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Phase-coherent charge transport in superconducting heterocontacts

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Motivated by recent tunneling experiments for high-temperature superconductors, we have examined the effect of a near interface pair potential suppression on the corresponding transport characteristics. Our model structure consists of a tunnel junction between a normal injector and a normal conductor-superconductor bilayer with an interface of a finite transmittance. To study conductance spectra of mesoscopic heterostructures, a simple approach to the coherent charge transport through a double-barrier system proposed by Büttiker has been generalized to the case of an anisotropic superconducting order parameter. The impact of local fluctuations of the normal interlayer thickness on Andreev bound states, as well as the magnetic-field effect have been studied. Numerical calculations were carried out for *s*- and *d*-wave order parameters. The model provides a clear physical understanding of the near zero-bias conductance features in such heterostructures. It is shown that zero- and finite-bias anomalous conductance peaks in oxide superconductors often interpreted in terms of surface bound states associated with *d*-wave pairing could be also explained within the *s*-wave scenario. Some criteria to distinguish between two types of the pairing symmetry are proposed.
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I. INTRODUCTION

The pairing mechanism and the symmetry of the order parameter remain key questions for the physics of high-temperature superconductors (HTSC). During the last few years, a series of experimental techniques have been developed to test the orbital symmetry of the Cooper pair wave function and a great amount of evidences supporting the *d*-wave character of superconductivity in HTSC compounds has appeared.[1] At the same time, some results have been explained within the *s*-wave scenario. Because of the controversial situation,[2] new phase-sensitive approaches capable to identify experimentally the order parameter symmetry of HTSC are needed and their results have to be analyzed from the viewpoint of the realistic structure of high-T_c oxides.

Recently, the appearance of zero-bias conductance peaks (ZBCP) in the *ab*-plane tunneling spectra of HTSC was interpreted as a consequence of the angle dependent sign change of a *d*-wave order parameter on the Fermi surface (the so-called surface midgap states).[3] Such anomalies have been repeatedly displayed in the conductance spectra of a variety of HTSC-based junctions with a normal (N) injector and a number of explanations were proposed before.[4] Now these features [at least, those that are developing just at the superconducting critical temperature T_c (Ref. 5)] are usually attributed to the *d*-wave state and thus the appearance of ZBCP is considered as a direct confirmation of the non-*s*-wave character of the pairing potential.[3,6,7] Unfortunately, from the realistic viewpoint the situation is not so simple because of the unusually complicated structure of the objects investigated. In particular, in many cases the HTSC surface is covered by a thin surface layer that can be semiconducting like the BiO plane in Bi2212, or a superconductor with reduced critical temperature, or even a normal conductor as it is discussed for the CuO-chain plane in the YBCO compound.[8,9] The problem relates also to degradation processes in the upper atomic layers forming an oxygen depletion layer on the cuprate surface[10] and thus greatly influencing the contact properties of HTSC interfaces.[11] With the aim to form HTSC-based tunnel junctions, a significant part of the investigation uses just this degraded surface layer as the native insulating (I) barrier. Because of the complicated nature of both interface layers and barrier properties,[12] at the moment the presence of spatially inhomogeneous (and in many cases heterogeneous) interfaces seems to be unavoidable for HTSC junctions used in spectroscopic investigations. In this context, as it was emphasized in Ref. 9, it would be of interest to know how the tunneling spectrum (in particular, its inner-gap region) for an ideal metal-insulator-superconductor (N-I-S) junction is changed when a superconducting bulk has a thin nonuniform normal (N') covering on its surface. The corresponding analysis of such a system is just the main goal of this work.

The outline of the paper is as follows. We start with a brief description of existing theoretical approaches to the ZBCP problem in normal-metal–superconductor heterocontacts (Sec. II) and present our model based on the realistic structure of an HTSC-based junction that includes a degraded near-surface layer [Fig. 1(a)]. In Sec. III a method of calculating the differential conductance of a ballistic normal-metal–superconductor contact is given. We repeat the main lines of previous papers on this subject and generalize them to the case of an anisotropic superconducting order parameter. To illustrate the effect of realistic factors on the subgap features in tunneling curves and to present a simple understanding of some results obtained before, we apply the conductance formula to the case of an N-I-N'/S structure with an *s*-wave superconductor and present numerical data for different situations (Sec. IV). We try to find out the main features of the anomalies caused by the formation of Andreev bound states in the middle N' layer that have to be rejected before the conclusion about the *d*-wave origin of HTSC oxides is made. It has to be emphasized that such analysis can be applied in all cases when a superconductor is covered by a normal coating, e.g., to niobium-based junctions with specific barrier properties (see Refs. 13,14, and references therein) or to the high-current injection in N/S point contacts.[15] In Sec. IV we also deal with a *d*-wave gap symmetry and derive some conclusions relating the influence of a thin normal layer near the *d*-wave superconductor surface on the tunneling curves. Finally, the last section provides our conclusions and presents some simple criteria to distinguish between two types of pairing symmetry.

II. MODELS OF THE SUPERCONDUCTING HETEROSTRUCTURES

As it was noted before, the features we are dealing with are near-zero-voltage anomalies in the conductance G versus voltage V characteristics of superconducting heterostructures. Earlier work has found the widespread presence of ZBCP in the *ab*-plane tunneling curves for HTSC-based junctions[4] and some evidences for finite-bias peaks have been obtained as well.[16,17] In principle, such features were registered and studied extensively long before the discovery of high-T_c cuprates. First observations concerned tunneling junctions with transition metals as contact electrodes,[18] or with paramagnetic impurities in the insulating layer.[19] These findings were interpreted within a theory taking into account an exchange-scattering interaction between tunneling electrons and localized magnetic moments[20] (resonant tunneling through discrete energy levels in small quantum dots could lead to similar anomalies as well[21]). But in some cases it was not clear whether a ZBCP is caused by the presence of magnetic moments in the barrier, or it appears as a result of specific metal/insulator interface conditions (a situation similar to that in HTSC materials). For example, conductance anomalies observed in niobium oxide junctions were considerably decreased when a thin aluminum layer was added to the Nb surface before oxidation.[22]

An interest in zero-bias features was revived in the early 1990s after the experiments for mesoscopic superconductor/semiconductor[23,24] and superconductor/normal-metal[25] heterocontacts where a large conductance enhancement at $V=0$ was found. This effect for conventional *s*-wave superconductors was interpreted in a number of theoretical works.[26–33] In Refs. 26–29 the authors started from the clean metal case and took into account the elastic scattering in the normal side of the contact treated in terms of the transmission matrices across the junction. In the publications[30,31] the scattering-matrix method was applied to the ballistic N-I-N-I-S system and the similarity between the conductance spectra of a double-barrier structure and that of a disordered N/S junction was emphasized. Another approach to the problem[32,33] is based on the quasiclassical formulation of superconductivity in dirty inhomogeneous systems where the impurity average has already been done and where the mean free path in the normal layer is much less than its length, as well as the effective coherence length. But, as it was stated by Yip,[33] most of the experiments made before may not satisfy this condition and for them the corresponding results have to be regarded only as qualitative ones.

Motivated by the problem of ZBCP in HTSC tunneling characteristics, we present here a model that is very simplified but nevertheless reflects the main properties of the corresponding junctions. It originates from a physically insightful approach to the phase-coherent charge transport in a double-barrier normal metallic structure proposed by Büttiker in Ref. 34 and generalized below to the case of an anisotropic superconducting order parameter. To take into account the inhomogeneity of the superconducting base electrode, we approximate it by introducing an N'/S bilayer as, for example, it was already done by Di Chiara *et al.*[35] [Fig. 1(a)]. But on the contrary to Ref. 35 where the strength of the coupling between the bilayer slabs was assumed to be weak and the main aim was to investigate the proximity-induced superconductivity in the normal film, we focus on the wave interference pattern in the middle N' interlayer. We shall remain within the steplike approximation for the superconducting order parameter[36,37] and consider Andreev-retroreflected processes at the N'/S boundary, as well as normal scattering events of any strength, although the self-consistency of the spatial variation of the pair potential will be ignored. A similar program (but only for a clean interface) was realized in Ref. 38 for a normal-metal–*s*-wave superconductor junction and in Ref. 7 for a *d*-wave symmetry. But in the case of a clean interface, normal reflections are not taken into account, which is not true for HTSC materials where the characteristics of the surface degraded layer do radically differ from those of the superconducting bulk[11] and their interface acts as an elastic scatterer. Below we consider an N'/S interface of an arbitrary transmittance. The next step towards more sophisticated treatment of the HTSC electrode nonhomogeneity is an account of the local variations of the N' layer thickness and their impact on the corresponding conductance curves. Recently the same model but for the simplest one-channel approximation was successfully applied by Poirier *et al.*[39] to interpret their results for the subgap conductance of superconductor-GaAs junctions at very low temperatures. For the first time, they observed experimentally the crossover from a high-temperature peak at $V=0$ to finite-voltage low-temperature anomalies and the restoration of ZBCP under the magnetic-field application. To understand the influence of a magnetic field on the inner-gap tunneling spectra of superconducting heterostructures is one of the aims of this paper. In the following we do not claim an exhaustive study of the charge transport in the heterostructures but intend rather to provide a clear physical interpretation of near-zero-bias conductance anomalies for realistic inhomogeneous superconducting structures.

III. FORMALISM

Let us consider a normal-metal–superconductor junction in the quasi-one-dimensional geometry with M propagating transverse modes and the x axis as the interface normal. A specular-scattering nonsuperconducting N' part of the system ($0<x<l$) is connected to a normal-metal reservoir. The superconductor ($x>l$) with the pair potential depending on the direction of the traveling quasiparticle, as well as on its energy is connected to a superconducting reservoir. Further we shall neglect the energy dependence of the pair potential and thus restrict ourselves to the weak-coupling theory when the wave vector is fixed on the Fermi surface. For a superconducting order parameter the usual step-function approximation[36,37] will be assumed and the self-consistency of its spatial variation will be ignored. The Fermi wave numbers k_F and other electronic parameters will be equal in the normal and superconducting regions of the mesoscopic system. The normal reflections at the N'/S interface are taking into account and only ballistic transport limited to specular scattering by the N' region will be considered.

The voltage V applied to the barrier shifts the chemical potentials in both reservoirs and causes the current I in response to V. At zero temperature, the differential conductance $G(V)=dI(V)/dV$ is given by the well-known Landauer-type formula[40,41]

$$G(V)=\frac{2e^2}{h}\sum_{n=1}^{M}\left[1-|R_n^{ee}(\varepsilon)|^2+|R_n^{he}(\varepsilon)|^2\right]\Big|_{\varepsilon=eV}, \tag{3.1}$$

where $R_n^{he}(\varepsilon)$ is the scattering amplitude for an electron in the nth mode with an energy ε incident from left in the normal region and reflected back as a hole, $R_n^{ee}(\varepsilon)$ is the corresponding amplitude for its reflection as an electron; below the minimal value of the superconducting energy gap the relation $|R_n^{ee}(\varepsilon)|^2+|R_n^{he}(\varepsilon)|^2=1$ is valid for any fixed value of ε. In our model the Andreev reflection at the turning points 1 and 2 [Fig. 1(b)] and the normal scattering at the N'/S boundary ($x=l$) are separated spatially. The corresponding distance of the order of the superconducting coherence length is very small for oxide superconductors in comparison with the width of the degraded layer and thus its effect may be ignored in the calculations.

In this section we shall express the scattering amplitudes in terms of the reflection r and transmission t characteristics of the nonsuperconducting transitional region and the superconducting surface. In the first case an electron is scattered always only in electron states. On the contrary, at the Andreev reflection it can be retroreflected also as a hole with the same mode index that will be characterized further by a wave-vector $\mathbf{k}$. In the Andreev approximation[36] the corresponding characteristics for an electron scattered into a hole and vice versa are equal to

$$r^{he(eh)}(\mathbf{k},\varepsilon)=r(\mathbf{k},\varepsilon)\exp[\mp i\varphi(\mathbf{k})+i\psi(\mathbf{k})]. \tag{3.2}$$

In Eqs. (3.2) and (3.3) $r(\mathbf{k},\varepsilon)=[\varepsilon-\mathrm{sgn}(\varepsilon)\sqrt{\varepsilon^2-\Delta^2(\mathbf{k})}]/\varepsilon$, $\varphi(\mathbf{k})$ is the direction-dependent phase of the order parameter,

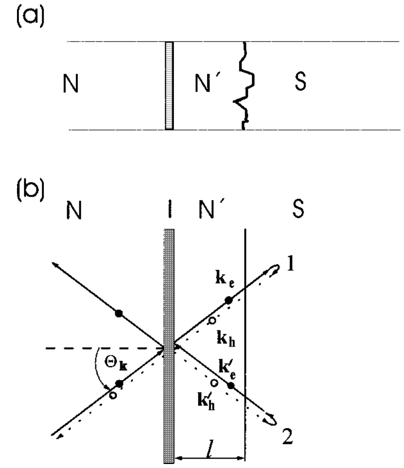


FIG. 1. Schematic view of the simulated model: (a) realistic structure of an HTSC-based junction including the degraded N' near-surface layer; (b) N-I-N'/S geometry used in the calculations. Points 1 and 2 are the turning points for the Andreev-like backscattering of the excitations penetrating into a superconductor at distances of the order of the superconducting coherence length. The numerical results presented below are obtained by averaging over fluctuations of the normal interlayer width l with a uniform distribution between two finite values.

and $\psi(\mathbf{k})$ describes the effect of a surface current in a superconductor. Below the minimal value of the superconducting energy gap we have the following relations:[37]

$$r^{he(eh)}(\mathbf{k},\varepsilon)=\exp\{-i\arccos[\varepsilon/|\Delta(\mathbf{k})|]\mp i\varphi(\mathbf{k})+i\psi(\mathbf{k})\}. \tag{3.3}$$

Usually $\psi(\mathbf{k})\equiv 0$ but if a stationary magnetic field B is applied parallel to the interface along the z axis the formula (3.2) and (3.3) provide the boundary conditions for a nonlinear differential equation determining the phase shift after the Andreev-like scattering.[13,14] Because in any case in the real experimental conditions the near-surface behavior of the superconducting properties is unknown we shall try only to estimate the corresponding value. For the field penetration depth λ much greater than the superconducting coherence length $\xi(\mathbf{k})$, as it is in HTSC compounds, the width of the layer with a depressed pair potential is of the order of $\xi(\mathbf{k})$ in the direction considered. Inside the superconductor the vector potential that can be chosen to be perpendicular to the interface normal as well as to the field direction equals to $A(x)=-B\lambda\exp[(l-x)/\lambda](x>l)$. Then the additional magnetic-field dependent part of the phase shift occurring by a charge after scattering from the superconductor is given by the following relation:

$$\psi^{(B)}=\int[\pm e\mathbf{A}(\mathbf{r})/\hbar]d\mathbf{s}=2\pi B/B_0(\mathbf{k})\sin(\Theta_\mathbf{k}), \tag{3.4}$$

where $B_0(\mathbf{k})=\Phi_0/[\lambda\xi(\mathbf{k})]$, $\Phi_0\equiv h/2e$, $\Theta_\mathbf{k}$ is the injection angle [see Fig. 1(b)], the sign of the phase shift (3.4) depends on the field orientation.

Let us now discuss the charge transport in the whole heterostructure shown in Fig. 1. An electron incident from the left reservoir with an energy ε and a wave vector $\mathbf{k}_e$ after transferring the normal transitional region with the transmission coefficient $t^e(\mathbf{k}_e,\varepsilon)$ will be reflected from the superconductor as a hole with the amplitude $r^{he}(\mathbf{k}_e,\varepsilon)$. The hole goes back along the path of the incoming electron with a wave vector $\mathbf{k}_h$ and then may get out of the N' region [the corresponding transmission amplitude is equal to $t_h(\mathbf{k}_h,\varepsilon)$] or could be reflected from it with a wave-vector $\mathbf{k}_h'$. The hole will be then retroreflected from the N'/S interface as an electron with an amplitude $r^{eh}(\mathbf{k}_h',\varepsilon)$ and will travel back with a wave-vector $\mathbf{k}_e'$. Again, the electron will leave the N' region or its wave vector will be transformed into $\mathbf{k}_e$ after normal scattering. These processes are schematically shown in Fig. 1(b) (see also Fig. 7 in Ref. 15). The scattering events described above together with multiple reflections from the middle N' layer [their amplitudes are equal to $r^h(\mathbf{k}_h,\varepsilon)$ and $r^e(\mathbf{k}_e',\varepsilon)$] determine $R^{he}(\mathbf{k}_e,\varepsilon)$ and $R^{ee}(\mathbf{k}_e,\varepsilon)$ scattering amplitudes. Taking into account all possible paths of the quasiparticle excitations, we obtain

$$R^{he}(\mathbf{k}_e,\varepsilon)=\frac{t^h(\mathbf{k}_h,\varepsilon)r^{he}(\mathbf{k}_e,\varepsilon)t^e(\mathbf{k}_e,\varepsilon)}{1-r^{he}(\mathbf{k}_e,\varepsilon)r^e(\mathbf{k}_e',\varepsilon)r^{eh}(\mathbf{k}_h',\varepsilon)r^h(\mathbf{k}_h,\varepsilon)}; \tag{3.5}$$

$$R^{ee}(\mathbf{k}_e,\varepsilon)=r^e(\mathbf{k}_e,\varepsilon)+\frac{t^e(\mathbf{k}_e',\varepsilon)r^{eh}(\mathbf{k}_h',\varepsilon)r^h(\mathbf{k}_h,\varepsilon)r^{he}(\mathbf{k}_e,\varepsilon)t^e(\mathbf{k}_e,\varepsilon)}{1-r^{eh}(\mathbf{k}_h',\varepsilon)r^h(\mathbf{k}_h,\varepsilon)r^{he}(\mathbf{k}_e,\varepsilon)r^e(\mathbf{k}_e',\varepsilon)}. \tag{3.6}$$

Here the amplitudes $t^{e(h)}$ and $r^{e(h)}$ include all phase shifts obtained during electron (hole) traveling through the normal transitional region. The scattering from the normal region assumed to be specular changes only the sign of the normal component k_x, whereas the Andreev reflection transforms the electron with k_x^e into a hole of k_x^h, where $k_x^e-k_x^h=2\varepsilon/\hbar v_{Fx}$, v_{Fx} is the x component of the Fermi velocity, the energy is conserved in all scattering processes.

To study the voltage region near $V=0$ we substitute the corresponding relations (3.3) into (3.5) and (3.6). For a given mode with a wave-vector $\mathbf{k}$ we obtain

$$G(\mathbf{k},V)=\frac{4e^2}{h}\left.\frac{|t^e(\mathbf{k}_e,\varepsilon)|^2|t^h(\mathbf{k}_h',\varepsilon)|^2}{1+|r^e(\mathbf{k}_e,\varepsilon)|^2|r^h(\mathbf{k}_h,\varepsilon)|^2-2\,\mathrm{Re}\{e^{i\phi}r^e(\mathbf{k}_e',\varepsilon)r^h(\mathbf{k}_h,\varepsilon)\}}\right|_{\varepsilon=eV} \tag{3.7}$$

with

$$\phi=-\arccos[\varepsilon/|\Delta(\mathbf{k}_e)|]-\arccos[\varepsilon/|\Delta(\mathbf{k}_h')|]+\varphi(\mathbf{k}_h')-\varphi(\mathbf{k}_e)+\psi(\mathbf{k}_h')+\psi(\mathbf{k}_e) \tag{3.8}$$

[compare with Eq. (16) in Ref. 31]. In the quasi-one-dimensional approximation the summation over the transverse modes can be replaced by an integration over the injection angle $\Theta_{\mathbf{k}}$, and for the normalized conductance spectrum we arrive at

$$\sigma(V)=\frac{G(V)}{g(V)}=\frac{\int d\Omega\, G(\mathbf{k},V)\cos\Theta_{\mathbf{k}}}{\int d\Omega\, g(\mathbf{k},V)\cos\Theta_{\mathbf{k}}} \tag{3.9}$$

with the normal-state value

$$g(\mathbf{k},V)=\frac{2e^2}{h}|t^e(\mathbf{k},eV)|^2. \tag{3.10}$$

For an isotropic s-wave superconductor where the effective pair potential does not depend on the wave-vector direction $\Delta(\mathbf{k})=\Delta_s=\mathrm{const}$, $\varphi(\mathbf{k}_h')=\varphi(\mathbf{k}_e)$ and in the absence of magnetic field ϕ is equal to $-\pi$ at $\varepsilon=0$. Taking into account that at the Fermi energy electron and hole excitations coincide we obtain the following expression:

$$G_s(\mathbf{k},0)=\frac{4e^2}{h}\frac{|t^e(k_F,0)|^4}{[1+|r^e(k_F,0)|^2]^2}=\frac{4e^2}{h}\frac{|t^e(k_F,0)|^4}{[2-|t^e(k_F,0)|^2]^2} \tag{3.11}$$

[see Eq. (5) in Ref. 28]. From Eq. (3.11) it follows that for any ballistic normal-superconducting junction the normalized zero-bias conductance height $\sigma_s(0)$ cannot exceed the value of two that is reached only for an ideal situation when $|t^e(k_F,0)|^2=1$ and $g(0)=2e^2M/h$. For anisotropic superconductors the situation changes dramatically. Let us illustrate it for a $d_{x^2-y^2}$-wave superconductor with a (110) oriented surface when the most pronounced anomalies in the conductance spectrum have been predicted.[3,6,7] In this case $\phi(\mathbf{k}_h')-\phi(\mathbf{k}_e)=\pi$ for any $\mathbf{k}_e$ and without magnetic fields $\phi=0$ for $\varepsilon=0$. Then from Eq. (3.7) for any interface between a normal injector and a d-wave superconductor we find that

$$G_d(\mathbf{k},0)=\frac{4e^2}{h}\frac{|t^e(\mathbf{k},0)|^4}{(1-|r^e(\mathbf{k},0)|^2)^2}=\frac{4e^2}{h}. \tag{3.12}$$

The formula looks like an expression for a normal-metal–s-wave superconductor junction with a clean interface and can be regarded as an example of a giant backscattering peak in superconducting heterostructures.[42] Because all modes bring in the same contribution, the zero-bias value of the differential conductance is equal to $G_d(0)=4e^2M/h$ whereas the normal-state conductance $g(0)$ can be very small for small transmission coefficients $|t^e(\mathbf{k},0)|^2$ and thus a giant zero-bias value $\sigma_d(0)$ appears for a (110) surface. Similar anomalies take place for other superconductor orientations, except the {100} tunneling direction. We have to emphasize that in the case of anisotropic superconductors this result strongly depends on the nature of the quasiparticle reflection from the N' region and any disorder will radically suppress the anomalous ZBCP.[43]

IV. CONDUCTANCE ANOMALIES IN N-I-N'/S JUNCTIONS

Now we shall specify the nature of the transitional N' layer. To make it closer to real HTSC experiments we have to take into account that: (i) the *ab*-plane tunneling usually studied is two-dimensional; (ii) the width l of the mediate N' layer has to be greatly nonuniform; (iii) l is known to be greater than the coherence length; (iv) the coherence length is two orders smaller than the *ab*-plane magnetic field penetration depth λ. The normal N' region will be described as a well between an electrostatic barrier and the scattering N'/S interface. Rough surfaces are assumed to have parallel plain regions of sizes more greater than the distance between them. A distribution of the last value is considered to be uniform between two finite values. Even if the surface regions are slightly nonparallel in real conditions it will not have any great effect because a finite mean free path or superconducting fluctuations in the N' region will prevent a charge to scatter many times from the interlayer boundaries.

The effect of the elastic scattering at the interface at $x=l$, as well as that of the electrostatic barrier at $x=0$ will be accounted by introducing repulsive potentials $V_{L,R}(x)$ of a delta-functional form (the main distinction between these two scattering planes is that the whole voltage bias is applied to the barrier but not to the N'/S boundary because of the shortenings in the interface). The corresponding reflection and transmission amplitudes do not depend on the energy and are given by[40]

$$r^{e}_{L,R}(\mathbf{k})=(r^{h}_{L,R}(\mathbf{k}))^{*}=-Z_{L,R}/(Z_{L,R}-i\cos\Theta_{\mathbf{k}}); \quad (4.1)$$

$$t^{e}_{L,R}(\mathbf{k})=(t^{h}_{L,R}(\mathbf{k}))^{*}=-\cos\Theta_{\mathbf{k}}/(Z_{L,R}-i\cos\Theta_{\mathbf{k}}), \quad (4.2)$$

where $Z_{L,R}=\int V_{L,R}(x)dx/\hbar v_{Fx}$ is the dimensionless barrier strength, L and R denote the characteristics of the left (I-N') and right (N'/S) boundaries of the N' layer, respectively.

If the normal transitional region may be simulated by a single potential barrier with a strength Z then the t's and r's in Eq. (3.7) are replaced by Eqs. (4.1) and (4.2), respectively, and the voltage behavior of $G(\mathbf{k},V)$ is determined exclusively by the energy dependence of ϕ (3.8). As it follows from Eq. (3.7), for an s-wave conductance spectrum there is a local minimum at $V=0$ that produces a zero-bias dip in the tunneling characteristics of conventional superconductors.[44] For one-dimensional geometry we obtain the well-known expression of Blonder-Tinkham-Klapwijk (BTK) theory (Table II in Ref. 40)

$$G_s(V)=\frac{4e^2}{h}\frac{\Delta_s^2}{\Delta_s^2(2Z^2+1)^2-4e^2Z^2(Z^2+1)V^2}, \quad |eV|<\Delta_s. \quad (4.3)$$

For a finite value of l and specular scattering interfaces reflection and transmission amplitudes describing the effect of the intermediate normal layer are given by

$$r^{e(h)}(\mathbf{k}_{e(h)},\varepsilon)=r_L^{e(h)}(\mathbf{k}_{e(h)})+\frac{[t_L^{e(h)}(\mathbf{k}_{e(h)})]^2r_R^{e(h)}(\mathbf{k}_{e(h)})\exp[2i\chi_{e(h)}(\mathbf{k}_{e(h)})]}{1-r_L^{e(h)}(\mathbf{k}_{e(h)})r_R^{e(h)}(\mathbf{k}_{e(h)})\exp[2i\chi_{e(h)}(\mathbf{k}_{e(h)})]}; \quad (4.4)$$

$$t^{e(h)}(\mathbf{k}_{e(h)},\varepsilon)=\frac{t_L^{e(h)}(\mathbf{k}_{e(h)})t_R^{e(h)}(\mathbf{k}_{e(h)})\exp[i\chi_{e(h)}(\mathbf{k}_{e(h)})]}{1-r_L^{e(h)}(\mathbf{k}_{e(h)})r_R^{e(h)}(\mathbf{k}_{e(h)})\exp[2i\chi_{e(h)}(\mathbf{k}_{e(h)})]}, \quad (4.5)$$

where $\chi_{e(h)}$ is the phase shift acquired by an electron (hole) traveling between two interlayer boundaries. Without magnetic fields $\chi_{e(h)}(\mathbf{k}_{e(h)})=k_x^{e(h)}l$. A stationary magnetic field parallel to the N'/S interface enters the amplitudes (4.4) and (4.5) via its penetration into the N' interlayer. For $0<x<l A(x)=-B\lambda[1+(l-x)/\lambda]$ and we easily find the corresponding magnetic-field dependent part of the phase shift,

$$\chi^{(B)}_{e(h)}=\pi Bl^2/[B_0(\mathbf{k})\lambda\,\xi(\mathbf{k})]\tan(\Theta_{\mathbf{k}}), \quad (4.6)$$

that is proportional to the magnetic flux entering the normal interlayer and thus is gauge invariant. Together with the impact of superconducting screening currents taking into account by Eq. (3.4) it determines the field effect on the interference pattern.

A. s-wave superconductors

We shall start with the case of an s-wave superconductor where $\Delta(\mathbf{k})\equiv\Delta_s$. If r and t characteristics of the intermediate normal region are energy independent we again obtain a zero-bias dip[38] with a value $\sigma(0)$ that is independent on l. The last result follows from the fact that at the Fermi energy electron and hole excitations coincide and the phase shift due to the electron traveling along the region of the width l is canceled by the corresponding quantity for a hole state into which the electron was transformed. If scattering characteristics of the normal region have a strong dependence on the energy, the conductance spectrum displays usually a sharp fine structure and changes from a peak at $V=0$ to a dip by small variations of the junction parameters (see, for example, the numerical simulations[31] for an N-I-N-I-S double-barrier junction). This behavior seems unrealistic and does not correspond to experimental findings that are, as a rule, smooth curves repeatedly reproduced in certain experimental conditions. In our opinion, the discrepancy can be removed if to take into account the real conditions mentioned above, especially, the nonuniformity of the intermediate normal layer width. The results of the averaging over fluctuations of l do not maintain any unusual sharp features and are shown in Fig. 2. It demonstrates how strongly the near-zero-bias feature depends on the interrelation between the N'/S interface scattering strength and the potential barrier height that can be a cause of dissimilar results obtained by different groups for the same objects. Really, with increasing the barrier height the subgap behavior of the conductance spectrum transforms

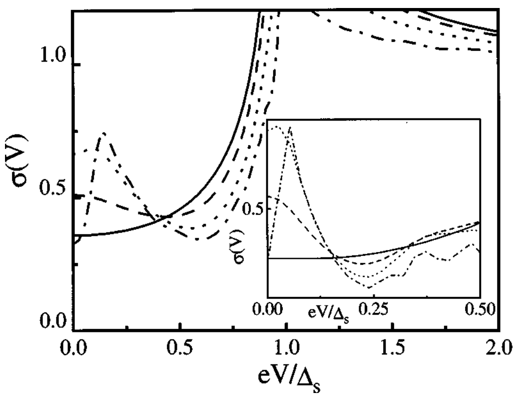
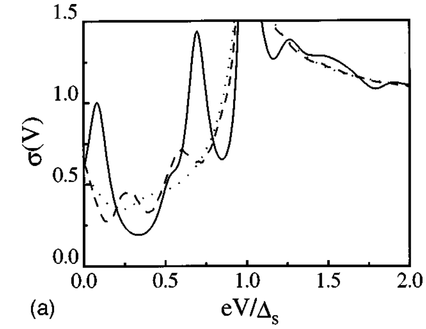


FIG. 2. The normalized conductance versus voltage for an N-I-N'/S junction with an s-wave superconductor. The width of the N' layer is uniformly distributed within the interval $l=(1.5–4.5)\xi_0$ and $l=(5–15)\xi_0$ (the inset). The scattering characteristics of the N'/S interface and the barrier are $Z_R=1.0$ and $Z_L=0.1$ (solid line), 0.5 (dashed line), 1.0 (dotted line), and 2.0 (dashed-dotted line), respectively; $k_F=10^{10}\,\mathrm{m}^{-1}$, $\xi_0=\hbar v_F/\pi\Delta_s$.

from an ideal tunneling curve for a conventional superconductor to a ZBCP and after that to a finite-bias peak near $V=0$.

Existence of two characteristic types of the conductance spectra observed in Fig. 2 can be understood by analyzing the initial curves. According to Eq. (3.7) there are two factors governing the conductance spectra of the N-I-N'/S system: the nominator and the denominator.[31] The first one is important when the main normal scattering occurs at the N'/S interface and the total transmittance of the system is not too high. For a quasi-one-dimensional situation near the Fermi energy we obtain $|t^h(\mathbf{k}_h',\varepsilon)|^2=|t^e(\mathbf{k}_e,-\varepsilon)|^2$ and the nominator that is a product of these two functions has a maximum at $V=0$ as a result of the electron-hole symmetry. The peak survives after the averaging procedure (Fig. 2) and even more, as it follows from Fig. 3(a) the averaging over l supports the existence of a zero-bias conductance enhancement. The denominator plays the main role in Eq. (3.7) for a weakly transparent barrier and for great distances l between the barrier and the N'/S interface. It results in appearance of bound states in the intermediate normal layer and, as a sequence, a system of conductance peaks with a dip pinned at zero voltage in $G(\mathbf{k},V)$ [Fig. 3(b)]. Positions of the states inside the energy gap depend on the thickness l and thus vary from point to point in our model. The most insensitive peak with respect to the value of l is the first one and just it remains after the averaging procedure together with a dip at $V=0$ [Fig. 3(b)]. The higher harmonics are smeared out or at least are not as prominent as the maximum nearest to $V=0$. For temperatures k_BT comparable with the first bound-state energy (about several K for a contact with a normal interlayer width l of about $10\xi_0$ and an s-wave superconductor with T_c near 100 K) the dip at $V=0$ will disappear and a ZBCP will be detected in the experimental curves. In Ref. 39 where l was estimated to be about 100 nm the corresponding temperatures were an order less. When the temperature increases further, a nonmonotonous behavior of a zero-voltage conductance should be observed with a maximum at finite temperature in accordance with the data of Refs. 39 and 45 for mesoscopic structures based on conventional superconductors.



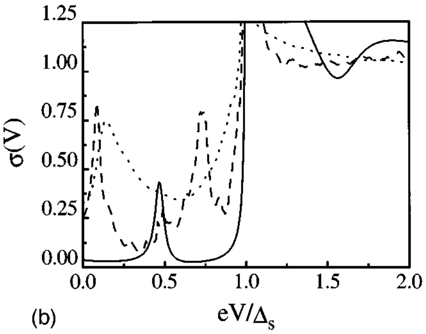
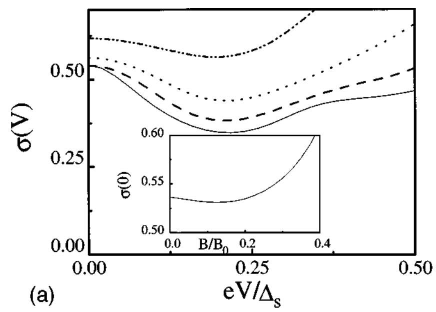


FIG. 3. The normalized conductance versus voltage for a one-dimensional N-I-N'/S junction with an s-wave superconductor (solid line), the multichannel two-dimensional system (dashed line), and the result of the averaging over the middle normal layer width (dotted line): $Z_R=1.0$ and $Z_L=0.5$, $l=(10\pm5)\xi_0$ (a), $Z_R=1.0$ and $Z_L=2.0$, $l=(3\pm1.5)\xi_0$ (b); $k_F=10^{10}\,\mathrm{m}^{-1}$, $\xi_0=\hbar v_F/\pi\Delta_s$.

Our model gives also clear physical insight into the phenomenon of reflectionless tunneling that was observed experimentally in semiconductor/superconductor[23,24] and normal-metal/superconductor[25] heterostructures. In Ref. 26 it was explained within a semiclassical description taking into account the presence of scatterers in the normal side that deterministically deflect electrons and holes. The authors found that the reduction in current caused by the scattering at the interface with a superconductor is compensated by the current increase due to multiple reflections in the disordered region (see also more correct results within the scattering-matrix approach obtained for a disordered N-I-S junction in Refs. 28 and 29). Within our approach the deterministic scattering back to the N'/S interface means that the strength of the tunneling barrier I is extremely great. In such case by doubling the value of l the system considered may be conventionally transformed into the $S/N'/S$ structure with two identical N'/S interfaces. The interference inside the N' layer yields striking Fabry-Perot-like resonances[46,47] independently on the scattering conditions at the interfaces. It is just the case of the reflectionless tunneling discussed in Ref. 26 with the only difference concerning positions of the resonances. Because of missing the additional phase changes (3.8) connected with the Andreev reflection a ZBCP was predicted in Ref. 26 whereas we obtain finite-voltage features (of course, for a great width of the disordered region they should appear near to $V=0$ and can be overlooked at sufficiently great temperatures).



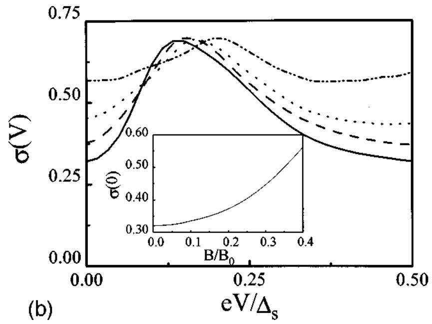


FIG. 4. Magnetic low-field effect on the normalized conductance spectrum for an N-I-N'/S junction with an s-wave superconductor. $B/B_0=0$ (solid line), 0.2 (dashed line), 0.3 (dotted line), 0.4 (dashed-dotted line); $B_0=\Phi_0/(\lambda\xi_0)$; $Z_R=1.0$, $Z_L=0.5$, $l=(10\pm5)\xi_0$ (a), and $Z_R=1.0$, $Z_L=2.0$, $l=(3\pm1.5)\xi_0$ (b); $k_F=10^{10}\,\mathrm{m}^{-1}$, $\lambda=100$ nm, $\xi_0=\hbar v_F/\pi\Delta_s$. The insets show zero-bias conductance behavior.

As it was established in Refs. 30 and 31 the differential conductance for a ballistic double-barrier N-I-N-I-S structure exhibits features similar to those obtained for a disordered N-S contact with a potential barrier at the interface. Our model even more resembles the last system because the averaging over different realizations of the impurity potential is modeled in our case by averaging over different distances l. In accordance with our simulations, Fig. 2 in Ref. 29 shows that for a high N-S interface resistance the conductance vs voltage curve has a peak at $V=0$ that is transformed into a minimum with a relative enhancement at finite voltages by increasing the N-S interface transparency (similar results were obtained by Yip[33] within the quasiclassical Green's-function approach).

The effect of low magnetic fields shown in Fig. 4 is governed by two factors (3.4) and (4.6) with the relative impact depending on the ratio $l^2/(\lambda\xi)$. Although the formal structures of Eqs. (3.4) and (4.6) coincide, for the s-wave symmetry the additional term equal to $-\pi$ appears in Eq. (3.8) and the phase shifts (3.4) and (4.6) provide opposite effects. Whereas the screening current enlarges the near-zero-bias conductance it diminishes under the influence of the field penetration into the intermediate N' region. For $l^2\ll\lambda\xi$, as it is in HTSC compounds, the low-field effect consists mainly in increasing the conductance value at $V=0$ (Fig. 4). The interplay between two factors resulting in a plateau in the $\sigma(0)$ vs B dependence can be clearly seen in Fig. 4(b). A significant effect in s-wave superconductors is achieved in fields above B_0 that can be as great as 10 T or even more for s-wave superconductors with parameters typical for HTSC compounds. For ordinary superconductors with a small field penetration depth l must be replaced with $l+\xi_0$ and for great

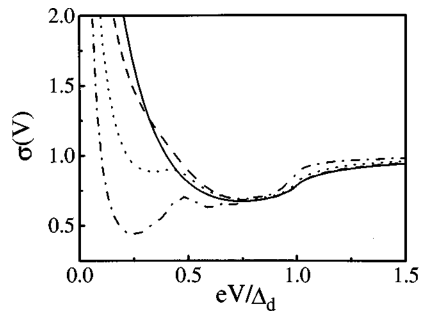


FIG. 5. The normalized conductance versus voltage for an N-I-N'/S junction with a (110) oriented d-wave superconductor: $Z_R=1.0$ and $Z_L=0.1$ (solid line), 0.5 (dashed line), 1.0 (dotted line), and 2.0 (dashed-dotted line), respectively, $l=(3\pm1.5)\xi_0$; $k_F=10^{10}\,\mathrm{m}^{-1}$, $\xi_0=\hbar v_F/\pi\Delta_d$.

values of the coherence length the fields needed can be reduced up to several hundreds of Tesla. The main impact of higher fields is to suppress as constructive, as destructive interference in the system and thus the conductance spectra become featureless that results in a peak in the field dependence of $\sigma(0)$. It corresponds to the results obtained by Yip[33] who studied the transport through a heterogeneous superconducting system with paramagnetic impurities on the normal side within the quasiclassical Green's-function technique and those of Marmorkos *et al.*[29] for a disordered N-S junction. They obtained that for some ''large'' voltages the differential conductance of the high-resistance barrier junction increases with increasing pair breaking. The simulations[29] also demonstrate the existence of a peak at a finite field value for appropriate parameters. The same results follow from our model. They are not presented in this paper because such fields seem to be unattainable in HTSC experiments although the peak can be observed for conventional superconductors (see the experimental data of Ref. 39 for a semiconductor-superconductor junction where a ZBCP observed below 4.2 K was transformed into finite-voltage peaks at lower temperatures and then restored at intermediate magnetic fields).

B. *d*-wave superconductors

Our numerical calculations for a $d_{x^2-y^2}$-wave superconductor where $\Delta(\mathbf{k})=\Delta_d\cos(2\Theta_{\mathbf{k}})$ (the angle $\Theta_{\mathbf{k}}$ is measured relative to the crystalline axis along which the d-wave order parameter reaches maximum) demonstrate the peak at zero-bias voltage in all cases except the (100) oriented surfaces. The ZBCP is maximal for the (110) oriented surface and equals to $G_d(0)=4e^2M/h$ for M different transport channels (Fig. 5). The effect appears because of: (i) the phase conjugation between electrons and holes at the Fermi energy, (ii) the fact that the Andreev-reflected hole always traces back the path of the incoming electron, (iii) zero phase difference between electron and Andreev-scattered hole excitations at the Fermi energy for the nodal tunneling direction [see Eq. (3.8) for $\psi(\mathbf{k}'_h)=\psi(\mathbf{k}_e)=0$], and (iv) the specular reflection from the normal region assumed above. The last restriction is crucial just for a d-wave state. As it follows from Eq. (3.8) a ZBCP appears when the states belonging to the neighboring lobes with the π phase shift are involved in the scattering process. If the surface roughness mixes the $\mathbf{k}_h$ and $\mathbf{k}'_h$ states from the same lobe then even in the {110} tunneling direction

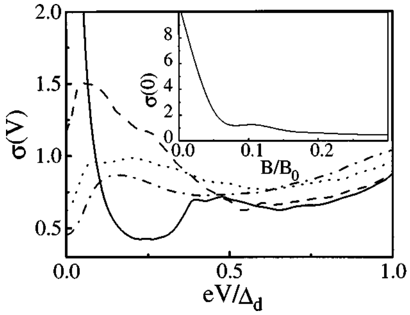


FIG. 6. Magnetic low-field effect on the normalized conductance spectrum of an N-I-N'/S junction with a (110) oriented d-wave superconductor. $B/B_0=0$ (solid line), 0.1 (dashed line), 0.2 (dotted line), 0.3 (dashed-dotted line); $B_0=\Phi_0/(\lambda\xi_0)$; $Z_R=1.0$, $Z_L=2.0$, $l=(3\pm1.5)\xi_0$; $k_F=10^{10}\,\mathrm{m}^{-1}$, $\lambda=100\,\mathrm{nm}$, $\xi_0=\hbar v_F/\pi\Delta_d$. The inset shows zero-bias conductance behavior.

an s-wave component will appear as well as a ZBCP for (100) oriented surfaces. Corresponding discussion within the framework of the quasiclassical Eilenberger equations was given in Refs. 43 and 48.

The well-known results[3,6] for d-wave superconductors without an intermediate normal region are changed if the scattering characteristics of the normal transitional part of the system strongly depend on the energy (as it was shown for the case of a clean N'/S interface in Ref. 7). The existing ZBCP becomes narrower with increasing the strength of the potential barrier (Fig. 5) and with increasing l (not shown here). At the appropriate conditions, two dips on both sides of the central huge peak appear. According to Cucolo[5] a simultaneous presence of these features in experimental curves is an intrinsic feature of ZBCP anomalies developed at $T=T_c$ and thus can be considered as an indication of the unusual nature of superconductivity in HTSC oxides. As in Refs. 3 and 6 the only orientation for which the conductance-voltage curve has no peak at $V=0$ corresponds to the $\{100\}$ tunneling direction. Its dependence on the parameters, as well as the magnetic-field effect is similar to that for an s-wave superconductor. On the contrary, for (110) oriented specimens the ZBCP always decreases with the magnetic field, and its fall is dramatic already at small B (Fig. 6). It is so because in the d-wave case there is no π factor in the phase shift (3.8) and on the contrary to the s-pairing situation both field effects (3.4) and (4.6) act in the same direction and greatly reduce the conductance value at $V=0$. The characteric fields are near $0.1B_0$ or about 1 T for the cuprate parameters. Together with decreasing the zero-bias conductance value finite-bias peaks appear in the d-wave conductance spectrum and shift to higher voltages with increasing fields (see also Ref. 48).

V. CONCLUSIONS

We have examined systematically the influence of a thin normal covering on the inner-gap conductance spectrum of superconductors. The circumstances introduced in a simple approach[34] describing coherent tunneling through a double-barrier structure include local variations of the intermediate N' layer width, the anisotropy of the superconducting order parameter, and the magnetic-field effect.

A general formula for the spectral conductance of the N-I-N'/S junction was analyzed for a superconductor in s- and d-wave cases. It was shown that in the s-wave state the averaging curves always reveal a ZBCP or two symmetrical maxima around $V=0$ divided by a dip centered at zero-bias voltage. These peaks are fingerprints of the first bound states in the initial (nonaveraged) characteristics. If the experiments are made at sufficiently great temperatures the fine structure will be smeared out and as in Ref. 39 the resulting characteristic will reveal a peak at $V=0$. For $d_{x^2-y^2}$-wave superconductors in all cases (except the tunneling in the $\{100\}$ direction) a ZBCP with robustness to the changes of the system parameters was obtained. If Andreev bound states are formed in the normal interlayer two symmetrical dips at finite voltages appear in accordance with the experimental results.[5] The ZBCP for a d-wave superconductor dramatically decreases and splits for magnetic fields below 1 T as it can be estimated for high-temperature cuprates. In real conditions the field values to reveal the effect discussed must be significantly lower because all conclusions made above are valid only below the first critical field.

We would like to point out that although the self-consistency of the order parameter[49] has been ignored, strictly speaking, the conductance spectrum near $V=0$ will not be modified significantly because the zero-bias value does not depend on the distance l and hence on the spatial behavior of the superconducting pair potential. This statement is more adequate to the case of a normal-metal/s-wave superconductor bilayer with a small-transmissivity scattering plane (see the corresponding simulations in Ref. 50). Calculations for d-wave superconductors made within the quasiclassical formalism[51] also demonstrate that main modifications arising due to the account of the self-consistency occur near the energy gap value. In the d-wave case more important simplification seems to be the near-parallel elastic-scattering planes bounded the degraded near-surface layer. In this sense our results for a ZBCP have to be considered only as extremal values achieved for ideal structures. In any case, the model presented cannot claim on quantitative results but, as we hope, it gives a clear physical understanding of the results obtained by more sophisticated methods, as well as, of the experimental data for heterostructures made from conventional superconductors.[39] It can be also easily generalized to more complicated situations as the point-contact spectroscopy when the translational symmetry is removed, inelastic events and phase randomization,[34] or a superconducting counter electrode (see corresponding experimental results in Ref. 52).

Concerning HTSC-based junctions, it is necessary to emphasize that an appearance of a ZBCP itself is insufficient for proving the unconventional nature of the pair potential symmetry in HTSC because unusual aspects of the conductance spectra for superconducting oxides could be in principle reproduced within the framework of the realistic model taking into account the presence of a randomly distributed normal layer between normal and s-wave superconducting bulks. Moreover, in Ref. 53 the authors have proved experimentally a correlation between the anomaly at $V=0$ in YBCO films and existence of a degraded surface layer in this material. To make a final conclusion about the order parameter symmetry, we need more sophisticated criteria to distinguish between two types of pairing symmetries. Below we propose such criteria for a ZBCP easily verified experimentally without any detailed comparison with the curves theoretically predicted:

(i) *The amplitude of the ZBCP*. As it was emphasized in Sec. III, the normalized conductance $\sigma(0)$ for an *s*-wave superconductor cannot exceed the factor of two and this value is realized only for an ideal N/S contact when the width of a near-zero-bias enhancement is of the order of the superconducting energy gap. A narrower and higher ZBCP with a greater height can serve as a first indication of the *d*-wave scenario. This result is of general character and does not relate directly to the original BTK theory.[40]

(ii) *The angular dependence of the ZBCP*. The next check on the unconventional pairing relates to the drastic effect of the tunneling direction on the ZBCP for a *d*-wave situation. First of all, it must be observed only for *ab*-tunneling curves and, if so, the most prominent anomaly has to appear for the misorientation angle α of the crystalline x axis equal to 45°. Of course, in a realistic situation it is difficult to be sure that the main tunneling direction coincides with the normal to the injector/superconductor interface. Thus it would be sufficient to deal with the highest zero-bias peak.

(iii) *The ZBCP dependence on the barrier parameters.* In Sec. III it was shown that for a *d*-wave orbital symmetry the zero-bias value for the misorientation angle equal to 45° is not influenced by modification of scattering properties of the insulating layer. On the contrary, in the *s*-wave case the conductance spectrum strongly depends on the barrier strength. This test can be realized, e.g., in scanning tunnel microscope spectroscopic experiments by changing the tip-specimen distance.

(iv) *The magnetic field effect on the ZBCP*. Because of the two-dimensional character of the charge transport in oxide superconductors the conductance spectrum has to modify radically for magnetic fields rotating in the plane perpendicular to the *ab*-plane. Any strong dependence upon the field orientation is not expected for the *s*-wave symmetry. The next check originates from the different impact of the fields in two kinds of superconductors that was discussed in Sec. IV. For low fields (they must be an order greater for *s*-wave materials than for the *d*-wave ones) it results in a ZBCP enhancement for an *s*-wave material (Fig. 4) and in a ZBCP suppression in the *d*-wave state accompanied with splitting into two symmetrical peaks (Fig. 6).

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[1] Reviews on experiments supporting a $d_{x^2-y^2}$ symmetry order parameter in HTSC compounds may be found in D. J. Van Harlingen, Rev. Mod. Phys. **67**, 515 (1995); B. G. Levi, Phys. Today **49**(1), 19 (1996); **50**(11), 19 (1997).

[2] An analysis of different experimental methods and controversial results obtained for the symmetry of the cuprate order parameter is presented by R. A. Klemm, in *High Temperature Superconductivity—Ten Years After Its Discovery*, edited by K. B. Garg and S. M. Bose (Narosa Publishing, New Delhi, 1998), p. 179.

[3] C.-R. Hu, Phys. Rev. Lett. **72**, 1526 (1994); J. Yang and C.-R. Hu, Phys. Rev. B **50**, 16 766 (1994).

[4] T. Walsh, Int. J. Mod. Phys. B **6**, 125 (1992).

[5] A. M. Cucolo, Physica C **305**, 85 (1998).

[6] Y. Tanaka and S. Kashiwaya, Phys. Rev. Lett. **74**, 3451 (1994); S. Kashiwaya, Y. Tanaka, M. Koyanagi, H. Takashima, and K. Kajimura, Phys. Rev. B **51**, 1350 (1995); S. Kashiwaya, Y. Tanaka, M. Koyanagi, and K. Kajimura, *ibid.* **53**, 2667 (1996).

[7] J. H. Xu, J. H. Miller, and C. S. Ting, Phys. Rev. B **53**, 3604 (1996).

[8] M. Ledvij and R. A. Klemm, Phys. Rev. B **52**, 12 552 (1995).

[9] K. Kitazawa, H. Sugawara, and T. Hasegawa, Physica C **263**, 215 (1996).

[10] J. Halbritter, Phys. Rev. B **46**, 14 861 (1992); **48**, 9735 (1993), and references therein.

[11] M. Grajcar, A. Plecenik, P. Seidel, V. Vojtanik, P. Kúš, and K.-U. Barholz, J. Low Temp. Phys. **106**, 277 (1997); P. Seidel, M. Grajcar, A. Plecenik, M. Belogolovskii, A. Matthes, and M. Zuzcak, Proc. SPIE **3480**, 67 (1998).

[12] R. Gross, L. Alff, A. Beck, O. M. Froehlich, D. Koelle, and A. Marx, IEEE Trans. Appl. Supercond. **7**, 2929 (1997).

[13] U. Gunsenheimer, A. Hahn, A. Krause, and S. V. Kuplevakhsky, Phys. Rev. B **54**, 6545 (1996).

[14] A. Hahn, S. Hofmann, A. Krause, and P. Seidel, Physica C **296**, 103 (1998).

[15] A. Hahn and K. Hümpfner, Phys. Rev. B **51**, 3660 (1995).

[16] J. Geerk, X. X. Xi, and G. Linker, Z. Phys. B **73**, 329 (1988).

[17] M. Covington, M. Aprili, E. Paraoanu, L. H. Greene, F. Xu, J. Zhu, and C. A. Mirkin, Phys. Rev. Lett. **79**, 277 (1997).

[18] A. F. G. Wyatt, Phys. Rev. Lett. **13**, 401 (1964).

[19] F. Mezei, Phys. Lett. **25A**, 534 (1967); A. F. G. Wyatt and D. J. Lythall, *ibid.* **25A**, 541 (1967).

[20] J. A. Applebaum, Phys. Rev. **154**, 633 (1966); J. A. Applebaum and L. Y. L. Shen, Phys. Rev. B **5**, 544 (1972).

[21] H. Schoeller and G. Schön, Physica B **203**, 423 (1994); J. König, J. Schmid, H. Schoeller, and G. Schön, Czech. J. Phys. **46**, Suppl S4, 2399 (1996).

[22] J. M. Rowell, M. Gurvitch, and J. Geerk, Phys. Rev. B **24**, 2278 (1981).

[23] A. Kastalsky, A. W. Kleinsasser, L. H. Greene, R. Bhat, F. P. Milliken, and J. P. Harbison, Phys. Rev. Lett. **67**, 3026 (1991); C. Nguyen, H. Kroemer, and E. L. Hu, *ibid.* **69**, 2847 (1992).

[24] P. H. C. Magnée, N. van der Post, P. H. M. Kooistra, B. J. van Wees, and T. M. Klapwijk, Phys. Rev. B **50**, 4594 (1994).

[25] P. Xiong, G. Xiao, and R. B. Laibowitz, Phys. Rev. Lett. **71**, 1907 (1993).

[26] B. J. van Wees, P. de Vries, P. Magnée, and T. M. Klapwijk, Phys. Rev. Lett. **69**, 510 (1992).

[27] Y. Takane and H. Ebisawa, J. Phys. Soc. Jpn. **60**, 3130 (1991); **61**, 1685 (1992); **61**, 2858 (1992); **62**, 1844 (1993).
[28] C. W. J. Beenakker, Phys. Rev. B **46**, 12 841 (1992).
[29] I. K. Marmorkos, C. W. J. Beenakker, and R. A. Jalabert, Phys. Rev. B **48**, 2811 (1993).
[30] J. A. Melsen and C. W. J. Beenakker, Physica B **203**, 219 (1994).
[31] G. B. Lesovik, A. L. Fauchère, and G. Blatter, Phys. Rev. B **55**, 3146 (1997).
[32] A. F. Volkov, A. V. Zaitzev, and T. M. Klapwijk, Physica C **210**, 21 (1993).
[33] S. Yip, Phys. Rev. B **52**, 15 504 (1995).
[34] M. Büttiker, IBM J. Res. Dev. **32**, 63 (1988).
[35] A. Di Chiara, F. Fontana, G. Peluso, and F. Tafuri, Phys. Rev. B **48**, 6695 (1993).
[36] A. F. Andreev, Zh. Eksp. Teor. Fiz. **46**, 1823 (1964) [Sov. Phys. JETP **19**, 1228 (1964)].
[37] R. Kümmel, Phys. Rev. B **10**, 2812 (1974).
[38] S. Chaudhuri and P. F. Bagwell, Phys. Rev. B **51**, 16 936 (1995).
[39] W. Poirier, D. Mailly, and M. Sanquer, Phys. Rev. Lett. **79**, 2105 (1997).
[40] G. E. Blonder, M. Tinkham, and T. M. Klapwijk, Phys. Rev. B **25**, 4515 (1982).
[41] C. J. Lambert, J. Phys.: Condens. Matter **3**, 6579 (1991).
[42] C. W. J. Beenakker, J. A. Melsen, and P. W. Brouwer, Phys. Rev. B **51**, 13 883 (1995).
[43] A. A. Golubov and M. Yu. Kupriyanov, JETP Lett. **67**, 501 (1998).
[44] E. L. Wolf, *Principles of Electron Tunneling Spectroscopy* (Oxford University, New York, 1985).
[45] P. Charlat, H. Courtois, Ph. Gandit, D. Mailly, A. F. Volkov, and B. Pannetier, Phys. Rev. Lett. **77**, 4950 (1996).
[46] G. E. Rittenhouse and J. M. Graybeal, Phys. Rev. B **49**, 1182 (1994).
[47] G. A. Gogadze and A. M. Kosevich, Fiz. Nizk. Temp. **24**, 716 (1998) [Low Temp. Phys. **24**, 540 (1998)].
[48] M. Fogelström, D. Rainer, and J. A. Sauls, Phys. Rev. Lett. **79**, 281 (1997).
[49] P. G. de Gennes, *Superconductivity of Metals and Alloys* (Benjamin, New York, 1966).
[50] R. A. Riedel, L.-F. Chang, and P. F. Bagwell, Phys. Rev. B **54**, 16 082 (1996).
[51] Yu. S. Barash, A. A. Svidzinsky, and H. Burkhardt, Phys. Rev. B **55**, 15 282 (1997).
[52] S. Sinha and K.-W. Ng, Phys. Rev. Lett. **80**, 1296 (1998).
[53] R. T. Kao, S. J. Wang, J. Y. Juang, K. H. Wu, T. M. Uen, Y. S. Gou, Physica C **282-287**, 1493 (1997).