

## Spin current in ferromagnet-insulator-superconductor junctions

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A theory of spin polarized tunneling spectroscopy based on a scattering theory is given for tunneling junctions between ferromagnets and  $d$ -wave superconductors. The spin filtering effect of an exchange field in the insulator is also treated. We clarify that the properties of the Andreev reflection are largely modified due to the presence of an exchange field in the ferromagnets, and consequently the Andreev reflected quasiparticle shows an evanescent-wave behavior depending on the injection angle of the quasiparticle. Conductance formulas for the spin current as well as the charge current are given as a function of the applied voltage and the spin polarization in the ferromagnet for arbitrary barrier heights. It is shown that the surface bound states do not contribute to the spin current and that the zero-bias conductance peak expected for a  $d$ -wave superconductor splits into two peaks under the influence of the exchange interaction in the insulator.

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### I. INTRODUCTION

The transport properties in hybrid structures between ferromagnets and superconductors have received considerable theoretical and experimental attention. Interest in such structures includes spin-dependent spectroscopy of superconductors and possible device applications. Since the Cooper pairs in spin singlet superconductors are formed between up and down spins, the high density of spin injection through a tunneling barrier induces a spin imbalance. This nonequilibrium state is expected to result in a suppression of the critical temperature and the critical current density in the superconductor. A large number of experimental studies on spin-polarized tunneling have already been performed using conventional metal superconductors such as Al and Nb about 20 years ago.<sup>1</sup> However, the recent discovery of so-called colossal magnetoresistance (CMR) in Mn oxides compound has aroused new interest in this field,<sup>2,3</sup> because hybrid structure fabrication of the spin-polarized ferromagnets with high- $T_c$  superconductors is now possible using these materials.<sup>4,5</sup>

On the other hand, the properties of ferromagnet-insulator-superconductor (F/I/S) and ferromagnet-ferromagnetic-insulator-superconductor (F/FI/S) junctions have been analyzed based on the assumption that the conductance spectra correspond to the density of states (DOS) of the superconductor weighted by the spin polarization.<sup>1,6,7</sup> A theory for FIS junctions based on a scattering method has been presented by de Jong and Beeneker,<sup>8</sup> new aspects of Andreev reflection have been revealed, and also detailed comparisons between theory and experiments have been accomplished.<sup>9,10</sup> However, these results are restricted to isotropic  $s$ -wave superconductors.

In contrast to  $s$ -wave superconductor cases, at the inter-

face of a  $d_{x^2-y^2}$ -wave superconductor, zero-energy states (ZES) are formed due to the interference effect of the internal phase of the pair potential.<sup>11</sup> Tunneling theory for  $d_{x^2-y^2}$ -wave superconductors has already been presented by extending the BTK formula<sup>12</sup> to include the anisotropy of the pair potential.<sup>13-15</sup> The theory predicts the existence of zero-bias conductance peak (ZBCP) which reflects the formation of the surface bound states on the  $d$ -wave superconductors. In this paper, an exchange interaction is introduced on the normal side of the junction and on the insulator in order to analyze the spin polarized tunneling effects. The bound-state condition and tunneling spectroscopy of ferromagnet- $d$ -wave-superconductor junctions have already been analyzed in two papers.<sup>16,17</sup> They have revealed several important features in charge transport. Here we will argue that the properties of the Andreev reflection<sup>18</sup> is largely modified due to the presence of the exchange interaction. In particular, the existence of an evanescent type of the Andreev reflection, which is referred to as virtual Andreev reflection (VAR), is explained (see Ref. 27). This process has significant roles on the transports especially for junctions between half-metallic ferromagnets and superconductors. The conductance formulas for the charge and the spin currents are presented based on the scattering method by fully taking account of the VAR process. The merit of formulas based on the scattering methods is that the conductance spectra can easily be calculated for arbitrary barrier heights cases without the restriction of the high-barrier limit. The spin current is, we believe, the most important physical quantity in spin injection devices based on the following two reasons: one is that the spin current gives a direct criterion to estimate the effect of the spin imbalance induced by the tunneling current, the other is that the charge and the spin conductivity may illuminate the

study of electron systems that undergo spin-charge separation, such as Tomonaga-Luttinger liquids and possibly underdoped high- $T_c$  superconductors.<sup>19–21</sup> We will also analyze ferromagnetic insulator effects, which include the spin-filtering effect,<sup>6,7</sup> due to the presence of an exchange field in the insulator. It is shown that a spin-dependent energy shift during the tunneling process induces a splitting of the ZBCP. Based on the detailed analysis of the conductance spectra, we propose a simple method to distinguish the broken time-reversal symmetry (BTRS) states induced at the surface<sup>22–24</sup> from spin-dependent tunneling effects. The implications of the ferromagnetic insulator effects on tunneling experiments of high- $T_c$  superconductors and a proposal for possible device applications are also presented.

## II. FORMULATION

For the model of formulation, a planar F/I/S junction with semi-infinite electrodes in the clean limit is assumed. A flat interface is assumed to be located at  $x=0$ , and the insulator for up [down] spin is described by a potential  $V_{\uparrow[\downarrow]}(\mathbf{x}) \{V_{\uparrow[\downarrow]}(\mathbf{x}) = (\hat{V}_0 - [+]\hat{U}_B)\delta(x)\}$ , where  $\delta(x)$ ,  $\hat{V}_0$ , and  $\hat{U}_B$  are the  $\delta$  function, a genuine barrier amplitude, and an exchange amplitude in the barrier, respectively. The effective mass  $m$  in the ferromagnet and in the superconductor are assumed to be equal. For the model of the ferromagnet, we adopt the Stoner model where the effect of the spin polarization is described by the one-electron Hamiltonian with an exchange interaction similar to the case of Refs. 8,16,17. For the description of the  $d_{x^2-y^2}$ -wave superconductor, we apply the quasiclassical approximation where the Fermi energy  $E_F$  in the superconductor is much larger than the pair potential following the model by Bruder.<sup>25,26</sup> The effective Hamiltonian (Bogoliubov–de Gennes equation) is given by

$$\begin{bmatrix} H_0(\mathbf{x}) - \rho U(\mathbf{x}) & \Delta(\mathbf{x}, \theta) \\ \Delta^*(\mathbf{x}, \theta) & -\{H_0(\mathbf{x}) + \rho U(\mathbf{x})\} \end{bmatrix} \begin{bmatrix} u(\mathbf{x}, \theta) \\ v(\mathbf{x}, \theta) \end{bmatrix} = E \begin{bmatrix} u(\mathbf{x}, \theta) \\ v(\mathbf{x}, \theta) \end{bmatrix}. \quad (1)$$

Here,  $E$  is the energy of the quasiparticle,  $U(\mathbf{x})$  is the exchange potential given by  $U\Theta(-x)$  ( $U \geq 0$ ), where  $\Theta(x)$  is the Heaviside step function,  $\rho$  is 1 (–1) for up (down) spins,  $\Delta(\mathbf{x}, \theta)$  is the pair potential, and  $H_0(\mathbf{x}) \equiv -\hbar^2 \nabla^2 / 2m + V(\mathbf{x}) - E_F$ . To describe the Fermi surface difference in  $F$  and  $S$ , we assume  $E_F = E_{FN}$  for  $x < 0$  and  $E_F = E_{FS}$  for  $x > 0$ . The pair potential  $\Delta(\mathbf{x}, \theta)$  is taken as  $\Delta(\theta)\Theta(x)$  for simplicity. The number of up (down) spin electrons is described by  $N_{\uparrow}$  ( $N_{\downarrow}$ ). The polarization and the wave vector of quasiparticles in the ferromagnet for up (down) spin are expressed as  $P_{\uparrow} \equiv N_{\uparrow} / (N_{\uparrow} + N_{\downarrow}) = (E_{FN} + U) / 2E_{FN}$  [ $P_{\downarrow} \equiv N_{\downarrow} / (N_{\uparrow} + N_{\downarrow}) = (E_{FN} - U) / 2E_{FN}$ ] and  $k_{N,\uparrow} = |\mathbf{k}_{N,\uparrow}| \equiv \sqrt{(2m/\hbar^2)(E_{FN} + U)}$  [ $k_{N,\downarrow} = |\mathbf{k}_{N,\downarrow}| \equiv \sqrt{(2m/\hbar^2)(E_{FN} - U)}$ ], respectively.<sup>8</sup>

We assume the quasiparticle injection of up spin electrons at an angle  $\theta_N$  to the interface normal as shown in Fig. 1. Four possible trajectories exist; they are Andreev reflection (AR), normal reflection (NR), transmission to superconductor as electronlike quasiparticles (ELQ), and transmission

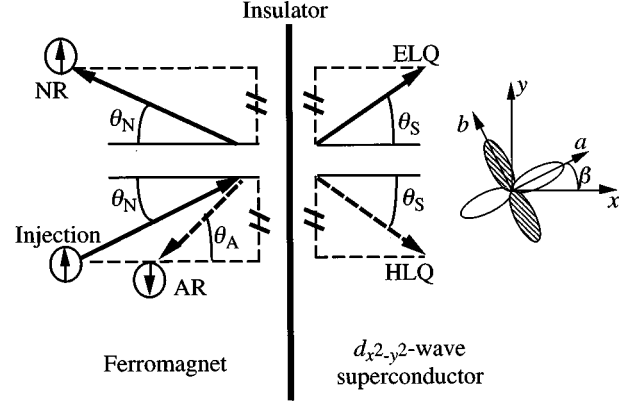


FIG. 1. Schematic illustration of the elastic reflection of quasiparticles in the FIS junction. For all trajectories, momenta parallel to the interface are conserved. This means that the retroreflection property of Andreev process is lost due to the exchange interaction. In the figure, the anisotropic pair potential of  $d_{x^2-y^2}$ -wave symmetry is also shown.

as holelike quasiparticles (HLQ). The spin direction is conserved for NR but not for AR. When the superconductor has  $d_{x^2-y^2}$ -wave symmetry, the effective pair potentials for ELQ and HLQ are given by  $\Delta_+ \equiv \Delta_0 \cos 2(\theta_S - \beta)$  and  $\Delta_- \equiv \Delta_0 \cos 2(\theta_S + \beta)$ , respectively, where  $\beta$  is the angle between the  $a$  axis of the crystal and the interface normal. Results for various pairing symmetries are obtained by setting proper values to  $\Delta_+$  and  $\Delta_-$  similarly to the previous formulas.<sup>13,15</sup> The wave vectors of ELQ and HLQ are approximated by  $k_S = |\mathbf{k}_S| \approx \sqrt{2mE_{FS}/\hbar^2}$  following the model by Andreev.<sup>18</sup> Since translational symmetry holds along the  $y$ -axis direction, the momentum components of all trajectories are conserved ( $k_{N,\uparrow} \sin \theta_N = k_{N,\downarrow} \sin \theta_A = k_S \sin \theta_S$ ).<sup>27</sup> Note that  $\theta_N$  is not equal to  $\theta_A$  except when  $U=0$ , which means retroreflectivity of AR is broken. Such novel behavior is a consequence of the fact that in the presence of an exchange field the BCS pairing is formed not strictly between states of equal but opposite  $k$  vectors, the so-called Fulde-Ferrell effect.<sup>28</sup> The wave function in the ferromagnet ( $x < 0$ ) for up (down) spin with injection angle  $\theta_N$  is described by

$$\begin{pmatrix} u(\mathbf{x}, \theta_N) \\ v(\mathbf{x}, \theta_N) \end{pmatrix} = e^{ik_{N,\uparrow} x} \begin{pmatrix} 1 \\ 0 \end{pmatrix} + a_{\uparrow[\downarrow]}(E, \theta_N) e^{ik'_{N,\uparrow} x} \begin{pmatrix} 0 \\ 1 \end{pmatrix} + b_{\uparrow[\downarrow]}(E, \theta_N) e^{ik'_{N,\uparrow} x} \begin{pmatrix} 1 \\ 0 \end{pmatrix}, \quad (2)$$

where the signs of the  $x$  components of  $\mathbf{k}_{N,\uparrow[\downarrow]}$  and  $\mathbf{k}'_{N,\uparrow[\downarrow]}$  are the reverse of each other. The reflection probabilities of the two processes are obtained by solving Eq. (1) and by connecting the wave function and its derivative at  $x=0$ .

Next, we will simply explain the Fermi surface effect by assuming up spin injection. Various kinds of reflection processes are expected depending on the values of  $E_{FN}$ ,  $E_S$ , and  $U$ . For example, when  $k_S < k_{N,\uparrow}$ , total reflection ( $|b_{\uparrow[\downarrow]}(E, \theta_N)|^2 = 1$ ) occurs when  $\theta_N > \sin^{-1}(k_S/k_{N,\uparrow}) \equiv \theta_{c1}$ .<sup>15,17</sup> In this case, the net currents of the spin and the charge from the ferromagnet to the superconductor vanish.

On the other hand, when  $k_{N,\downarrow} < k_S < k_{N,\uparrow}$ , the  $x$  component of the wave vector in the AR process ( $\sqrt{k_{N,\downarrow}^2 - k_S^2 \sin^2 \theta_S}$ ) becomes purely imaginary for  $\theta_{c1} > \theta_N > \sin^{-1}(k_{N,\downarrow}/k_{N,\uparrow}) \equiv \theta_{c2}$ . In this case, although transmitted quasiparticles from ferromagnet to superconductor do propagate, the Andreev reflected quasiparticles do not propagate (VAR process). A finite amplitude of the evanescent AR process still exists ( $|a_{\uparrow}(E, \theta_N)|^2 > 0$ ) and the net currents of the spin and the charge from the ferromagnet to the superconductor do not vanish. It is easy to check the conservation laws for the charge, the excitation, and the spin on the VAR process fol-

lowing the method presented in Ref. 12. The existence of the VAR process has not been treated in the one-dimensional model<sup>8</sup> because it is a peculiar feature of a two or three dimensional F/S interface.

The conductance of the junctions are obtained by extending previous formula to include the effect of spin.<sup>12,8,17</sup> In the following, consider a situation where  $k_{N,\downarrow} < k_S < k_{N,\uparrow}$ . To analyze the transport properties of an F/I/S junction, two kinds of conductance spectrum are introduced. The conductance for the charge current is defined by the charge flow induced by the up (down) spin quasiparticle injection and is given by

$$\hat{\sigma}_{q,\uparrow\downarrow}(E, \theta_N) \equiv \text{Re} \left[ 1 + \frac{\lambda_2}{\lambda_1} |a_{\uparrow\downarrow}(E, \theta_N)|^2 - |b_{\uparrow\downarrow}(E, \theta_N)|^2 \right] \quad (3)$$

(for  $0 < |\theta_N| < \theta_{c2}$ ),

$$= \frac{4\lambda_1 [4\lambda_2 |\hat{\Gamma}_+|^2 + (1 + \lambda_2)^2 + Z_{\downarrow\uparrow}^2 - |\hat{\Gamma}_+ \hat{\Gamma}_-|^2 \{(1 - \lambda_2)^2 + Z_{\downarrow\uparrow}^2\}]}{|(1 + \lambda_1 + iZ_{\uparrow\downarrow})(1 + \lambda_2 - iZ_{\downarrow\uparrow}) - (1 - \lambda_1 - iZ_{\uparrow\downarrow})(1 - \lambda_2 + iZ_{\downarrow\uparrow}) \hat{\Gamma}_+ \hat{\Gamma}_-|^2} \quad (4)$$

(for  $\theta_{c2} < |\theta_N| < \theta_{c1}$ ),

$$= \frac{4\lambda_1 (1 - |\hat{\Gamma}_+ \hat{\Gamma}_-|^2) \{1 + (-\kappa_2 + Z_{\downarrow})^2\}}{|(1 + \lambda_1 + iZ_{\uparrow}) \{1 - i(\kappa_2 + Z_{\downarrow})\} - (1 - \lambda_1 - iZ_{\uparrow}) \{1 + i(\kappa_2 + iZ_{\downarrow})\} \hat{\Gamma}_+ \hat{\Gamma}_-|^2} \quad (5)$$

(for  $\theta_{c1} < |\theta_N| < \pi/2$ ),

$$= 0,$$

where

$$Z_{\uparrow\downarrow} = \frac{Z_{0,\uparrow\downarrow}}{\cos \theta_S}, \quad Z_{0,\uparrow\downarrow} = \frac{2m(\hat{V}_0 - [ + ] \hat{U}_B)}{\hbar^2 k_S}, \quad \hat{\Gamma}_{\pm} = \Gamma_{\pm} \exp(\mp i \phi_{\pm}), \quad \exp i \phi_{\pm} = \frac{\Delta_{\pm}}{|\Delta_{\pm}|},$$

$$\Gamma_{\pm} = \frac{E - \sqrt{E^2 - |\Delta_{\pm}|^2}}{|\Delta_{\pm}|}, \quad \lambda_1 = \frac{k_{N,\uparrow\downarrow} \cos \theta_N}{k_S \cos \theta_S}, \quad \lambda_2 = \frac{k_{N,\downarrow\uparrow} \cos \theta_A}{k_S \cos \theta_S}, \quad \kappa_2 = i \lambda_2 = \frac{\sqrt{k_S^2 \sin^2 \theta_S - k_{N,\downarrow}^2}}{k_S \cos \theta_S}.$$

The conductance for the spin current is defined by the spin imbalance induced by the up (down) spin quasiparticle injection,

$$\hat{\sigma}_{s,\uparrow\downarrow}(E, \theta_N) \equiv \text{Re} \left[ 1 - \frac{\lambda_2}{\lambda_1} |a_{\uparrow\downarrow}(E, \theta_N)|^2 - |b_{\uparrow\downarrow}(E, \theta_N)|^2 \right] \quad (6)$$

(for  $0 < |\theta_N| < \theta_{c2}$ ),

$$= \frac{4\lambda_1 [-4\lambda_2 |\Gamma_+|^2 + (1 + \lambda_2)^2 + Z_{\downarrow\uparrow}^2 - |\hat{\Gamma}_+ \hat{\Gamma}_-|^2 \{(1 - \lambda_2)^2 + Z_{\downarrow\uparrow}^2\}]}{|(1 + \lambda_1 + iZ_{\uparrow\downarrow})(1 + \lambda_2 - iZ_{\downarrow\uparrow}) - (1 - \lambda_1 - iZ_{\uparrow\downarrow})(1 - \lambda_2 + iZ_{\downarrow\uparrow}) \hat{\Gamma}_+ \hat{\Gamma}_-|^2} \quad (7)$$

(for  $\theta_{c2} < |\theta_N| < \theta_{c1}$ ),

$$= \frac{4\lambda_1 (1 - |\hat{\Gamma}_+ \hat{\Gamma}_-|^2) \{1 + (-\kappa_2 + Z_{\downarrow})^2\}}{|(1 + \lambda_1 + iZ_{\uparrow}) \{1 - i(\kappa_2 + Z_{\downarrow})\} - (1 - \lambda_1 - iZ_{\uparrow}) \{1 + i(\kappa_2 + iZ_{\downarrow})\} \hat{\Gamma}_+ \hat{\Gamma}_-|^2} \quad (8)$$

(for  $\theta_{c1} < |\theta_N| < \pi/2$ ),

$$= 0.$$

The Andreev reflected quasiparticles positively contribute to the charge current, but since their spins are reversed, they have negative contribution to the spin current. The second terms on the right-hand side of Eqs. (3) and (6) do not have finite contribution on net current in the VAR process, since the corresponding  $\lambda_2$  is purely imaginary. The normalized total conductance spectra for the charge current  $\sigma_q(E)$  and the spin current  $\sigma_s(E)$  are given by

$$\sigma_q(E) = \sigma_{q,\uparrow}(E) + \sigma_{q,\downarrow}(E), \quad (9)$$

$$\sigma_{q,\uparrow[\downarrow]}(E) = \frac{1}{R_N} \int_{-\pi/2}^{\pi/2} d\theta_N \cos \theta_N \hat{\sigma}_{q,\uparrow[\downarrow]}(E, \theta_N) P_{\uparrow[\downarrow]} k_{N,\uparrow[\downarrow]}, \quad (10)$$

$$\sigma_s(E) = \sigma_{s,\uparrow}(E) - \sigma_{s,\downarrow}(E), \quad (11)$$

$$\sigma_{s,\uparrow[\downarrow]}(E) = \frac{1}{R_N} \int_{-\pi/2}^{\pi/2} d\theta_N \cos \theta_N \hat{\sigma}_{s,\uparrow[\downarrow]}(E, \theta_N) P_{\uparrow[\downarrow]} k_{N,\uparrow[\downarrow]}, \quad (12)$$

where

$$R_N = \int_{-\pi/2}^{\pi/2} d\theta_N \cos \theta_N [\hat{\sigma}_{N,\uparrow}(\theta_N) P_{\uparrow} k_{N,\uparrow} + \hat{\sigma}_{N,\downarrow}(\theta_N) P_{\downarrow} k_{N,\downarrow}], \quad (13)$$

$$\hat{\sigma}_{N,\uparrow[\downarrow]}(\theta_N) = \frac{4\lambda_1}{|1 + \lambda_1 + iZ_{\uparrow[\downarrow]}|^2}.$$

In the above,  $R_N$ ,  $\sigma_{q,\uparrow[\downarrow]}(E)$ , and  $\sigma_{s,\uparrow[\downarrow]}(E)$  correspond to the conductance when the superconductor is in the normal state and the spin-resolved normalized conductance spectra for charge and spin, respectively. The conductance spectra for negative  $E$  are obtained by replacing  $\hat{\Gamma}_{\pm}$  by  $-1/\hat{\Gamma}_{\pm}^*$  in the above formulation. The net polarization  $J_p(eV)$  as a function of the bias voltage  $V$  is given by

$$J_p(eV) = \frac{\int_{-\infty}^{\infty} dE \sigma_s(E) \{f(E-eV) - f(E)\}}{\int_{-\infty}^{\infty} dE \sigma_q(E) \{f(E-eV) - f(E)\}}, \quad (14)$$

where  $f(E)$  is the Fermi-distribution function. Since the convolution with  $f(E)$  gives only a smearing effect in the conductance spectra, the temperature is set to zero in the following discussions.

In the above formulation, we have neglected the self-consistency of the pair potential in order to get analytical formulas.<sup>29</sup> However, the present formula is easily extended to include this effect simply by replacing  $\Gamma_{\pm}$  with  $\Gamma_{\pm}(x)|_{x=0}$ , where  $\Gamma_{\pm}(x)$  follows the Riccati equations described by

$$\begin{aligned} \frac{d}{dx} \hat{\Gamma}_{\pm}(x) &= \frac{1}{i\hbar^2 k_F \cos \theta_S} [-\Delta_{\pm}(x) \hat{\Gamma}_{\pm}^2(x) - \Delta_{\pm}^*(x) \\ &+ 2E \hat{\Gamma}_{\pm}(x)], \end{aligned} \quad (15)$$

$$\begin{aligned} \frac{d}{dx} \hat{\Gamma}_{-}(x) &= \frac{1}{i\hbar^2 k_F \cos \theta_S} [-\Delta_{-}^*(x) \hat{\Gamma}_{-}^2(x) - \Delta_{-}(x) \\ &+ 2E \hat{\Gamma}_{-}(x)]. \end{aligned} \quad (16)$$

Here the spatial dependence of the pair potential is assumed as  $\Delta_{\pm}(x)$  (functions of  $x$ ).

The most important differences in the present formula from previous ones are (i) a formula for the nonlinear spin current, (ii) the capability to treat the ferromagnetic insulator effects based on the scattering method, (iii) the introduction of the breakdown in the retro-reflectivity of the AR process and consequently the vanishing of the propagating AR (VAR process). In particular, the concept of the VAR process is a new physical process presented in this paper. If we did not accept the existence of this process, the total reflection independent of  $E$  is naively expected. Since finite transmission is possible in this angle region ( $\theta_{c2} < |\theta| < \theta_{c1}$ ) above  $T_c$ , this total reflection would induce a sudden decrease of the conductance just below  $T_c$  for highly polarized ferromagnets junctions. As far as we know, no trends for such an effect have been reported thus far. This fact may be direct evidence for the existence of the VAR process. The VAR process is shown to have an important role on the Josephson current in superconductor-ferromagnet-superconductor junctions, because the evanescent wave carries a net Josephson current in this configuration.<sup>30</sup>

### III. RESULTS

#### A. Effects of polarization

In this subsection, to reveal the influence of the polarization on the tunneling conductance spectra, we assume a FIS junction by setting  $\hat{U}_B = 0$  ( $Z_{0,\uparrow} = Z_{0,\downarrow} = Z_0$ ). At first, let us discuss several analytical results obtained from the above formulation in order to check the validity of the formula. When  $U = 0$ , the ferromagnet reduces to a normal metal, and as expected  $\sigma_q(E)$  reproduces the results of Refs. 13,15, and  $\sigma_s(E)$  vanishes. For half-metallic ferromagnets ( $U = E_{FN}$ ), the Fermi surface for the down spins has shrunk to zero. In this case, the VAR process occurs for all  $\theta_N$ . Under the condition of VAR,  $\hat{\sigma}_q(E, \theta_N) = \hat{\sigma}_s(E, \theta_N)$  applies, which corresponds to the fact that the tunneling current is completely spin polarized. Furthermore, the conductance spectra in the energy gap ( $E < |\Delta_+|, E < |\Delta_-|$ ) become completely zero [ $\sigma_q(E, \theta_N) = \sigma_s(E, \theta_N) = 0$ ]. In the tunneling limit ( $H \rightarrow \infty$ ) and in the absence of VAR,  $\hat{\sigma}_{q,\uparrow[\downarrow]}(E, \theta)$  gives the angle resolved surface DOS of an isolated superconductor. Then  $\sigma_q(E)$  converges to the surface DOS weighted by the tunneling probability distribution.<sup>15</sup> At this limit, we can reproduce a well-known result that the ratio of the peak heights in the spin-resolved spectra directly reflect the polarization in the ferromagnet.<sup>1</sup> On the other hand,  $\sigma_{s,\uparrow[\downarrow]}(E, \theta)$  reduces to a function similar to the surface DOS, but where the divergence at the energy levels of the surface bound states is missing.

The calculated results based on above formula are presented for  $d_{x^2-y^2}$ -wave superconductors. In the following, we assume  $E_{FN} = E_{FS}$ . Figures 2 and 3 show the conductance spectra of charge current for the transparent limit ( $Z_0$

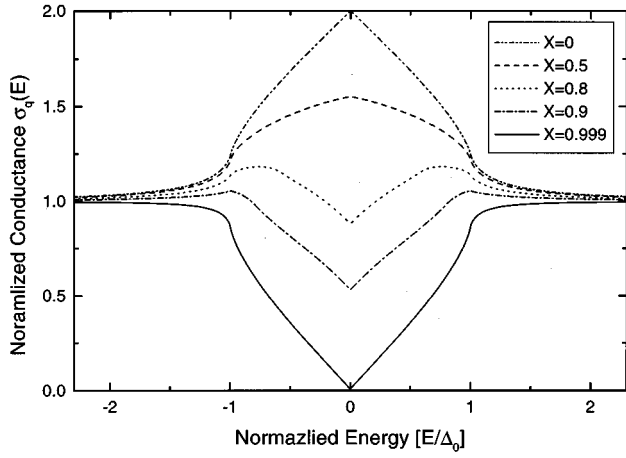


FIG. 2. The normalized conductance spectra for the charge current  $\sigma_q(E)$  as the function of  $X \equiv U/E_{FN}$  with  $\beta=0$  and  $Z_0=0$  (the transparent limit). As  $X$  becomes larger, the peak around the zero-bias level is largely suppressed.

$=0$ ,  $\beta=0$ ) and high-barrier case ( $Z_0=5$ ,  $\beta=\pi/4$ ) as the function of exchange interaction  $X (\equiv U/E_{FN})$ . For  $X=0$ , results in Ref. 13 are reproduced. However, as  $X$  increases, the conductance inside the gap ( $|E| < \Delta_0$ ) is largely reduced for both cases. Especially, the ZBCP disappears for the half-metallic ferromagnet case. Since the spin polarization has such a drastic influence on the ZBCP, the height of ZBCP can be used in principle as a measurement of the magnitude of the spin polarization. Figure 4 shows the difference of the spin current and the charge current when  $X=0.85$ ,  $Z_0=5$ , and  $\beta=\pi/4$ . It is clear that the ZBCP is not present for the spin current. This corresponds to the fact that the charge current components corresponding to the ZES are carried by condensed Cooper pairs in the superconductor, and therefore they do not contribute to the spin imbalance. As a result, the spin current becomes relatively insensitive to the orientation of the junctions. Figure 5 shows the conductance spectra for the spin current as the function of spin polarization ( $Z_0=5$ ). It is clear that the spin current increases as  $X$  becomes larger. Note that  $\sigma_s(E)$  is larger than unity around  $E=\Delta_0$  when  $X \approx 1$ . This corresponds to the fact that the peak in the DOS has an influence even for the spin current.

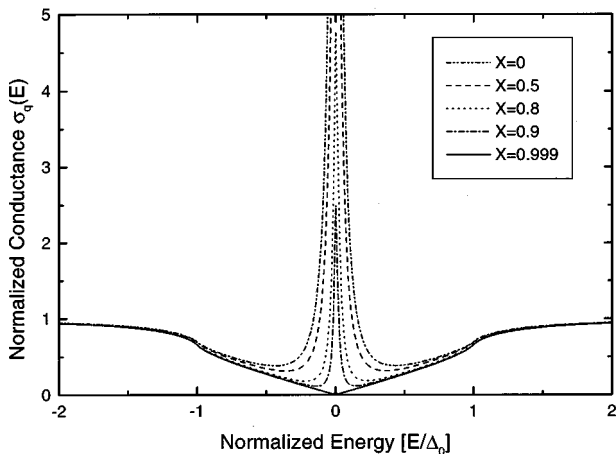


FIG. 3. The normalized conductance spectra for the charge current  $\sigma_q(E)$  as the function of  $X$  with  $\beta=\pi/4$  and  $Z_0=5$ . As  $X$  becomes larger, the height of the ZBCP is largely reduced.

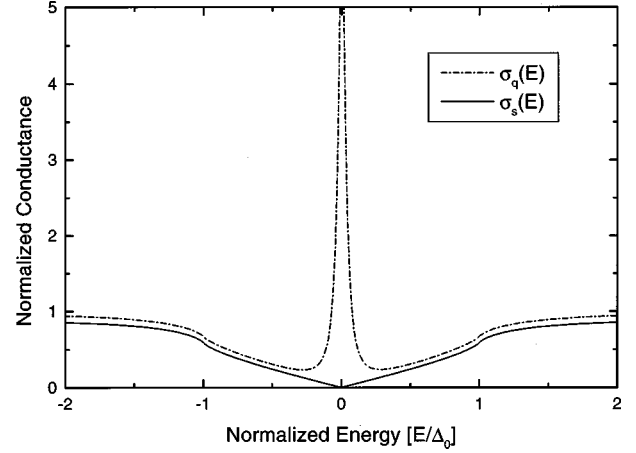


FIG. 4. The comparison between the normalized conductance spectrum for the charge current  $\sigma_q(E)$  and that for the spin current  $\sigma_s(E)$  for  $X=0.7$ ,  $\beta=\pi/4$ , and  $Z_0=5$ . Since the ZBCP originates from the current carried by the surface bound states, the peak disappears for the spin current.

Next, the net polarization  $J_p(eV)$  is calculated for  $d_{x^2-y^2}$ -wave superconductors as a function of the orientation ( $\beta$ ) when  $T=0$ . Four lines of Fig. 6 show the results for various values of barrier parameter when  $eV=2\Delta_0$ . It is clear that the orientational effect is much smaller compared to the effect of  $Z_0$ . In the same figure, results for  $s$ -wave superconductors ( $\Delta_+=\Delta_-=\Delta_0$  independent of  $\theta_N$ ) are also shown as closed dots. The large deviations of  $d_{x^2-y^2}$  wave from  $s$  wave for small values of  $Z_0$  are originated from the distribution of the pair amplitude in  $k$  space. As the barrier parameter becomes larger, the spin injection efficiency becomes insensitive to the symmetry of the pair potential.

### B. Spin filtering effects and the ZBCP splitting

It has been experimentally verified that a ferromagnetic semiconductor used as the insulator in tunneling junctions works as a ferromagnetic barrier. Since the transmission probabilities for up and down spins are not equal, a spin-filtering effect is expected to be realized.<sup>6,7</sup> Also it has been

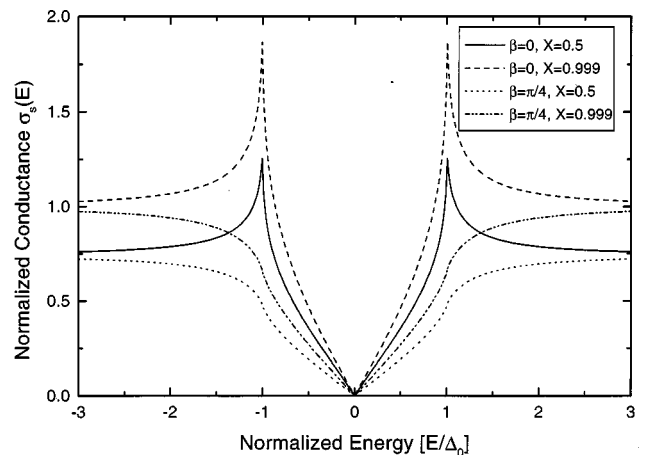


FIG. 5. The normalized conductance spectra for the spin current  $\sigma_s(E)$  as a function of  $X$  with  $\beta=0$  and  $Z_0=5$ . As  $X$  becomes larger, the spin current is increased. Note that the peak at  $E=\Delta_d$  is larger than unity when  $X$  is close to 1.

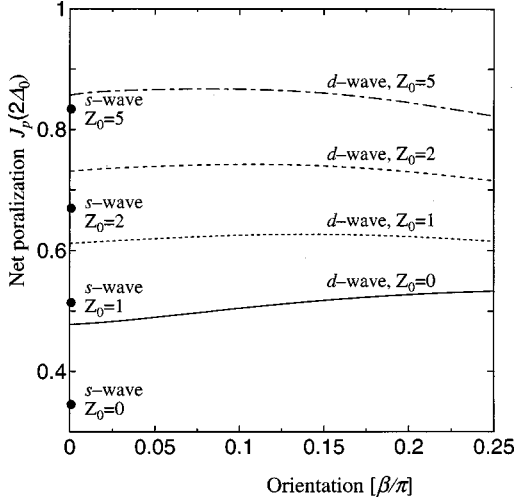


FIG. 6. Orientational dependencies of  $J_p(E)$  for  $d_{x^2-y^2}$ -wave superconductors for  $E=2\Delta_0$  and  $X=0.7$  are plotted for various  $Z_0$  values. Closed dots in the figures correspond to those for  $s$ -wave superconductors. The large deviations of the  $d_{x^2-y^2}$  wave from  $s$  wave for small values of  $Z_0$  are originated from the distribution of the pair amplitude in  $k$  space.

theoretically verified that a ferromagnetic insulator placed in the vicinity of superconductor induces a spin splitting on the DOS of  $s$ -wave superconductors.<sup>31</sup> In the following, we will analyze the influence of the exchange interaction existing inside the insulator on the transport properties based on the formulation described in Sec. II.

Figure 7 shows the response of the conductance spectra  $\sigma_q(E)$  on the exchange interaction in the insulator when  $X=0$ . ZBCP splittings are obtained for finite exchange amplitude ( $U_B$ ) cases. As  $U_B$  is increased and consequently as the difference between  $Z_{0,\uparrow}$  and  $Z_{0,\downarrow}$  becomes larger, the amplitude of the splitting becomes larger and the two peaks become broader and smaller. The peaks in the gap disappear when the difference between  $Z_{0,\uparrow}$  and  $Z_{0,\downarrow}$  becomes prominent. To see more clearly these trends, the spin-resolved con-

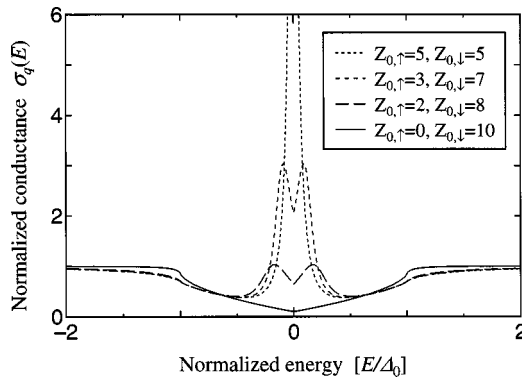


FIG. 7. The effects of ferromagnetic insulator on the charge current for  $X=0$  and  $\beta=\pi/4$ . When  $Z_{0,\uparrow}=5$  and  $Z_{0,\downarrow}=5$ , a large ZBCP exists. The difference in  $Z_{0,\uparrow}$  and  $Z_{0,\downarrow}$  induces the peak splitting ( $Z_{0,\uparrow}=3$  and  $Z_{0,\downarrow}=7$ ). As the difference becomes larger, the ZBCP split into two peaks and the amplitude of the splitting becomes larger and the peaks become smaller and broader ( $Z_{0,\uparrow}=2$  and  $Z_{0,\downarrow}=8$ ). Finally, the peaks in the gap disappear as the difference becomes prominent ( $Z_{0,\uparrow}=0$  and  $Z_{0,\downarrow}=10$ ).

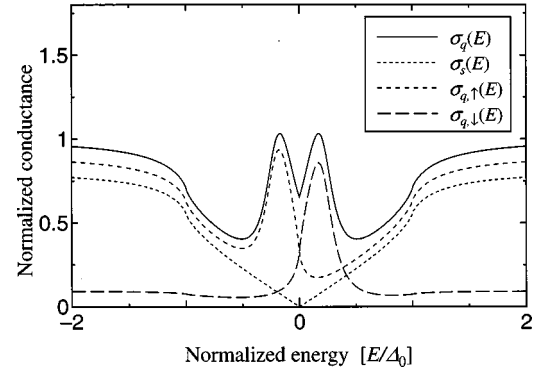


FIG. 8. Four types of normalized conductance for  $Z_{0,\uparrow}=2$ ,  $Z_{0,\downarrow}=8$ ,  $X=0$ , and  $\beta=\pi/4$ . The normalized charge current  $\sigma_q(E)$  has splitted peaks. The peak of the lower energy and higher energy are originated from the up spin component  $\sigma_{q,\uparrow}(E)$  and down spin component  $\sigma_{q,\downarrow}(E)$ . These peaks do not appear in the spin current conductance  $\sigma_s(E)$ .

ductance spectra  $\sigma_{q,\uparrow[\downarrow]}(E)$  and  $\sigma_s(E)$  for  $Z_{0,\uparrow}=2.5$  and  $Z_{0,\downarrow}=7.5$  are plotted in Fig. 8. The spectra for up (down) spins are shifted for lower (higher) energy level. Furthermore,  $\sigma_s(E)$  becomes finite even though  $X=0$  in the ferromagnet. In order to check the effect of the polarization, Fig. 9 shows the response of the charge current as a function of polarization  $X$  for a fixed barrier parameter. The spin polarization in the ferromagnet induces the imbalance of the peak heights, thus the ratio of the splitted peak heights can be used as a criteria for the spin polarization.

These results are interpreted as follows. (i) The peaks corresponding to the up (down) spin components are shifted because of the energy gain (loss) during the tunneling process. (ii) Since this energy gain (loss) has  $k$  dependence,<sup>31</sup> the peak becomes broader comparing to the magnetic-field induced peak splitting (see below). (iii) The amplitude of the peak splitting depends on the genuine barrier amplitude  $\hat{V}_0$  as well as the exchange amplitude  $\hat{U}_B$ . For example, the splitted peaks merge into a single peak at the tunneling limit ( $\hat{V}_0 \rightarrow \infty$ ) even if  $\hat{U}_B$  is kept constant. (iv) The current corresponding to the ZBCP is carried by the Cooper pair in the superconductor as described in the previous subsection. This

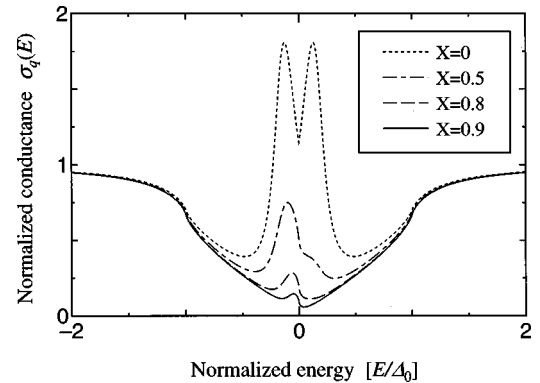


FIG. 9. The normalized conductance spectra  $\sigma_q(E)$  as a function of  $X$  for  $Z_{0,\uparrow}=2.5$ ,  $Z_{0,\downarrow}=7.5$ , and  $\beta=\pi/4$ . As  $X$  becomes larger, the higher energy peak becomes smaller. Thus, the polarization can be estimated from the ratio of two peak heights. When  $X$  approaches 1, the two peaks disappear.

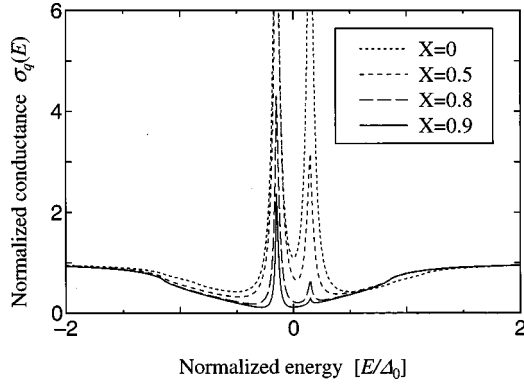


FIG. 10. The normalized conductance spectra  $\sigma_q(E)$  in an applied magnetic field ( $\mu_g H/\Delta_0=0.15$ ) as a function of  $X$  for  $Z_{0,\uparrow}=5$ ,  $Z_{0,\downarrow}=5$ , and  $\beta=\pi/4$ . As  $X$  becomes larger, the peak with higher energy is largely reduced. When  $X$  approaches 1, the two peaks disappear.

corresponds to the fact that the AR process is the second-lowest order tunneling process which requires both up and down spins tunneling. Hence, as  $Z_{0,\downarrow}$  becomes larger and as the tunneling probability for down spins are suppressed, the conductance peaks and the AR process are rapidly reduced even if  $Z_{0,\uparrow}$  is kept zero. (v) The spin current is increased as  $\tilde{U}_B$  is raised from zero even if  $X$  in the ferromagnet is kept at zero (unpolarized). This feature directly corresponds to the spin-filtering effect that the spin-selective tunneling occurs due to the presence of the exchange field in the insulator.

Next, various types of the ZBCP splitting expected for  $d$ -wave superconductors and their polarization effects are analyzed. Mainly two possibilities other than the ferromagnetic insulator effects have been proposed for the origins of the ZBCP splitting on high- $T_c$  superconductor junctions. One is the Zeeman effect due to an applied magnetic field and the other is the inducement of the BTRS states such as a  $d_{x^2-y^2}+is$  wave. The conductance spectra in an applied magnetic field is calculated from above formula by simply using the relation

$$\sigma_{q[s]}(E) = \sigma_{q[s],\uparrow}(E - \mu_B H) + [-] \sigma_{q[s],\downarrow}(E + \mu_B H), \quad (17)$$

where  $\mu_B H$  is the Zeeman energy. Calculated charge conductance spectra for a  $d_{x^2-y^2}$  wave superconductor as a function of  $X$  are shown in Fig. 10. The amplitude of the splitting is linear to the applied field independent of the barrier heights. Moreover, since the energy shift induced by the magnetic field does not have  $k$  dependence, the broadening of the peaks are not present. The ratio of the splitted peak heights simply reflects the polarization in the ferromagnet, which is consistent with the results by Tedrow and Meservey.<sup>1,32</sup> On the other hand,  $\sigma_q(E)$  for a  $d_{x^2-y^2}+is$  wave superconductor is calculated by setting  $\Delta_{\pm} = \Delta_0 \cos 2(\theta_{\pm} \mp \beta) + i\Delta_s$ . Calculated charge conductance spectra for various  $X$  values are shown in Fig. 11. The amplitude of the splitting is almost equivalent to the amplitude of the  $s$ -wave component. The shape of the spectrum without the polarization ( $X=0$ ) is quite similar to that shown in Fig. 10 ( $X=0$ ). As  $X$  becomes larger, the heights of the two peaks are reduced, which is consistent with that shown in Fig. 3.

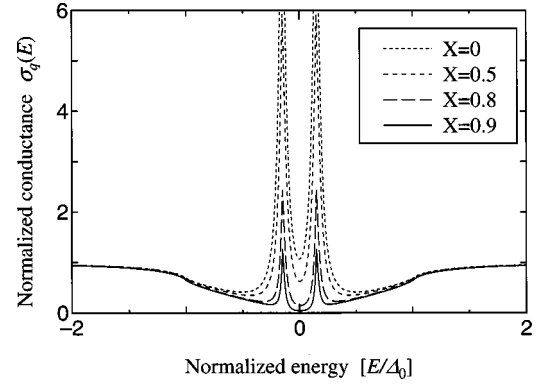


FIG. 11. The normalized conductance spectra  $\sigma_q(E)$  for the BTRS states ( $\Delta_s/\Delta_0=0.15$ ) as a function of  $X$  for  $Z_{0,\uparrow}=5$ ,  $Z_{0,\downarrow}=5$ , and  $\beta=\pi/4$ . As  $X$  becomes larger, the heights of the two peaks are reduced symmetrically.

On the other hand, differently from Figs. 9 and 10, since the peak splitting is not induced by spin-dependent effects in this case, the polarization in the ferromagnets does not yield an imbalance in the peak heights. Thus the heights of the two peaks are reduced symmetrically.

The responses of the ZBCP on the variation of the polarization and the applied magnetic field are summarized as follows. (i) The peak splitting due to the ferromagnetic insulator and the Zeeman effect are spin dependent. Therefore, the polarization in the ferromagnet induces the asymmetrical splitting of the ZBCP. (ii) The amplitude of the peak splitting is linear to the applied field in the case of the Zeeman effect. However, it is nonlinear in the cases of the ferromagnetic insulator effects<sup>7</sup> and the BTRS states.<sup>23</sup> In particular, the peak splittings are expected even in the absence of the applied field for these two cases. (iii) The multiplication of the BTRS states with the Zeeman effect induces an additional peak splitting, that is, the ZBCP splits into four peaks. However, the multiplication of the Zeeman with the ferromagnetic insulator effects does not yield the additional peak splittings.

The experimental observations of the ZBCP splitting have been reported for normal metal–high- $T_c$  superconductor junctions.<sup>14,24,33,34</sup> It is a really interesting experiment to observe the same features by using ferromagnets–high- $T_c$  superconductor junctions in order to distinguish the spin-dependent effects from the BTRS states inducement. Recently, Sawa *et al.* have detected an asymmetric magnetic field response in  $\text{La}_{0.67}\text{Sr}_{0.33}\text{MnO}_3\text{-YBa}_2\text{Cu}_3\text{O}_{7-\delta}$  junctions.<sup>35</sup> The qualitative features on the magnetic field responses of their junctions are consistent with F/FI/S with the  $d_{x^2-y^2}$  wave explained here. A detailed comparison between the above formulas and their experiments is expected.

Finally, a simple proposal is given for a possible device application utilizing the ferromagnetic insulator effects. The thickness of the insulator is the order of 1 nm in usual tunneling junctions. Since the controlling of properties in such a thin layer requires high technology, as far as we know, not so many experimental trials have been accomplished thus far. However, as shown in this paper, a small change in the insulator property causes a drastic change on the transport properties. Therefore, the controlling of the barrier properties is one of the most promising methods to create new func-

tional devices. For example, consider an F/FI/S junction with a  $d_{x^2-y^2}$ -wave superconductor ( $\beta = \pi/4$ ). The sharp ZBCP is drastically modified as the difference between  $Z_{\uparrow}$  and  $Z_{\downarrow}$  becomes larger as shown in Fig. 7. This means that, for a fixed bias voltage, a large response in current is expected due to a small variation in the exchange interaction in the insulator. This response is applicable for the high-sensitive magnetization measurement of a thin insulating film by inserting the film into a junction as a tunneling barrier. If the exchange interaction is sensitive to the external field, this effect can be used as a magnetic sensor. Alternatively, if the magnetization of the insulator shows hysteresis on the external field variation, a memory function can be realized. The current gain of the junction as a function of the external field is largely enhanced by using a superconductor–ferromagnetic-insulator–superconductor junction with a  $d$  wave, because negative-conductance regions are expected just beside the ZBCP in this configuration.<sup>36,37</sup> Differently from conventional superconducting memories based on a flux-quantum logic, a large-scale integration circuit may be possible based on the present principle.

#### IV. SUMMARY

In this paper, the conductance spectra for the charge and the spin currents under the influence of the exchange interaction have been calculated based on the scattering method. The influence of the spin polarization on the transport properties has been clarified. It is shown that the retroreflectivity

of the standard Andreev reflection process is broken in the presence of an exchange field and that the surface bound states due to superconducting pair potentials do not contribute to the spin current. Next, the ferromagnetic insulator including the spin-filtering effect are analyzed. It is shown that the spin polarization gives asymmetric peak splitting. Moreover, various features in the splitting of ZBCP due to the ferromagnetic insulator, the Zeeman splitting, and the BTRS states effects are analyzed in detail. It is shown that the spin-polarized tunneling gives important information for identifying the origin of the ZBCP splitting. By comparing the present analysis with experimental data, we expect that the mechanism of the peak splitting in high- $T_c$  superconductors will be well identified. In the present model, we have neglected the effects of spin-orbit scattering<sup>38</sup> and the nonequilibrium properties of superconductors.<sup>39,40</sup> Inclusion of these effects would be necessary for a complete theory. The formulation for triplet superconductors will be presented in another publication.<sup>41</sup>

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