm$, i.e., G_{ds}, G_{cs}, and G_{ns}. These functions each depend on one of the variables $\omega_{\mp} = \omega \mp \mu_B^* H$, whose sign is chosen to be opposite to the sign of the subscript s.

2.2. Current–Voltage Characteristics

2.2.1. Asymmetric junctions. Let us now consider the current–voltage characteristics for asymmetric junctions. In this case, the bias voltage V between the ferromagnet and normal or superconducting CDW metal is reckoned from the CDW metal: $V \equiv V_{FM} - V_{CDWM}$, where V_{FM} and V_{CDWM} are the potentials of the ferromagnet and the CDW metal, respectively.

It is assumed that, at a magnetic field strength H high enough to induce the experimentally observed splitting of the peaks in the density of states, all domains in the ferromagnet are completely aligned with the magnetic field [2]. Moreover, we assume that the initial polarization P of a quasiparticle in the bulk of the ferromagnet is conserved during tunneling. This means that the influence of the FM–I interface on the tunneling current is completely ignored. Generally speaking, this strong assumption is made for simplicity and does not affect the generality of the results obtained. However, we recognize that the effects associated with the interface and its disordering can turn out to be important for specific applications [20].

The current $J(V)$ between the ferromagnet and the CDW metal is calculated with the use of the Bardeen–Cooper–Schrieffer dependence for the dielectric gap

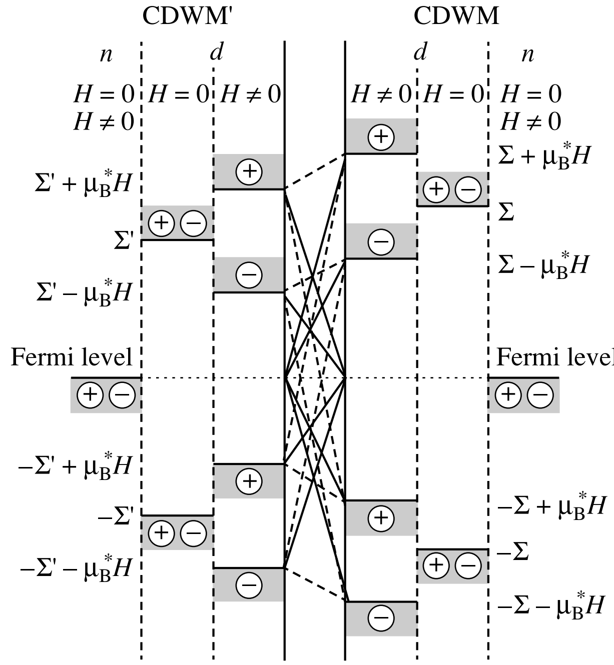

Fig. 1. Schematic diagram of the energy levels of quasiparticles in electrodes of the tunnel junction between normal dissimilar partially dielectrized metals with charge density waves at a zero bias voltage $V = 0$ in magnetic fields $H = 0$ and $H \neq 0$. The spectra of quasiparticles for incongruent (n) and congruent (d) sections of the Fermi surface are depicted separately on both sides of the insulating barrier. Designations: Σ' and Σ are the CDW gaps in the left and right electrodes, respectively, and μ_B^* is the effective Bohr magneton. Hatched regions of the congruent sections of the Fermi surface indicate the positions of the spin subbands with respect to the edges of the corresponding CDW gap and the levels of these edges on the energy scale. The signs "+" and "–" correspond to the spin orientations in the subbands aligned with the direction of the magnetic field and opposite to it, respectively. The Fermi level in the incongruent section is assumed to be independent of the magnetic field H (see the text). The scheme of the possible tunneling transitions of quasiparticles without a spin flip is shown. Solid and dashed lines represent the quasiparticle current components dependent on and independent of the magnetic field H, respectively.

$\Sigma(T)$ and the method developed for conventional superconductors [21]. In our case, the total current $J(V)$ can be represented as the sum of different terms $J_{i\pm}$ that have the same form

$$J_{i\pm} \propto \mathrm{Re} \int_{-\infty}^{\infty} d\omega' \int_{-\infty}^{\infty} d\omega \frac{\mathrm{Im}\, G_{i\pm}^{\mathrm{CDWM}}(\omega'_{\mp}) G_{\pm}^{\mathrm{FM}}(\omega_{\mp})}{\omega' - \omega + eV + i0} \quad (6)$$

and correspond to different combinations of spin-dependent time Green's functions. The form of the Green's functions $G_{\pm}^{\mathrm{FM}}(\omega)$ is well known. However, instead of one pair of the functions $G_{\pm}^{\mathrm{BCS}}(\omega)$ for the Bardeen–Cooper–Schrieffer superconductor, there appear six functions $G_{i\pm}^{\text{CDWM}}(\omega)$. These functions can be obtained from the following temperature Green's functions for the CDW metal:

$$\mathcal{G}_n^{\pm}(\mathbf{p}, \omega_n) = \frac{i\omega_n \mp \mu_B^* H + \xi_3(\mathbf{p})}{(i\omega_n \mp \mu_B^* H)^2 - \xi_3^2(\mathbf{p}) - \Delta^2}, \quad (7)$$

$$\mathcal{G}_d^{\pm}(\mathbf{p}, \omega_n) = \frac{i\omega_n \mp \mu_B^* H + \xi_1(\mathbf{p})}{(i\omega_n \mp \mu_B^* H)^2 - \xi_1^2(\mathbf{p}) - \Delta^2}, \quad (8)$$

$$\mathcal{G}_c^{\pm}(\mathbf{p}, \omega_n) = \frac{\tilde{\Sigma}}{(i\omega_n \mp \mu_B^* H)^2 - \xi_1^2(\mathbf{p}) - \Delta^2}. \quad (9)$$

Here, $D^2 \equiv \Delta^2 + \Sigma^2$, $\omega_n = (2n + 1)\pi T$ ($n = 0, \pm 1, \pm 2, \ldots$), T is the temperature, and $k_B = 1$ is the Boltzmann constant. Using the procedure previously developed for similar problems, we derive the following relationship for the current:

$$J(V) = \sum_{f = n, d, c;\, s = -, +} J_{fs}(V), \quad (10)$$

where

$$J_{n\pm} = \frac{(1-\mu)(1 \mp P)}{4eR} \int_{-\infty}^{\infty} d\omega K(\omega, V, T)|\omega_{\mp}| f(\omega_{\mp}, \Delta), \quad (11)$$

$$J_{d\pm} = \frac{\mu(1 \mp P)}{4eR} \int_{-\infty}^{\infty} d\omega K(\omega, V, T)|\omega_{\mp}| f(\omega_{\mp}, D), \quad (12)$$

$$J_{c\pm} = \frac{\mu(1 \mp P)\Sigma\cos\varphi}{4eR} \times \int_{-\infty}^{\infty} d\omega K(\omega, V, T)\,\text{sgn}(\omega_{\mp}) f(\omega_{\mp}, D), \quad (13)$$

$$K(\omega, V, T) = \tanh\frac{\omega}{2T} - \tanh\frac{\omega - eV}{2T}, \quad (14)$$

$$f(\omega, X) = \frac{\theta(|\omega| - X)}{\sqrt{\omega^2 - X^2}}, \quad (15)$$

$X = \Delta$ or D, R is the resistance of the junction in the "doubly normal" state (i.e., at temperatures higher than the temperatures T_c and T_d), and $\theta(x)$ is the Heaviside function. Note once again that the signs in the variables $\omega_{\mp}$ involved in the integrands for the current components are opposite to those for the subscripts s.

It should also be taken into account that the current components depend on the phase φ of the dielectric order parameter $\tilde{\Sigma} = \Sigma e^{i\varphi}$, whereas the thermodynamic properties of CDW superconductors are degenerate with respect to the phase φ [22]. Moreover, we assume that the total contribution of quasiparticles from each section of the Fermi surface to the total tunneling current is proportional to the density of states for the given section. This means that directional tunneling, which, in principle, is possible [23], is absent.

The main difference between the problem under consideration and its analogs for "pure" superconductivity is the appearance of the terms $J_{c\pm}$. These terms result from the interband Green's function $\mathcal{G}_c$ [see expression (9)], which describes the pairing of electrons with holes in different congruent sections of the Fermi surface [9, 15]. The form of these terms differs from that of the other terms relating to the conventional normal Green's functions $\mathcal{G}_d$ and $\mathcal{G}_n$ [see expressions (7) and (8)]. The Green's function $\mathcal{G}_c$, to a considerable extent, is an analog of the Gor'kov–Green's function $\mathcal{F}$, which, however, determines the Josephson current rather than the quasiparticle tunneling current. The appearance of terms (13) leads to a strong asymmetry of the current–voltage characteristics for asymmetric tunnel junctions involving CDW metals [9] in contrast to the symmetric current–voltage characteristics for similar asymmetric junctions based on conventional superconductors. In an incommensurate CDW metal, the order parameter phase φ is arbitrary [16]. For a commensurate charge density wave, we have the order parameter phase $\varphi = 0$ or π. Therefore, the corresponding equations describe, in particular, the tunneling between excitonic insulators [15]. In principle, there can arise an opposite situation where the tunneling contact area is large enough to cover several sections with different order parameter phases. In this case, terms (13) turn out to be averaged to some extent and the asymmetry of the current–voltage characteristic is substantially less pronounced.

In specific experiments, the averaging of the phase can occur or be absent. If the averaging takes place, the observed properties of the tunnel junction can change radically. It should be noted that, in CDWM–I–CDWM junctions, the "coherent" current component proportional to $\cos(\varphi_1 - \varphi_2)$ [12, 24] should also be averaged, as is the case with the other components dependent on the phase. However, if the phase φ in the CDWM electrode is constant in the spatial region responsible for the tunneling current in the FM–I–CDWM junction, the current component J_c should have the initial unaveraged form (13). It should be emphasized that the predicted spin splitting is retained irrespective of the phase averaging even in the case where the current is governed by a set of different tunneling matrix elements [24].

2.2.2. Symmetric junctions. For a junction containing dissimilar CDW metals (the CDWM'–I–CDWM junction) in the magnetic field H, the structure of the features of the conductance $G(V)$ is more complex when both CDWM electrodes are normal metals rather than superconductors. This unexpected situation occurs because there is no spin splitting if the Fermi surfaces of both electrodes are completely dielectrized, regardless of the nature of the collective energy gap (gaps).

In our case, the spin splitting can be observed only in the case where the dielectrization of the Fermi surface is partial and $\Delta = 0$. Without going into detail on the form of the contributions from different sections of the Fermi surface to the total current, we will restrict our analysis to the simple case of symmetric junctions with identical CDWM electrodes. Within the Bilbro–McMillan model [14], the symmetry implies that the equalities $\Sigma = \Sigma'$ and $\mu = \mu'$ are satisfied, even though the phases φ and φ' for formally identical states with the same energy can differ from each other. This phenomenon is the spontaneous breaking of the symmetry with respect to the phases of the order parameters of charge density waves [9]. This breaking is similar to the known phenomena in degenerate magnetic and other many-particle systems [25–27].

The current–voltage characteristics obtained for the CDWM–I–CDWM symmetric junctions are considerably simpler than those observed in the general case. Indeed, the positions of the features dependent on the amplitudes of the gaps in the CDWM' and CDWM electrodes coincide and a number of preintegral factors become equal to each other. Therefore, the entire set of components of the quasiparticle tunneling current through the CDWM–I–CDWM symmetric junction with due regard for the possible violation of the phase equality can be represented in the following form:

$$J_{dd\pm} = \frac{\mu^2}{4eR}\int_{-\infty}^{\infty} d\omega K(\omega, V, T)|\omega_{\mp}|f(\omega_{\mp}, D)|\omega_{\mp} - eV| \times f(\omega_{\mp} - eV, D), \quad (16)$$

$$J_{cc\pm} = \frac{(\mu\Sigma)^2\cos\varphi'\cos\varphi}{4eR}\int_{-\infty}^{\infty} d\omega K(\omega, V, T)\,\mathrm{sgn}(\omega_{\mp}) \times f(\omega_{\mp}, D)\,\mathrm{sgn}(\omega_{\mp} - eV)f(\omega_{\mp} - eV, D), \quad (17)$$

$$J_{nn\pm} = \frac{(1-\mu)^2}{4eR}\int_{-\infty}^{\infty} d\omega K(\omega, V, T)|\omega_{\mp}|f(\omega_{\mp}, \Delta) \times |\omega_{\mp} - eV|f(\omega_{\mp} - eV, \Delta), \quad (18)$$

$$J_{dn\pm} = \frac{\mu(1-\mu)}{4eR}\int_{-\infty}^{\infty} d\omega K(\omega, V, T)|\omega_{\mp}|f(\omega_{\mp}, D) \times |\omega_{\mp} - eV|f(\omega_{\mp} - eV, \Delta), \quad (19)$$

$$J_{nd\pm} = \frac{\mu(1-\mu)}{4eR}\int_{-\infty}^{\infty} d\omega K(\omega, V, T)|\omega_{\mp}|f(\omega_{\mp}, \Delta) \times |\omega_{\mp} - eV|f(\omega_{\mp} - eV, D), \quad (20)$$

$$J_{cn\pm} = \frac{\mu(1-\mu)\Sigma\cos\varphi'}{4eR}\int_{-\infty}^{\infty} d\omega K(\omega, V, T)\,\mathrm{sgn}(\omega_{\mp}) \times f(\omega_{\mp}, \Sigma)|\omega_{\mp} - eV|f(\omega_{\mp} - eV, \Delta), \quad (21)$$

$$J_{nc\pm} = \frac{\mu(1-\mu)\Sigma\cos\varphi}{4eR}\int_{-\infty}^{\infty} d\omega K(\omega, V, T)|\omega_{\mp}| \times f(\omega_{\mp}, \Delta)\,\mathrm{sgn}(\omega_{\mp} - eV)f(\omega_{\mp} - eV, \Sigma). \quad (22)$$

At $T = 0$, all the current components can be expressed through elliptic integrals. At $T \neq 0$, it is necessary to perform numerical calculations.

3. NUMERICAL CALCULATIONS

In subsequent calculations, all the results obtained will be represented in the form of dependences of the dimensionless conductance RdJ/dV of the FM–I–CDWM or CDWM–I–CDWM junctions on the dimensionless bias voltage eV/Σ_0, where $\Sigma_0 \equiv \Sigma(T = 0)$. The other dimensionless parameters of the problem are the reduced external magnetic field $h = \mu_B^* H/\Sigma_0$, the reduced temperature $t = T/\Sigma_0$, and the polarization P.

3.1. Asymmetric Junctions

The splitting of the features of the conductance $G(V)$ in an external magnetic field for the FM–I–CDW normal metal junction at $\varphi = 0$ is illustrated in Fig. 2. It is clearly seen from this figure that, unlike the current–voltage characteristics for tunnel junctions with conventional superconductors, the dependence of the conductance $G(V)$ is highly asymmetric. Mathematically, this results from an almost complete mutual compensation of the features corresponding to the terms $G_d(V)$ and $G_c(V)$ at bias voltages of one sign and their enhancement at bias voltages of the other sign. In the magnetic field H, the peaks of the density of states are split, as is the case with superconductors [2]. The spin splitting depends substantially on the polarization P but is observed only in one branch of the current–voltage characteristic (at $V > 0$ in the case of $\varphi = 0$). The other branch contains only residual fragments of the features associated with the energy gap.

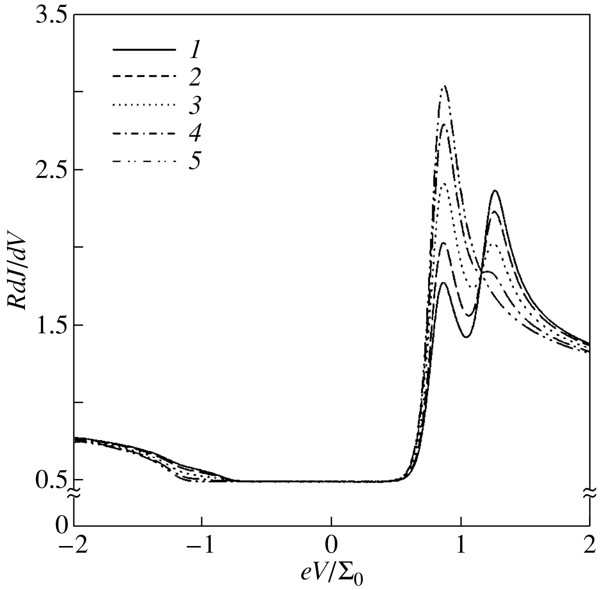

Fig. 2. Dependences of the dimensionless conductance $RdJ/dV = RG(V)$ on the dimensionless bias voltage eV/Σ_0 in a tunnel junction between the ferromagnet and the normal CDW metal in the magnetic field H for different polarizations of the ferromagnet P = (*1*) 0, (*2*) 0.2, (*3*) 0.5, (*4*) 0.8, and (*5*) 1.0. The other dimensionless parameters of the problem are $\mu = 0.5$, $\varphi = 0$, $t = 0.05$, and $h = 0.2$. The dimensionless parameters and the junction resistance R are defined in the text.

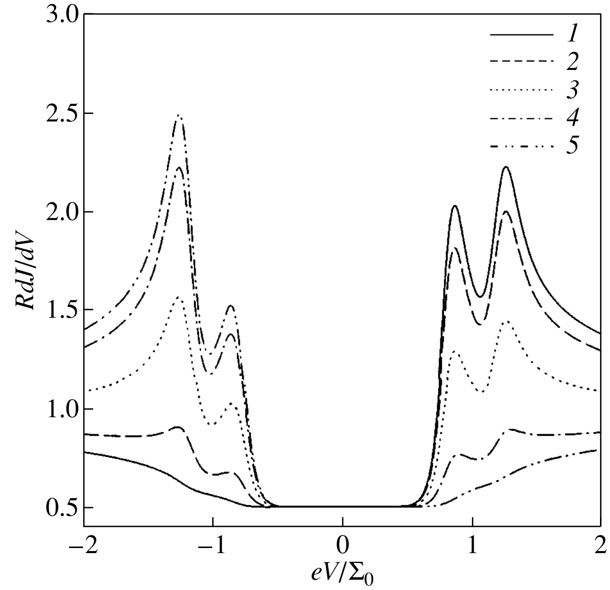

Fig. 3. Dependences of the dimensionless conductance $RG(eV/\Sigma_0)$ of the FM–CDWM tunnel junction for different phases of the order parameter of the charge density wave φ = (*1*) 0, (*2*) $\pi/4$, (*3*) $\pi/2$, (*4*) $3\pi/4$, and (*5*) π. The other dimensionless parameters of the problem are $\mu = 0.5$, $P = 0.2$, $t = 0.05$, and $h = 0.2$.

Similar dependences for a junction of the same type at fixed values of h and P and different phases of the order parameter of the charge density wave are shown in Fig. 3. It can be seen from this figure that the phase of the order parameter of the charge density wave substantially affects the form of the current–voltage characteristic. If a set of microtips considerably contributing to the total tunneling current were to be characterized by a random phase distribution, the corresponding terms (13) would be compensated upon averaging and, consequently, there should appear a curve at $\varphi = \pi/2$. Another interesting feature (Fig. 3) is the violation of the generalized symmetry relationships obtained earlier:

$$J(V, -\Sigma, P = 0) = -J(-V, \Sigma, P = 0), \qquad (23)$$

$$G(V, -\Sigma, P = 0) = G(-V, \Sigma, P = 0). \qquad (24)$$

As can be seen from Fig. 3, this violation occurs at any polarization $P \neq 0$.

A combined distortion of the density of states in the presence of two order parameters Δ and $\tilde{\Sigma}$ leads to an interesting correlation between different contributions to the conductance $G(V)$. This is illustrated in Fig. 4 for the case where the CDW metal is a superconductor, $\varphi = 0$, and $\delta_0 \equiv \Delta_0/\Sigma_0 = 0.5$. Here, Δ_0 stands for the amplitude of the bare superconducting gap when the electron–hole pairing is absent ($T_d = \Sigma = 0$). It should be emphasized that the inequality $H < H_p^{\mathrm{CDWM}}$ always holds true in our calculations. As can be seen from Fig. 4, both types of features observed in the current–voltage characteristics due to the gap-induced (square-root) singularities of the density of states at the energies D and Δ are split. This leads to a combination of the results obtained in the framework of the Tedrow–Meservey method and our approach. Note also that only the conductance peaks associated with the gap Δ appear in the negative branch of the current–voltage characteristic.

The current–voltage characteristics of the FM–I–CDW superconductor junction (the same as in Fig. 4) for different phases φ are shown in Fig. 5. It can be seen from this figure that, as in the current–voltage characteristics depicted in Fig. 3, the features associated with the charge density waves "pass" from the positive voltage branch into the negative voltage branch with an increase in the phase φ. In view of continuing discussions regarding the nature of the pseudogap in high-temperature superconducting oxides [9, 28, 29], it is of interest to attempt to reveal the simultaneous spin splitting of both the gap and pseudogap features in the quasiparticle current–voltage characteristics of cuprates with structures similar to those whose properties are illustrated in Figs. 4 and 5.

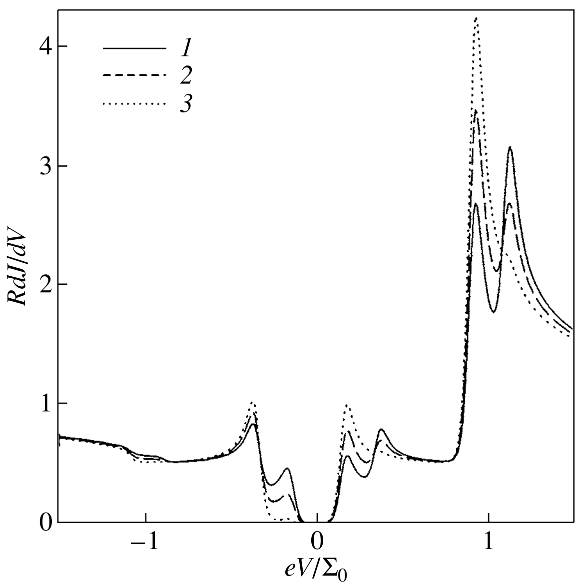

Fig. 4. Dependences of the dimensionless conductance $RG(eV/\Sigma_0)$ of the FM–CDWS tunnel junction for polarizations P = (*1*) 0.1, (*2*) 0.5, and (*3*) 0.9. The dimensionless parameters of the problem are $\mu = 0.5$, $\varphi = 0$, $t = 0.02$, $h = 0.1$, and $\delta_0 = \Delta_0/\Sigma_0 = 0.5$, where Δ_0 is the dimensionless superconducting gap for the CDWS electrode at zero temperature in the absence of dielectrization of the electron spectrum.

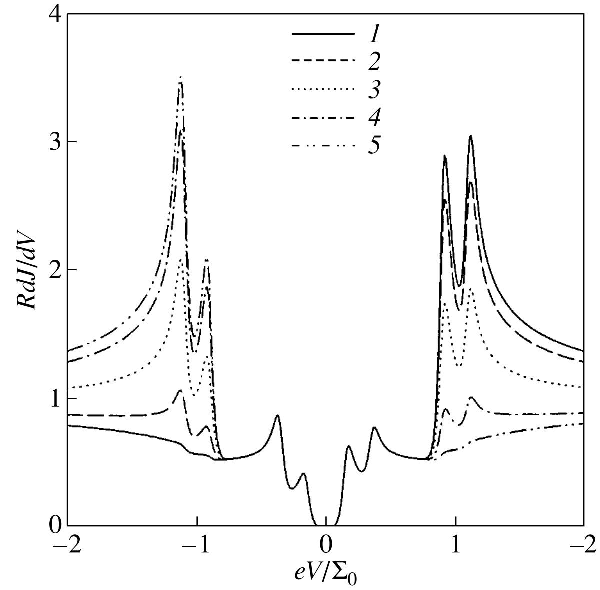

Fig. 5. Dependences of the dimensionless conductance $RG(eV/\Sigma_0)$ of the FM–CDWS tunnel junction for different phases of the order parameter of the charge density wave φ = (*1*) 0, (*2*) $\pi/4$, (*3*) $\pi/2$, (*4*) $3\pi/4$, and (*5*) π. The dimensionless parameters of the problem are $P = 0.2$, $t = 0.02$, $h = 0.1$, and $\delta_0 = 0.5$.

3.2. Symmetric Junctions

Figure 6 shows three typical current–voltage characteristics for the CDWM–I–CDWM symmetric tunnel junction at different order parameter phases φ_r for the right electrode ($\varphi_l = 0$ is taken for the phase of the left electrode). In particular, the current–voltage characteristic at $\varphi_r = 0$ or π corresponds to the excitonic insulator state. The current–voltage characteristic at $\varphi_r = \pi/2$ describes two physically different situation: (i) the situation with tunneling between Peierls insulators at the aforementioned intermediate phase and (ii) the situation where the tunneling current is collected from several one-dimensional chains and the contribution dependent on the phase φ_r completely disappears after the averaging. It should be remembered that the charge density waves are considered in the self-consistent field approximation, i.e., without regard for any truly pseudogap effects, such as the partial decrease in the density of states in strongly correlated systems due to fluctuations of the order parameter of charge density waves above the critical temperature T_d [30]. Similar features of the pseudogap type in the tunneling between chain conductors were investigated, for example, by Matveenko and Brazovskii [31].

According to the theory described in the previous section, the peaks of the conductance $G(V)$ in Fig. 6 are split as a result of the paramagnetic effect. These peaks correspond to the transition of quasiparticles from the degenerate sections of the Fermi surface of the left elec-

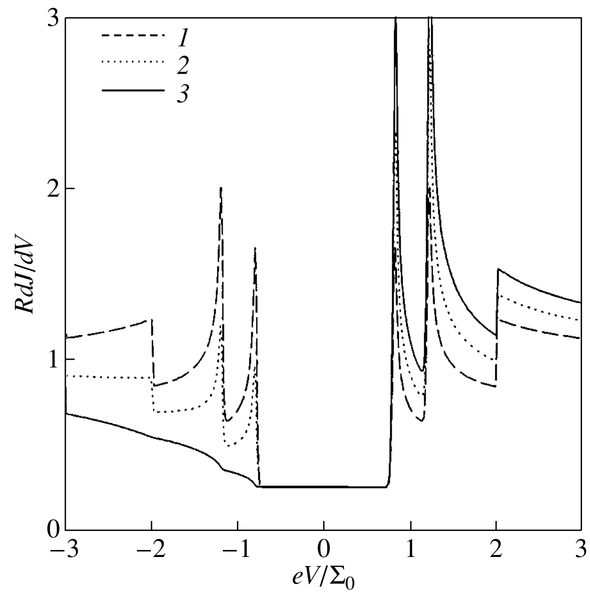

Fig. 6. Dependences of the dimensionless conductance $RG(eV/\Sigma_0)$ of the CDWM–CDWM nominally symmetric tunnel junction for different phases of the order parameter of the charge density wave in the right electrode φ_r = (*1*) 0, (*2*) $\pi/2$, and (*3*) π. The phase of the order parameter of the charge density wave in the left electrode is $\varphi_l = 0$. The dimensionless parameters of the problem are $\mu = 0.5$, $t = 0.01$, and $h = 0.2$.

trode to the nondegenerate sections of the Fermi surface of the right electrode, and vice versa (Fig. 1). This splitting is retained at arbitrary phases of the order parameters $\tilde{\Sigma}$ for both CDWM electrodes. The features that are attributed to the transitions of quasiparticles from the degenerate section of the Fermi surface of one CDWM electrode to the degenerate section of the Fermi surface of the other CDWM electrode remain unsplit.

4. CONCLUSIONS

Thus, we predicted two new types of tunnel junctions in which the peaks of the differential conductance should be split by a magnetic field. The splitting has a paramagnetic (spin) nature and occurs only when the following two conditions are satisfied: (i) the Fermi surface of the material of one electrode contains a correlation gap, and (ii) a correlation gap is absent at least in part of the Fermi surface of the other electrode. In this case, the electron spectrum of the first electrode can be dielectrized completely or partially.

Tunnel junctions of the first type are asymmetric metal–I–CDWM (CDWS) junctions. If the metal electrode is a ferromagnet, its polarization P affects the dependences $J(V)$ and $G(V)$. There exist several possible candidates for the formation of the appropriate structure. These are organic CDW metals, such as α-$(ET)_2M\mathrm{Hg(SCN)}_4$ (M = K, Tl, Rb) [32] and $Per_2[M(mnt)_2]$ (M = Au, Pt) [11].

The emergence of the superconductivity at temperatures $T < T_c < T_d$ in any specific CDW compound (see review [9]) can serve as a clear indication that this material is a metal rather than an insulator and, hence, can exhibit Zeeman splitting for an asymmetric configuration of the experiment. For example, low-dimensional metals with a CDW instability, such as $NbSe_3$, Nb_3Te_4, $Li_{0.9}Mo_6O_{17}$, $Tl_2Mo_6Se_6$, layered dichalcogenides, alloys with $A15$ and $C15$ structures, $Lu_5Ir_4Si_{10}$, $P_4W_{14}O_{50}$, tungsten bronzes doped with alkali metals, and $BaPb_{1-x}Bi_xO_3$ solid solutions, can be used as good objects for observation of the splitting. Another class of suitable materials includes superconducting cuprates.

Tunnel junctions of the second type that can exhibit the predicted effect are CDWM–I–CDWM symmetric junctions (it is significant that both electrodes should be normal metals!). It is worth noting that the splitting under consideration is absent for symmetric junctions between conventional Bardeen–Cooper–Schrieffer superconductors. The predicted splitting for this configuration is associated with the incomplete dielectrization of the electron spectrum of CDW metals. This situation is typical of many low-dimensional Peierls metals with an incommensurate charge density wave.

The calculations performed have demonstrated that the current–voltage characteristics obtained for tunnel junctions of both types depend on the phase of the order parameter of the charge density waves in the electrodes.

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