Point-contact spectroscopy in Co-doped CaFe$_2$As$_2$: nodal superconductivity and topological Fermi surface transition

R S Gonnelli$^1$, M Tortello$^1$, D Daghero$^1$, R K Kremer$^2$, Z Bukowski$^{3,4}$, N D Zhigadlo$^3$ and J Karpinski$^3$

$^1$Dipartimento di Scienza Applicata e Tecnologia, Politecnico di Torino, corso Duca degli Abruzzi 24, 10129 Torino (TO), Italy
$^2$Max-Planck-Institut für Festkörperforschung, D-70569 Stuttgart, Germany
$^3$Laboratory for Solid State Physics, ETHZ, CH-8093 Zurich, Switzerland
$^4$Institute of Low Temperature and Structure Research, Polish Academy of Sciences, PO Box 1410, 50-422 Wroclaw, Poland

E-mail: renato.gonnelli@polito.it

Received 22 January 2012, in final form 6 March 2012
Published 12 April 2012
Online at stacks.iop.org/SUST/25/065007

Abstract
We performed point-contact Andreev-reflection spectroscopy measurements in Ca(Fe$_{1-x}$Co$_x$)$_2$As$_2$ single crystals with effective $x = 0.060 \pm 0.005$. The spectra of ab-plane contacts show a zero-bias maximum and broad shoulders at about 5–6 meV. Their fit with the three-dimensional Blonder–Tinkham–Klapwijk (BTK) model (making use of an analytical expression for the Fermi surface that mimics the one calculated from first principles) shows that this compound presents a large isotropic gap on the quasi-2D electronlike Fermi surface sheets and a smaller anisotropic (possibly nodal) gap on the 3D holelike Fermi surface pockets centered at the Z point in the Brillouin zone. These results nicely fit into the theoretical picture for the appearance of nodal superconductivity in 122 compounds.

(Some figures may appear in colour only in the online journal)

1. Introduction
Among the members of the 122 family of Fe-based superconductors, the CaFe$_2$As$_2$ system is particularly challenging from both the theoretical and the experimental points of view. With respect to Sr-122 and Ba-122, which have been widely studied in the past and whose phase diagrams have been already determined, the Ca-122 system presents some peculiarities (namely, a high sensitivity of the cell structure to chemical and physical pressure, and a strong dependence of the physical properties of single crystals on the growth procedure [1]) that have hindered its investigation and caused conflict between data and calculations by different groups.

Recently, the phase diagram of Ca(Fe$_{1-x}$Co$_x$)$_2$As$_2$ as a function of $x$ has been determined independently by Hu et al [1] and Harnagea et al [2]. Despite some small quantitative disagreement on the Co content, the two phase diagrams agree well with each other. In particular, unlike in the sister compounds Sr(Fe$_{1-x}$Co$_x$)$_2$As$_2$ and Ba(Fe$_{1-x}$Co$_x$)$_2$As$_2$, the superconducting phase appears abruptly (at $x = 0.03$ in [1]) where the high-temperature magnetic and structural transitions are still present, and the critical temperature gradually decreases with increasing $x$.

In this paper we present the results of point-contact Andreev-reflection spectroscopy (PCARS) measurements in Ca(Fe$_{1-x}$Co$_x$)$_2$As$_2$ single crystals with $x = 0.060 \pm 0.005$, aimed at determining the number, the amplitude and the symmetry of the superconducting order parameter(s) (OP) in this compound. These pieces of information, which were lacking up to now, are of particular interest if compared to the predictions concerning the dependence of the gap symmetry of 122 compounds on the shape of the Fermi surface (FS) (and thus on the height of the pnictogen above the Fe layer,
Figure 1. Some experimental conductance curves (not normalized) as a function of temperature. The point-contact resistance is about 3 Ω. Inset: the zero-field-cooling DC magnetization of one of our crystals.

Indeed, we will show that the PCARS results give evidence of a large isotropic gap and a small nodal (or fully anisotropic) gap, and that density functional theory (DFT) calculations of the Fermi surface show a topological transition in the holelike Fermi surface sheets occurring at the same doping content. We speculate that the two facts may be correlated and that the symmetry of the OP on the closed holelike pockets at \( x \geq 0.06 \) can be modeled as a 3D d-wave one, with two nodal planes intersecting each other along the \( \Gamma-Z \) line.

2. Experimental results

The \( \text{Ca(Fe}_{1-x}\text{Co})_2\text{As}_2 \) single crystals used in this work were grown from Sn flux, as described in [4]; they looked very similar to those shown in [2], i.e. they were plate-like, with mirror surfaces and the \( c \) axis perpendicular to the plate. The inset to figure 1 shows the DC magnetization of one of the crystals as a function of temperature. The transition sets in at \( T^\text{on}_c = 20 \) K, and an effective \( T^\text{eff}_c = 17 \) K can be determined by extrapolating the linear part of the curve. The rather broad transition is common to all the state-of-the-art single crystals grown at present [1, 2, 5, 6]. The broadening might be related to chemical inhomogeneities over nanoscopic length scales [2] that are not detectable by energy- or wavelength-dispersive x-ray spectroscopies [2, 1], or might be intrinsic to the Ca-122 system because of its high sensitivity to chemical and physical pressure [2]. In any case, the large transition width is not detrimental to point-contact measurements since PCARS is a local probe not only of the superconducting order parameter but also of the critical temperature. The temperature of the best contacts always fell between \( T^\text{on}_c \) and \( T^\text{eff}_c \). By comparison with the phase diagram in Hu et al [1], we can conclude that the real Co content of our crystals is \( x = 0.060 \pm 0.005 \).

The point-contact measurements were performed by using the ‘soft’ technique described elsewhere [7], i.e. by putting a small drop of Ag paste (\( \phi \approx 50 \mu \text{m} \)) on a fresh side surface of the crystal (exposed by breaking it) so that the current is mainly injected along the \( ab \) planes. Contacts made in this way are more mechanically and thermally stable than those made by using the standard needle-anvil technique (i.e. by gently pressing a sharp metallic tip against the same surface). The \( J-V \) characteristics of the normal metal/superconductor junctions obtained in this way were measured and numerically differentiated to obtain the conductance curves, \( dI/dV \) versus \( V \).

Figure 1 shows the temperature dependence, from 2.21 K up to 20.12 K, of the raw conductance curves of an \( ab \)-plane point contact whose normal-state resistance was \( R_N \approx 3 \) Ω. The Andreev signal disappears completely at \( T^A_c \approx 18 \) K, leaving a slightly bent but rather symmetric normal-state conductance. Despite the rather low resistance, the conduction through the contact is very likely to be ballistic. This is witnessed by the fact that the high-energy tails of the curves coincide with one another (which means that there is no contribution from the Maxwell term in the contact resistance [7]), by the absence of features associated with the current-induced suppression of superconductivity (such as characteristic dips [8]) and by the fact that the Andreev critical temperature \( T^A_c \) is not suppressed with respect to the bulk \( T_c \). The small resistance of the contact can indeed be reconciled with the evidence of ballistic conduction if multiple parallel nanocontacts are formed between Ag grains in the paste and the crystal surface, which is not only very reasonable but practically unavoidable in our point contacts.

Even at a first glance, the shape of the curves in figure 1 looks not to be compatible with a nodeless gap (or multiple nodeless gaps, as observed in most 1111 compounds [9] and in 122 compounds such as \( \text{Ba(Fe}, \text{Co})_2\text{As}_2 \) [10] and \( \text{(Ba, K)Fe}_2\text{As}_2 \) [11]). It must be noted that this shape is common to all the conductance spectra in the Andreev-reflection regime we obtained in different \( ab \)-plane contacts, so that it appears to be an intrinsic feature of the material and not an artifact. A narrow maximum at zero bias, such as that shown in the curves in figure 1, has been observed in many unconventional superconductors, including iron pnictides [12–14]; however, this kind of feature can have either intrinsic or extrinsic origins. For example, it could be the hallmark of a Josephson supercurrent—but this requires either the formation of some superconductor-insulator-superconductor (SIS) junction, or intrinsic Josephson effects as in \( c \)-axis tunneling in layered materials, both excluded here. Alternatively, the zero-bias conductance peak (ZBCP) could also be due to non-perfectly ballistic conduction (already excluded for our contacts), or local damage or deformation of the surface due to the pressure applied by the tip (see [7] and references therein), but this is clearly impossible with our technique. Finally—and we claim this is the case in our measurements—it could be a direct consequence of the anisotropy of the superconducting order parameter.

If the OP has nodal lines crossing the Fermi surface, as in the d-wave or in the nodal \( s \) symmetries proposed for 1111 pnictides in suitable conditions [15], its change of sign can give rise—for suitable directions of current injection—to constructive interference between holelike and
electronlike quasiparticles that results in zero-energy bound states localized at the interface [16]. These, in turn, manifest themselves in a zero-bias peak in the Andreev-reflection conductance. In principle, also the 3D nodal symmetry of the OP proposed for the isovalent-doped Ba-122 system [3, 17] could give rise to these effects. A similar peak can also appear in diffusive metal/nodal superconductor junctions as a result of the coherent Andreev reflection in the former [18].

We have, however, recently shown [9] that, in the presence of a large spectral broadening (as in our contacts on Ca(Fe, Co)₂As₂) it becomes impossible to distinguish a true zero-bias conductance peak (due to a nodal OP) from the zero-bias maximum that is the natural consequence of a OP with zeros on the Fermi surface. Strictly speaking, we can thus conclude that the shape of the conductance curves of figure 1 indicates that at least some component of the OP is zero somewhere on the Fermi surface—although we cannot say whether it changes sign or not.

To extract quantitative information on the number, amplitude and symmetry of the OPs (with the aforementioned limitations), it is necessary first to normalize the experimental conductance curves and then to compare them with the predictions of suitable models for Andreev reflection. The first step is straightforward in this case, since the normal-state conductance is well defined and, judging from the superposition of the high-energy tails of the curves, seems not to depend very much on temperature below $T^\Lambda_0$. Therefore, we will divide every conductance curve at $T < T^\Lambda_0$ by a curve measured just above $T^\Lambda_0$. In the case of figure 1, we used the curve at 18.19 K. As for the fit of the normalized conductance curves, one might choose for simplicity a two-band 2D-BTK model [16]. This model is actually based on the assumption that the Fermi surface is approximately cylindrical, but in 122 compounds a more or less pronounced warping of the Fermi cylinders is present already in the parent compounds, and increases with doping [17]. In section 3 we will show that the 2D approximation for the Fermi surface is not justified in Ca(Fe₀.₉₀Co₀.₀₄)₂As₂ so that the more complicated two-band 3D-BTK model must be used instead.

3. Calculation of the Fermi surface

To gain further insight into the real band structure of Ca(Fe₁−ₓCoₓ)₂As₂ we performed first-principles DFT calculations by using the Elk full-potential linearized augmented-plane wave (FP-LAPW) Code (http://elk.sourceforge.net/) within the generalized gradient approximation (GGA) approach for the exchange correlation potential. A virtual-crystal approximation (VCA) was used to mimic the Co doping, owing to the fact that VCA works particularly well for Fe–Co substitutions. In the absence of direct experimental information on the low-temperature lattice constants of Co-doped CaFe₂As₂, we used the values $a = b = 3.925$ Å and $c = 11.356$ Å, calculated according to [19] for CaFe₂As₂ at low temperature and null pressure. They were obtained as follows. First, we calculated the lattice parameters of the orthorhombic phase at pressure $P = 0$ from the total energy minimization and by considering the stripe antiferromagnetic order. Then, in order to calculate the Fermi surface at $P = 0$ in the tetragonal phase, we made the structure tetragonal by using for the lattice constant $a$ the average of $a$ and $b$ determined in the orthorhombic phase. Incidentally, these values of the theoretical lattice parameters are in good agreement with the experimental ones measured in the tetragonal phase of the parent compound at 300 K (and pressure $P = 0.8–1$ GPa [20]). Owing to the small doping dependence of the lattice constants at high temperature [1], we assumed these values to be a good first approximation to the real ones at the doping content of our interest ($x < 0.08$). Starting from the calculated equilibrium phase, and always considering the antiferromagnetic phase, an optimized internal parameter $z_{A_2} = 0.7306$ (such that $h_{A_2} = 1.309$ Å) was obtained. The charge density was thus integrated over $8 \times 8 \times 4$ $k$ points in the Brillouin zone and the band structures, as well as the Fermi surfaces, were calculated in the non-magnetic body-centered tetragonal phase.

Figures 2(a)–(c) show the Fermi surface (plotted on a $40 \times 40 \times 40$ $k$-points grid) of Ca(Fe₁−ₓCoₓ)₂As₂ for $x = 0.04$, $x = 0.06$ and $x = 0.08$, respectively. The evolution of the holelike Fermi surface towards a full 3D nature on increasing $x$ is clear. In the $x = 0.04$ case (panel (a)), three holelike FS sheets are present: two inner closed pockets centered at $Z$ and an outer, strongly warped cylinder extending along the $\Gamma$–$Z$ line. At $x = 0.06$ the smaller pocket disappears and the warped cylinder undergoes a topological transition splitting into separated closed surfaces centered at $Z$. It is worth noting that this trend is qualitatively robust against the details of the calculations, although some refinement (e.g. the relaxation of

![Figure 2](https://example.com/figure2.png)
the cell parameters or their small doping dependence) could slightly change the doping content at which the complete separation of the two closed sheets around Z occurs.

4. Analysis of the experimental data

In these conditions, it is necessary to use the complete 3D-BTK model [9] to fit our conductance curves. We first schematized the actual Fermi surface of Ca(Fe$_{0.94}$Co$_{0.06}$)$_2$As$_2$ by means of a one-sheeted hyperboloid of revolution (that simulates the electronlike FS sheets centered in X) and an oblate spheroid (to simulate the 3D holelike FS centered in Z). Figure 2(d) shows one half of the two model Fermi surfaces (corresponding to the upper half of the real Brillouin zone shown in panel (b)). The proportions between the size of the two FSs have been respected, though the distance between them (which is irrelevant in the following calculation) has been increased for clarity with respect to the actual distance shown in figure 2(b). For this reason the Brillouin zone sketched in figure 2(d) appears wider that the actual one and should be just considered as a guide to the eye.

Before proceeding, we had to make a guess about the symmetry of the two OPs residing on the two FSs. According to recent theoretical predictions, in the Ba-122 system an evolution from a pure ±$\pi$ gap symmetry towards a peculiar 3D nodal one is expected when the pnictogen height $z_{\text{As}}$ is reduced, for example by isovalent P doping [3]. It is interesting to note that the reduction in $z_{\text{As}}$ simply leads to a larger holelike Fermi surface around the Z point, similar to the one depicted in figure 2(a). In these conditions, a three-dimensional sign change of the OP takes place within this Fermi surface near the edge of the Brillouin zone [3, 17], while the gap on the electron Fermi surface always remains fully open. When the holelike FS is further deformed into separated closed surfaces, as in our case, the symmetry of the OP relevant to it must evolve as well. We guess that the result of this evolution can be expressed as a 3D d-wave OP, whose expression in the reciprocal space is $\Delta(\theta, \phi) = \Delta \cos(2\theta) \sin^2(\phi)$, $\theta$ being the azimuthal angle in the $(k_x, k_y)$ plane of the reciprocal space and $\phi$ the inclination angle. This OP features two orthogonal nodal planes intersecting along the $\Gamma$–Z line. As for the electronlike FS sheets, we assume that the relevant OP keeps a s-wave symmetry. The amplitude of the two OPs are shown in figure 2(d) as gridded surfaces.

Now the theoretical conductance curve can be calculated, within some reasonable assumptions, by means of equation 9 of [9]. The two bands will be indicated by the subscripts 1 (holelike band) and 2 (electronlike band). There are seven degrees of freedom in the problem: the OP amplitudes $\Delta_1$ and $\Delta_2$, the barrier parameters $Z_1$ and $Z_2$, the broadening parameters $\Gamma_1$ and $\Gamma_2$, plus the angle $\alpha$ between the direction of current injection (here in the basal plane) and the lobes of the d-wave gap $\Delta_1$. The relative weight of the two bands in the conductance is not an adjustable parameter but directly follows from the geometry of the problem (i.e. the shape of the relevant Fermi surface sheets and the direction of current injection) and from the choice of the $Z$ parameters.

Figure 3(a) shows some of the experimental conductance curves of figure 1, normalized and compared to the best-fitting curves (lines) calculated within the aforementioned model. The parameters that allow fitting the low-temperature curve are the following: $\Delta_1 = 1.7$ meV, $\Delta_2 = 5.3$ meV, $Z_1 = 0.4$, $Z_2 = 0.185$, $\Gamma_1 = 1.66$ meV and $\Gamma_2 = 6.4$ meV. The high values of the $\Gamma$ parameters are due to the smallness of the Andreev signal. Such large values could be avoided by introducing an additional scaling parameter $S$, sometimes used in these cases [21], for instance to account for a fraction of injected current that does not give rise to Andreev reflection. To reproduce the ZBCP a misorientation angle $\alpha = \pi/8$ was necessary. This value is quite reasonable because the side surfaces of the crystals, on which the point contacts are made, are rather irregular and thus the angle of current injection in the basal plane is not well defined. Indeed, it
can be shown that the conductance curve one would obtain by averaging over all angles \( \alpha \) between 0 and \( \pi/4 \) is best described by using a single \( \alpha \approx \pi/8 \). Note that the values of the gaps given by this fit are very similar to those obtained by using a simplified 2D-BTK model as in [9], but the Z values are systematically smaller because of the Z-enhancing effect discussed elsewhere [9]. Moreover, as already explained, similar results could be obtained by fitting the experimental normalized curves with a s-wave OP on one band and a fully anisotropic s-wave OP (with zeros on the Fermi surface) on the other band. It is also worth noting that a simpler model with a smaller number of parameters fails to reproduce the experimental conductance curves. The left inset to figure 3(a) shows that a single-gap d-wave model (with four parameters) can only reproduce either the zero-bias maximum (blue line) or the higher energy portion of the conductance curve (red line). Though the latter curve is better, the one-band fit is still less accurate than the two-band one; moreover, it would give an unreasonably large value of the gap.

Figure 3(b) shows the temperature dependence of the normalized conductance curves (symbols) of a different point contact, whose normal-state resistance was \( R_N \approx 3.4 \Omega \). As in panel (a), the experimental data are compared to the 3D-BTK best-fitting curves calculated for a large isotropic gap and a smaller d-wave gap. The fitting parameters at the lowest temperature are \( \Delta_1 = 1.6 \text{ meV}, \Delta_2 = 5.3 \text{ meV}, Z_1 = 0.35, Z_2 = 0.265, \Gamma_1 = 2.2 \text{ meV}, \Gamma_2 = 6.8 \text{ meV} \) and finally \( \alpha = \pi/7.5 \). The values of the broadening parameters are even higher than in the previous case, because the signal is smaller, but the values of the gaps are in very good agreement with those obtained from the curves in figure 3(a). Some small additional structures are also visible in the lowest-temperature curves at about 8 and 12 meV that are not reproduced by the fit. These features might be related to a strong electron–boson interaction, as shown in [9, 10]. However, in this case we did not observe them in a reproducible way, moreover they disappear soon with increasing temperature and the Andreev signal is too small to draw any conclusion in this sense. It is worth mentioning that such features contribute very little to the uncertainty of the gap values, which mostly arises from the smallness of the Andreev signal, as discussed below.

As shown in figures 3(a) and (b), we were able to fit the conductance curves at any temperature up to \( T^A_c \) by using the same model and keeping the barrier parameters \( (Z_1, \text{ and } Z_2) \) and the angle \( \alpha \) independent of temperature. The amplitudes of the gaps, \( \Delta_1 \) and \( \Delta_2 \), are shown as a function of temperature in the right insets to figures 3(a) and (b). Error bars indicate the spread of gap values obtained in different fits of the same curve, with different combinations of the remaining parameters \( (Z_1, \Gamma_1 \text{ and } \alpha) \). Because of the smallness of the Andreev signal, this spread at low temperature is of the order of 0.2 meV on \( \Delta_1 \) and 0.3 meV on \( \Delta_2 \), and increases on heating. For the sake of comparison, the same insets also show two Bardeen–Cooper–Schrieffer (BCS)-like trends—approximately expressed as \( \Delta_1(T) = \Delta_1(0) \tanh(1.74 \sqrt{(T^A_c/T) - 1}) \), where \( \Delta_1(0) \) is the low-temperature gap and \( T^A_c \) the experimental critical temperature of the contact. While the data of panel (a) look fairly compatible with these trends, some suppression of the gaps are observed in the inset to panel (b) in the proximity of \( T^A_c \). The gap ratios turn out to be \( 2\Delta_1/k_B T^A_c = 2.2-2.3 \) (much smaller than the BCS value for d-wave superconductors) and \( 2\Delta_2/k_B T^A_c = 6.85-7.25 \), which is instead suggestive of a strong electron–boson coupling.

5. Conclusions

We have presented the results of point-contact Andreev-reflection measurements in state-of-the-art Ca(Fe\(_{1-x}\),Co\(_x\))\(_2\)As\(_2\) single crystals with effective \( x = 0.060 \pm 0.005 \). The point contacts were made on the side surface, i.e. the current was mainly injected along the \( ab \) planes. We have shown that the shape of the PCARS spectra can be nicely explained within a multiband (multigap) scenario where a large isotropic gap \( \Delta_2 \) coexists with a smaller nodal (or fully anisotropic) gap \( \Delta_1 \). These two gaps can be associated with the electronlike and holelike Fermi surface sheets, respectively, on the basis of recent theoretical papers [17, 3] that predict the development of 3D nodes in the OP of the holelike FS when the size of the latter around the \( Z \) points of the Brillouin zone increases (for example, in Ba–122, as a consequence of P doping). We have shown by first-principles DFT calculations that Co doping in CaFe\(_2\)As\(_2\) has a similar, but even more pronounced, effect on the holelike FS, which actually undergoes a topological transition at about \( x \approx 0.06 \) from a slightly warped cylinder along the \( \Gamma–Z \) line to separate ellipsoids centered around the \( Z \) points. In these conditions, we have assumed that the symmetry of the nodal gap \( \Delta_1 \) can be described by a 3D d-wavefunction, while \( \Delta_2 \) keeps a s-wave symmetry. The fit of the experimental spectra with a two-band 3D-BTK model using a model Fermi surface that simulates the calculated one gives \( \Delta_2 = 5.3 \pm 0.3 \text{ meV} \) and \( \Delta_1 = 1.6 \pm 0.2 \text{ meV} \).

In conclusion, our results seem to confirm and extend the existing predictions about the evolution of the gap symmetry as a function of isovalent doping, and might suggest a connection between the topological transition of the holelike FS and the appearance of nodal points in the gap along the \( \Gamma–Z \) line. Further point-contact measurements in other Fe-based compounds and with the current injected in different directions could further substantiate this picture.

Acknowledgments

The work at Politecnico di Torino was supported by the PRIN Project No. 2008XWLF9-005. The work at ETHZ was supported by the NCCR Project MaNEP. The authors thank G Profeta very much for suggestions and enlightening discussions. RSG acknowledges the Max Planck Institute for Solid State Research in Stuttgart, where the measurements were carried out.

References

[1] Hu R, Ran S, Bud’ko S, Straszheim W E and Canfield P C 2011 Single crystal growth and superconductivity of Ca(Fe\(_{1-x}\),Co\(_x\))\(_2\)As\(_2\) unpublished (arXiv:1111.7034)
[2] Harnagea L et al 2011 Phys. Rev. B 83 094523
[3] Suzuki K, Usui H and Kuroki K 2011 J. Phys. Soc. Japan 80 013710
[4] Matusiak M, Bukowski Z and Karpinski J 2010 Phys. Rev. B 81 020510(R)
[5] Abdel-Hafiez M, Harnagea L, Singh S, Stockert U, Wurmehl S, Klingeler R, Wolter A U B and Büchner B 2011 Calorimetric study of the superconducting and normal state properties of Ca(Fe1−xCo)x2As2 unpublished (arXiv:1109.3135)
[6] Kumar N, Nagalakshmi R, Kulkarni R, Paulose P L, Nigam A K, Dhar S K and Thamizhavel A 2009 Phys. Rev. B 79 012504
[7] Daghero D and Gonnelli R 2010 Supercond. Sci. Technol. 23 043001
[8] Sheet G, Mukhopadhyay S and Raychaudhuri P 2004 Phys. Rev. B 69 134507
[9] Daghero D, Tortello M, Ummarino G and Gonnelli R S 2011 Rep. Prog. Phys. 74 124509
[10] Tortello M, Daghero D, Ummarino G A, Stepanov V A, Jiang J, Weiss J D, Hellstrom E E and Gonnelli R S 2010 Phys. Rev. Lett. 105 237002
[11] Szabó P, Pribulová Z, Pristáš G, Bud’ko S L, Canfield P C and Samuely P 2009 Phys. Rev. B 79 012503
[12] Lu X et al 2010 Supercond. Sci. Technol. 23 054009
[13] Yates K A, Cohen L F, Ren Z A, Yang J, Lu W, Dong X L and Zhao Z X 2008 Supercond. Sci. Technol. 21 092003
[14] Samuely P, Szabó P, Pribulová Z, Tillman M, Bud’ko S L and Canfield P C 2009 Supercond. Sci. Technol. 22 014003
[15] Kuroki K, Usui H, Onari S, Arita R and Aoki H 2009 Phys. Rev. B 79 224511
[16] Kashiwaya S and Tanaka Y 2000 Rep. Prog. Phys. 63 1641–724
[17] Graser S, Kemper A F, Maier T A, Cheng H P, Hirschfeld P J and Scalapino D J 2010 Phys. Rev. B 81 214503
[18] Tanaka Y, Nazarov Y V, Golubov A A and Kashiwaya S 2004 Phys. Rev. B 69 144519
[19] Colonna N, Profeta G, Continenza A and Massidda S 2011 Phys. Rev. B 83 094529
[20] Mittal R et al 2011 Phys. Rev. B 83 054503
[21] Naidyuk Y G, Kvitnitskaya O E, Tiutrina L V, Yanson I K, Behr G, Fuchs G, Drechsler L N, Nenkov K and Schults L 2011 Phys. Rev. B 84 094516