Transport properties in ferromagnetic Josephson junctions between triplet superconductors

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Abstract
Charge and spin Josephson currents in a ballistic superconductor–ferromagnet–superconductor junction with spin-triplet pairing symmetry are studied using the quasiclassical Eilenberger equation. The gap vector of superconductors has an arbitrary relative angle with respect to magnetization of the ferromagnetic layer. We clarify the effects of the thickness of the ferromagnetic layer and the magnitude of the magnetization on the Josephson charge and spin currents. We find that the 0–π transition can occur when the misorientation angle between the exchange field of the ferromagnetic layer and the d-vector is smaller than π/4. We also show how spin current flows due to misorientation between the exchange field and the d-vector.

(Some figures in this article are in colour only in the electronic version)

1. Introduction

Since both superconductivity and ferromagnetism are antagonistic ordered phases of matter, for a few decades after the discovery of superconductivity the interplay between superconductivity and ferromagnetism was not a subject of intensive research interest. However, recently the study of the interplay between superconductor and ferromagnet has been revived and, in particular, the proximity effect and Andreev reflection in superconductor–ferromagnet junctions has attracted much attention [1–5].

When a singlet superconductor is placed in proximity to a ferromagnetic layer, one can find triplet pairing correlations in the ferromagnetic layer and this component of pairing provides interesting effects [6–12]. On the other hand, spin-triplet superconductors have been extensively investigated due to their anomalous features [13–35]. The proximity of a spin-triplet superconductor and a ferromagnet is more interesting than the singlet case in the sense that both the spin-triplet superconductor and ferromagnet have a magnetic nature.

In [36–38] the authors demonstrated the presence of both charge and spin currents in the systems consisting of a thin barrier of ferromagnet sandwiched between two triplet superconductors. They obtained a spin current due to the coupling of the ferromagnetic moment to the spin of the triplet Cooper pairs. Also, they found a 0–π transition in the triplet superconductor–ferromagnet–superconductor (SFS) Josephson junction generated by the misalignment of the d-vectors of the triplet superconductors with respect to the magnetic moment of the ferromagnet. In contrast to [36–38], where the ferromagnet is modeled as a δ function potential, in this paper we allow for arbitrary length of ferromagnet, which is a more realistic modeling, based on the quasiclassical Eilenberger equation.

In this paper, we investigate a triplet SFS structure with arbitrary misalignment between magnetization of the ferromagnet and gap vector of superconductors (figure 1), by changing the exchange field and thickness of the ferromagnetic layer. We find that the 0–π transition can occur when the misorientation angle between the exchange field of the ferromagnetic layer and d-vector is smaller than π/4.

The organization of this paper is as follows. In section 2 the quasiclassical equations for Green’s functions are presented. In section 3, we calculate Green’s functions of the system analytically, presenting the formulas for the Green’s functions to calculate the charge and spin current densities. In section 4, we show the results of spin and charge currents. The paper will be finished with the conclusions in section 5.

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2. Formalism and basic equations

We consider a clean SFS structure with a homogeneous ferromagnet of thickness \( l \) and bulk triplet superconductors (see figure 1). The ferromagnetic layer is characterized by an exchange field \( h \). There is a misorientation angle \( \alpha \) between the exchange field of the ferromagnet and order parameters of superconductors (\( \vec{d} \)-vectors). Interfaces are fully transparent and magnetically inactive. The thickness \( l \) is larger than the Fermi wavelength and smaller than the elastic mean free path. Then, we can use a quasiclassical description in the clean limit, and apply the Eilenberger equation in this limit [39–41]:

\[
\hbar v_F \nabla \tilde{g} + \left[ \varepsilon_m \tilde{\sigma}_3 + i(\tilde{\Delta} + \tilde{\mathbf{h}}), \tilde{g} \right] = 0, \tag{1}
\]

with the normalization condition \( \tilde{g}^\dagger \tilde{g} = \mathbb{I} \). Here, \( h \) is Planck’s constant, \( \varepsilon_m = \pi T (2m + 1) \) are discrete Matsubara energies, \( T \) is the temperature, \( v_F \) is the Fermi velocity and \( \tilde{\sigma}_3 = \tilde{\tau}_3 \otimes \mathbb{I} \) in which \( \tilde{\tau}_3 \) is the third Pauli matrix in particle–hole space. \( \tilde{\sigma}_j \) \((j = 1, 2, 3)\) denote Pauli matrices in spin space in the following. The Matsubara propagator \( \tilde{g} \) can be written in the standard form:

\[
\tilde{g} = \begin{pmatrix}
\tilde{g}_1 + \tilde{g}_1 \cdot \tilde{\sigma} & (\tilde{g}_2 + \tilde{g}_2 \cdot \tilde{\sigma}) i \tilde{\sigma}_2 \\
(\tilde{g}_2 \cdot \tilde{\sigma}) (\tilde{g}_3 + \tilde{g}_1 \cdot \tilde{\sigma}) & \tilde{g}_4 - \tilde{\sigma}_2 (\tilde{g}_1 \cdot \tilde{\sigma}) \tilde{\sigma}_2 \\
\end{pmatrix} \tag{2}
\]

where \( \tilde{\sigma} = (\tilde{\sigma}_1, \tilde{\sigma}_2, \tilde{\sigma}_3) \) is the vector of Pauli matrices in spin space. Also, we can write \( \tilde{\mathbf{h}} \) as follows:

\[
\tilde{\mathbf{h}} = \begin{pmatrix}
\mathbf{h} \cdot \tilde{\sigma} \\
0 \\
(\mathbf{h} \cdot \tilde{\sigma})^\ast \\
\end{pmatrix}. \tag{3}
\]

We use the Eilenberger equation for both ferromagnetic and superconducting materials. For superconductors we set \( \mathbf{h} = 0 \). The matrix structure of the off-diagonal self-energy \( \tilde{\Delta} \) in the Nambu space is

\[
\tilde{\Delta} = \begin{pmatrix}
0 & (\mathbf{d} \cdot \tilde{\sigma}) i \tilde{\sigma}_2 \\
(\tilde{\sigma}_2 (\mathbf{d}^\ast \cdot \tilde{\sigma}) i \tilde{\sigma}_2 & 0 \\
\end{pmatrix}. \tag{4}
\]

In the ferromagnet we set \( \tilde{\Delta} = 0 \) and \( \mathbf{h} = (0, \sin \alpha, \cos \alpha) \). In this paper, the unitary state, \( \mathbf{d} \times \mathbf{d}^\ast = 0 \), is investigated.

Solutions of equation (1) have to satisfy the conditions for Green’s functions in the bulks of the superconductors:

\[
\tilde{g} (\pm \infty) = \begin{pmatrix}
\frac{\varepsilon_m \tilde{\sigma}_3 + i \tilde{\Delta}_{1,2}}{\sqrt{\varepsilon_m^2 + |\tilde{d}_{1,2}|^2}} \\
\end{pmatrix}, \tag{5}
\]

\[
\tilde{d}_{1,2} = \tilde{d}_{L,R} (\tilde{v}_F) \exp \left( \pm \frac{i \phi}{2} \right), \tag{6}
\]

where \( \phi \) is the external phase difference between the order parameters of the bulks of superconductors. Equation (1) has to be supplemented by the continuity conditions for Green’s functions at the interfaces between ferromagnet and superconductors. For all quasiparticle trajectories, the Green’s functions satisfy the boundary conditions both in the right and left bulks as well as at the interfaces.

3. Analytical results of Green’s functions

In principle, Eilenberger equations (1) should be supplemented with a self-consistent equation for the gap vector. While these numerical self-consistent calculations can be used to investigate the spatial variation of gap vector in the superconductors and also the dependence of gap vector on temperature, in this paper we do not apply them. In our analytical calculations, we consider a step-like function of gap vector which is zero in the ferromagnet and has a finite value in the superconducting part. It has been shown that the absolute value of a self-consistently determined order parameter decreases near the interface between superconductor and normal metal, while its dependence on the direction of Fermi velocity remains unaltered [42]. In [42–45] a qualitative agreement between self-consistent and non-self-consistent calculations for unconventional Josephson junctions is found. Also, it is observed that the results of non-self-consistent calculation in [46] are coincident with the experimental data in [47]. Also, the results by non-self-consistent treatment in [48] are similar to the experimental results in [49]. Consequently, we believe that the non-self-consistent approximation does not change results qualitatively. Therefore, in our calculations, a simple model of the constant order parameter up to the interface is considered to obtain qualitative results. When the Fermi velocity in the superconductor is much larger than that in the ferromagnet, the so-called rigid boundary condition, namely non-self-consistent treatment, is justified. Using such an approximation, the analytical expressions for the charge and spin currents in the ferromagnet and superconductors can be obtained for a specified form of the order parameter. The expressions for the charge and spin currents are the following [45, 50]:

\[
\mathbf{j}_c (\mathbf{r}) = 2i \pi e \hbar B T N (0) \sum_m \{ \mathbf{v}_F g_1 (\tilde{v}_F, \mathbf{r}, \varepsilon_m) \}, \tag{7}
\]

\[
\mathbf{j}_s (\mathbf{r}) = i \pi \hbar B T N (0) \sum_m \{ \mathbf{v}_F (\tilde{\mathbf{e}}, \mathbf{g}; (\tilde{v}_F, \mathbf{r}, \varepsilon_m)) \}. \tag{8}
\]
where \( k_B \) is the Boltzmann constant, \( \hat{e}_z = (\hat{x}, \hat{y}, \hat{z}) \) and \( N(0) \) is the density of states in the normal state at the Fermi energy. We consider the order parameter as follows:

\[
d_{a}(T, \hat{v}_F, \mathbf{r}) = d_{a}(T, \hat{\nu}, \mathbf{r})\hat{z}. \tag{9}
\]

Solving the Eilenberger equation in the ferromagnetic region and the superconductors under the continuity of solutions across the interfaces, \( x = 0, l \) and the boundary conditions at the bulks, we obtain the Green’s functions in the ferromagnet. The Green’s functions are given by

\[
g_{1F} = \frac{\eta(p^2 - 1)}{D} \left( p^2 + 2p[\beta + \sin^2 \alpha(2\beta^2 - \beta - 1)] + 1 \right), \tag{10}
\]

and

\[
g_{3F} = \frac{\eta p}{D} \left( A \cos \left( \frac{2\beta x}{v} \right) + B \sin \left( \frac{2\beta x}{v} \right) \right) \left[ \cos \alpha \left[ A \sin \left( \frac{2\beta x}{v} \right) - B \cos \left( \frac{2\beta x}{v} \right) \right] + C \sin \alpha \left[ A \cos \left( \frac{2\beta x}{v} \right) - B \sin \left( \frac{2\beta x}{v} \right) \right] + C \cos \alpha \right], \tag{11}
\]

where

\[
A = \sin 2\alpha [\beta - 1] [p^2(2\beta + 1) - 2p(\beta \sin^2 \alpha + \cos^2 \alpha) - 1],
\]

\[
B = \sin 2\alpha \sin(2hl/\nu) [p^2(2\beta - 1) + 2p(\beta \sin^2 \alpha + \cos^2 \alpha) + 1],
\]

\[
C = \cos^2 \alpha \sin(2hl/\nu) [p^2 + 2p(\sin^2 \alpha - \beta - 1) + 1],
\]

\[
D = p^2 + 2p[\beta(\sin^2 \alpha + \cos^2 \alpha)^2 + 2\beta^2 - 1] + 4p\beta(\beta^2 - 1) \sin^2 \alpha + \cos^2 \alpha + 1,
\]

\[
\beta = \cos(2hl/\nu) \eta = \text{sgn}(v_\alpha)
\]

\[
p = \frac{d_1d_2^* \exp(-2\nu_{nl}/\nu)}{(\varepsilon_m + \eta\Omega_1)(\varepsilon_m + \eta\Omega_2)} \Omega_{1,2} = \sqrt{\varepsilon_m^2 + [d_{1,2}]^2}.
\]

For parallel alignment of gap vector and exchange field, \( \alpha = 0 \), Green’s functions reduce to

\[
g_{1F} = \frac{\eta(p^2 - 1)}{p^2 + 2p\beta + 1}, \quad g_{3F} = \frac{2\eta p \sin(2hl/\nu)}{p^2 + 2p\beta + 1}. \tag{12}
\]

It is seen from this expression that Green’s functions and hence charge and spin currents depend on the exchange field through \( \beta \).

Also, for the perpendicular configuration \( \alpha = \frac{\pi}{2} \), we obtain

\[
g_{1F} = \frac{\eta(p - 1)}{p + 1}, \quad g_{3F} = 0 \tag{13}
\]

where the Green’s functions are independent of the exchange field. As a result, we do not see any 0–\( \pi \) transition in this case [51].

4. Numerical results of charge and spin currents

Here, we have considered \( p \)– wave superconductors with the model of order parameter as

\[
d(T, \hat{v}_F, \mathbf{r}) = \Delta(T) \left( k_x + ik_y \right) \tag{14}
\]

which would be realized in Sr$_2$RuO$_4$ [16]. Note that for the other component of the order parameter \( d(T, \hat{v}_F, \mathbf{r}) = \Delta(T)(k_x - ik_y) \) [16], the results remain the same since the order parameter is included in the Green’s functions only in the form \( p = \frac{d_1d_2^* \exp(-2\nu_{nl}/\nu)}{(\varepsilon_m + \eta\Omega_1)(\varepsilon_m + \eta\Omega_2)} \) and \( p \) is the same for both choices of order parameters. The function \( \Delta(T) \) describes the dependence of the order parameter \( d \) on the temperature \( T \). We have used Green’s functions in equations (10) and (11) and calculated charge and spin currents for our specified model of the order parameter. For representative values of \( h, l \) and \( \alpha \), we have shown charge and spin currents in figures 2–5 as a function of phase difference between superconductors. In the figures, \( \xi(T) = \frac{h}{\Delta(T)} \) denotes a superconductor coherence length at the bulk of the superconductor.

In figures 2 and 3 we have plotted the charge current as a function of phase difference for different misorientation between the \( d \)-vector and the exchange field, different thicknesses of the ferromagnet and different values of the exchange field \( h \). It is seen that, except for the case of \( \alpha = \pi/2 \), by increasing \( h \) or \( l \), the Josephson charge current decreases. This is reasonable because the ferromagnet suppresses the proximity effect and hence the Josephson current. Also, in figures 2 and 3, for the case of \( \alpha = \frac{\pi}{2} \) it is obtained that the charge current is independent of the exchange field of the ferromagnet but depends on the thickness of the ferromagnet, which is understandable from the analytical expressions of the current. By increasing the thickness of the ferromagnet, the charge current decreases. In addition, for small misorientation angle \( \alpha \) in figures 2 and 3, it is shown that, by increasing \( h \) or \( l \), 0–\( \pi \) transition occurs.

While charge current is independent of position in the ferromagnet, spin current is dependent on position and is not conserved. In fact, spin current at the left and right interfaces have opposite sign, and in the middle of the ferromagnetic layer, the spin current vanishes. We have plotted three polarization components of spin current at the left interface between superconductor and ferromagnetic layer in figures 4, 5 and 8. Note that in [36–38], spin current is evaluated in the bulk superconductor, and hence spin current is polarized along the \( x \)-direction.

In our formalism, there are in general three components \( (S_x, S_y, S_z) \) of spin current due to the presence of the ferromagnetic layer with finite thickness. However, we find that spin current with spin polarization parallel to the exchange field is absent. As seen in figures 4 and 5 for \( \alpha = \pi/4 \), there are \( S_y \) and \( S_z \) components with the same magnitude but opposite sign, such that spin current with spin polarization parallel to the exchange field is vanishing. We also find that all components of spin current vanish at \( \alpha = 0 \) and \( \pi/2 \). As shown in figure 4, spin current is strongly dependent on the exchange field, so for a weak ferromagnet (for example \( h = 0.5\Delta \)), spin current is negligibly small.

For some values of the phase difference, spin current exists in the absence of charge current, as seen in figures 4 and 5. Also, there is a local maximum or minimum in spin current in \( \phi = 0 \) or \( \pi \). This can be explained as follows. Charge currents are odd functions but spin currents are even functions of the phase difference, because of the symmetry under the time reversal. Therefore, the slope of the spin current becomes zero at \( \phi = 0 \) and \( \pi \).
Figure 2. The $x$ component of the charge current versus the phase difference $\phi$ for $l = \xi/2$, $T = 0.05T_C$, and different $h$ and $\alpha$. Currents are given in units of $j_{0,e} = \frac{eN(0)v_F}{\Delta(0)}$.

Figure 3. The $x$ component of the charge current versus the phase difference $\phi$ for $h = 5\Delta$, $T = 0.05T_C$, and different $l$ and $\alpha$. 

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Figure 4. The normal component (x component) of the spin current at the left interface, versus the phase difference $\phi$, for $l = \xi/2$, $T = 0.05T_C$, different $h$ and $\alpha = \pi/4$. Currents are calculated in units of $j_{0x} = \frac{\pi}{2} \hbar N(0)v_F/\Delta_1(0)$.

Figure 5. The $x$ component of the spin current at the left interface, versus the phase difference $\phi$ for $h = 5\Delta$, $T = 0.05T_C$, $\alpha = \pi/4$ and different $l$. 
Figure 6. The $x$ component of the critical charge current versus $h$ for $T = 0.05T_c$ and different $l$ and different $\alpha$.

Figure 7. The $x$ component of the critical charge current versus the thickness of the ferromagnetic layer at $T = 0.05T_c$ for different $h$ and different $\alpha$. 
Figure 8. The $x$ component of the critical charge and spin currents versus $\alpha$ for $T = 0.05T_C$. Left panels are for $l = \xi$ and right panels are for $l = 2\xi$.

For charge current, we see that by increasing the thickness of the ferromagnetic layer the current decreases with oscillation because of decreasing of quantum coherency, but for the spin current the situation is quite different. As a result of the exchange field, a thicker layer has a greater effect on the spin of quasiparticles and thus a larger spin current because the spin current flowing in the junction is due to the presence of the ferromagnet.

In contrast to currents for $\alpha = 0$ which are sensitive to the exchange field, for $\alpha = \pi/2$, the results are independent of the exchange field because in this case the spin polarization and exchange field of the ferromagnetic layer are parallel. Then, we can show that Josephson currents can be divided into spin-up and spin-down contributions, and Josephson currents through the spin-up (spin-down) pairing channel are identical to those in a singlet superconductor–normal metal–singlet superconductor junction, where the normal metal has a slightly larger (smaller) Fermi wavevector than in the superconductor [51]. Therefore, the currents for $\alpha = \pi/2$ are independent of the exchange field. Thus, charge current in this case is only dependent on the thickness of the ferromagnetic layer and spin current is absent.

By changing the parameters, we can find a tendency toward the $0-\pi$ transition at $\alpha = 0$. The critical currents as a function of $h$ and $l$ are shown in figures 6 and 7 for different misorientations $\alpha = 0, \pi/4$ and $\pi/2$. Here, the critical charge and spin currents are defined as $j_{ce} = \text{Max}_\phi [j_c(\phi) = j_s(\phi^*)]$ and $j_s = j_s(\phi^*)$, respectively. The $0-\pi$ transition can be seen for $0 \leq \alpha < \pi/4$ for large $h$ and $l$. Note that the mechanism of the $0-\pi$ transition predicted in this paper is the same as that in [36–38]. In fact, spin-dependent phase shifts are naturally included in our theory.

Critical charge and spin currents as a function of $\alpha$ are plotted in figure 8. In [38], it is shown that $0-\pi$ transition can occur for $0 \leq \alpha \leq \pi/4$ for a $\delta$ function type ferromagnetic barrier. We find that this is the case even in the presence of a ferromagnetic layer with finite thickness instead of a $\delta$ function type barrier. As shown in figure 8, for finite thickness of ferromagnetic layer, the $0-\pi$ transition can be found by changing the misorientation angle. The $0-\pi$ transition is accompanied by the jump of the spin current due to the phase jump at the $0-\pi$ transition point [12]. Even when charge current shows a monotonous dependence on $\alpha$, spin current can show non-monotonous $\alpha$ dependence. By increasing the temperature of the system, the $0-\pi$ transition occurs for thinner ferromagnetic layers (see figure 9).
5. Conclusions

In this paper, we have investigated transport properties in triplet SFS Josephson junctions and clarified the effects of thickness of ferromagnetic exchange field and the misorientation angle between the exchange field and d-vector of superconductors on the Josephson charge and spin currents. We found that the 0–π transition can occur when the misorientation angle between exchange field and d-vector is smaller than π/4. This is consistent with the previous works based on tunneling Hamiltonian and δ-function tunneling barrier analyses [36–38]. When exchange field and d-vector are perpendicular to each other, Josephson currents are independent of the exchange field. Also, we showed how spin current flows due to misorientation between the exchange field and d-vector and found that it is absent when two vectors are parallel or perpendicular to each other.

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References

[1] Buzdin A I 2005 Rev. Mod. Phys. 77 935
[2] Bergeret F S, Volkov A F and Efetov K B 2005 Rev. Mod. Phys. 77 1321
[3] Volkov A F and Efetov K B 2008 Phys. Rev. B 78 024519
[4] Zareyan M, Belzig W and Nazarov Yu V 2002 Phys. Rev. B 65 184505
[5] Braude V and Blanter Ya M 2008 Phys. Rev. Lett. 100 207001
[6] Bergeret F S, Volkov A F and Efetov K B 2001 Phys. Rev. Lett. 86 4096
[7] Kadigrobov A, Shechter R I and Jonson M 2001 Europhys. Lett. 50 394
[8] Eschrig M, Löfwander T, Champel Th, Cuevas J C and Schön G 2007 J. Low Temp. Phys. 147 457
[9] Eschrig M and Löfwander T 2008 Nature Phys. 4 138
[10] Yokoyama T, Tanaka Y and Golubov A A 2007 Phys. Rev. B 75 134510
[11] Linder J, Yokoyama T, Sudbø A and Eschrig M 2009 Phys. Rev. Lett. 102 107008
[12] Alidoust M, Linder J, Rashidi G, Yokoyama T and Sudbø A 2010 Phys. Rev. B 81 214504
[13] Maeno Y, Hashimoto H, Yoshida K, Nishizaki S, Fujita T, Bednorz J G and Lichtenberg F 1994 Nature 372 533
[14] Ishida K, Mukuda H, Kitaoka Y, Asayama K, Mao Z Q, Morii Y and Maeno Y 1998 Nature 396 658
[15] Luke G M et al 1998 Nature 394 558
[16] Mackenzie A F and Maeno Y 2003 Rev. Mod. Phys. 75 657
[17] Nelson K D, Mao Z Q, Maeno Y and Liu Y 2004 Science 306 1151
[18] Asano Y, Tanaka Y, Sigrist M and Kashiwaya S 2003 Phys. Rev. B 67 184505
[19] Tou H, Kitaoka Y, Ishida K, Asayama K, Kimura Y, Yamamoto E, Haga Y and Maezawa K 1998 Phys. Rev. Lett. 80 3129
[20] Muller V, Roth Ch, Maurer D, Scheidt E W, Lers K, Bucher E and Bömel H E 1987 Phys. Rev. Lett. 58 1224
[21] Qian Y J, Xu M F, Schenstrom A, Baum H P, Ketterson J B, Hinks D, Levy M and Sarma B K 1987 Solid State Commun. 63 599
[22] Abrikosov A A 1983 J. Low Temp. Phys. 53 359
[23] Fukuyama H and Hasegawa Y 1987 J. Phys. Soc. Japan 56 877
[24] Lebed A G, Machida K and Otsuki M 2000 Phys. Rev. B 62 795
[25] Saxena S S et al 2000 Nature 406 587
[26] Pfeiferer C, Uhlarz M, Hayden S M, Vollmer R, Lohneysen H V, Bernhoeft N R and Lonzarich G G 2001 Nature 412 58
[27] Aoki D, Huxley A, Ressouche E, Braithwaite D, Flouquet J, Brison J, Lhotel E and Paulsen C 2001 Nature 413 613
[28] Bolech C J and Giamarchi T 2005 Phys. Rev. Lett. 97 147002
[29] Bolech C J and Giamarchi T 2004 Phys. Rev. B 70 054518
[30] Asano Y 2005 Phys. Rev. B 72 092508
[31] Asano Y 2006 Phys. Rev. B 74 220501(R)
[32] Mahmoodi R, Kolesnichenko Yu A and Shevchenko S N 2002 Fiz. Nizk. Temp. 28 262
[33] Mahmoodi R, Kolesnichenko Yu A and Shevchenko S N 2002 Low Temp. Phys. 28 184 (Engl. Transl.)
[34] Rashedi G and Kolesnichenko Yu A 2007 Physica C 451 31
[35] Rashedi G and Kolesnichenko Yu A 2005 Supercond. Sci. Technol. 18 482
[36] Rashedi G, Rahnavard Y and Kolesnichenko Yu A 2010 Fiz. Nizk. Temp. 36 262
[37] Rashedi G, Rahnavard Y and Kolesnichenko Yu A 2010 Low Temp. Phys. 36 205 (Engl. Transl.)
[38] Rashedi G and Kolesnichenko Yu A 2005 Fiz. Nizk. Temp. 31 634
[39] Rashedi G and Kolesnichenko Yu A 2005 Low Temp. Phys. 31 481 (Engl. Transl.)
[40] Rahnavard Y, Rashedi G and Yokoyama T 2010 J. Phys.: Condens. Matter 22 415701
[36] Kastening B, Morr D K, Manske D and Bennemann K 2006
  *Phys. Rev. Lett.* **96** 047009

[37] Brydon P M R, Kastening B, Morr D K and Manske D 2008
  *Phys. Rev. B* **77** 104504

[38] Brydon P M R and Manske D 2009 *Phys. Rev. Lett.* **103** 147001

[39] Eilenberger G 1968 *Z. Phys.* **214** 195

[40] Kulik I O and Omelyanchuk A N 1978 *Fiz. Nizk. Temp.* **4** 296

[41] Kupriyanov M Yu 1981 *Fiz. Nizk. Temp.* **7** 700

[42] Barash Yu S, Bobkov A M and Fogelström M 2001 *Phys. Rev. B* **64** 214503

[43] Amin M H S, Coury M, Rashkeev S N, Omelyanchouk A N and Zagoskin A M 2002 *Physica B* **318** 162

[44] Viljas J K and Thuneberg E V 2002 *Phys. Rev. B* **65** 64530

[45] Viljas J 2000 arXiv:cond-mat/0004246

[46] Faraii Z and Zareyan M 2004 *Phys. Rev. B* **69** 014508

[47] Freamat M and Ng K-W 2003 *Phys. Rev. B* **68** 060507

[48] Yip S-K 1999 *Phys. Rev. Lett.* **83** 3864

[49] Backhaus S, Pereverzev S, Simmonds R W, Loshak A, Davis J C and Packard R E 1998 *Nature* **392** 687

[50] Serene J and Rainer D 1983 *Phys. Rep.* **101** 221

[51] Powell B J, Annett J F and Győrffy B L 2003 *J. Phys. A: Math. Gen.* **36** 9289