Resonance states near a quantum magnetic impurity in single-layer FeSe superconductors with $d$-wave symmetry

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Abstract

In this work, we investigate the local density of states (LDOS) near a magnetic impurity in single-layer FeSe superconductors. The two-orbital model with spin–orbit coupling proposed in Agterberg et al (2017 Phys. Rev. Lett. 119 267001) is used to describe the FeSe superconductor. In the strong coupling regime, two impurity resonance peaks appear with opposite resonance energies in the LDOS spectral function. For strong spin–orbit coupling, the superconducting gap in this model is $d$-wave symmetric with nodes, the spatial distributions of the LDOS at the two resonance energies are fourfold symmetric, which reveals typical characteristic of $d$-wave pairing. When the spin–orbit coupling is not strong enough to close the superconducting gap, we find that the spatial distribution of the LDOS at one of the resonance energies manifests $s$-wave symmetry, while the pairing potential preserves $d$-wave symmetry. This result is consistent with previous experimental investigations.

Keywords: pairing symmetry, magnetic impurity, iron-based superconductivity

(Some figures may appear in colour only in the online journal)
between the impurity state and the superconducting state with wave-vector $\mathbf{k} = (k_x, k_y)$. Hereafter we set the hybridization to be short-ranged, so that it is wave-vector independent, $V_\mathbf{k} = V_0\sqrt{c_{\mathbf{k} s}^\dagger c_{\mathbf{k} s}}$ and $(c_{\mathbf{k} s})$ are the creation and annihilation operators of the electron state in single-layer FeSe with wave-vector $\mathbf{k}$ and spin-$s$. $N$ refers to the total number of wave-vectors in summation. The last term in equation (1) describes the free Hamiltonian of the superconducting state [46],

$$
H_{sc} = \sum_{\mathbf{k}} \sum_{s=\uparrow, \downarrow} \left\{ \varepsilon_0(\mathbf{k}) \tau_s \sigma_0 + \gamma_s(\mathbf{k}) \tau_s \sigma_0 + \tau_s [\gamma_s(\mathbf{k}) \sigma_s + \gamma_s(\mathbf{k}) \sigma_s] \right\} \psi_\mathbf{k} + \frac{1}{2} \sum_{\mathbf{k}} \left\{ \psi_\mathbf{k}^\dagger \left[ \Delta_d \tau_0 + \Delta_c \tau_0 \right] i \sigma_y \psi_\mathbf{k} + \text{h.c.} \right\}. \tag{4}
$$

Here $\psi_\mathbf{k} = \left( c_{\mathbf{k}1,\uparrow}, c_{\mathbf{k}2,\uparrow}, c_{\mathbf{k}1,\downarrow}, c_{\mathbf{k}2,\downarrow} \right)^T$ (the superscript $T$ refers to matrix transpose) is the four-component spinor description of the electron states with two orbital degrees of freedom described by $\tau_{x,y}$. Pauli matrices and two spin degrees of freedom described by $\sigma_{x,y}$. Pauli matrices. $\Delta_0$ and $\sigma_0$ are the 2 × 2 identity matrices in orbital and spin space, respectively. We need to emphasize that, as demonstrated in [46], the two effective orbitals are generally $k$-dependent linear combinations of the $x^2 - y^2$ and $(x, y, z)$ orbitals centered on Fe sites. h.c. in equation (4) means the Hermitian conjugate. $\varepsilon_0(\mathbf{k}) \pm \gamma_s(\mathbf{k})$ yield the two elliptic electron Fermi surface pockets near the M-point. They are given in the following tight-binding form [47],

$$
\varepsilon_0(\mathbf{k}) = t_1 \cos(k_x a) + \cos(k_y a) - \epsilon, \tag{5}
$$

$$
\gamma_s(\mathbf{k}) = t_2 \cos\left( \frac{k_x a}{2} \right) \cos\left( \frac{k_y a}{2} \right), \tag{6}
$$

$\gamma_s(\mathbf{k})$ and $\gamma_s(\mathbf{k})$ in equation (4) represent the spin–orbit couplings, they are given by,

$$
\gamma_s(\mathbf{k}) = -t_3 \sin(k_x a), \quad \gamma_s(\mathbf{k}) = -t_3 \sin(k_y a). \tag{7}
$$

The superconducting gap terms are given by $\Delta_d = \Delta_s \sin(k_x a) \sin(k_y a)$ and $\Delta_c = \Delta_0$ with $\Delta_0 = 11$ meV and $\Delta_s = -2.34$ meV. $\epsilon$ in equations (5)–(7) is the lattice constant, it is set to be 4 Å in this work [34, 48]. The other parameters are chosen as follows, $t_1 = 171.875$ meV, $t_2 = 150$ meV, $\epsilon = -288.75$ meV, $t_3 = v_{so}/a$, and $v_{so}$ represents the spin–orbit coupling strength. One can check that, in the continuous limit, the model Hamiltonian (4) tends to that given in [46]. Figure 1 shows the Fermi surfaces and superconducting gap on the Fermi surfaces for different spin–orbit couplings, one can find that they are consistent with those given in [46].

Now we study the model Hamiltonian (1). In the strong Coulomb interaction limit, $U \rightarrow \infty$, the double occupied state of electrons on the impurity can be excluded. This limit may be represented by introducing an auxiliary boson operator $b$ to reformulate the creation and annihilation operators of the impurity states, $(d^\dagger_i, d_i) = (f^\dagger_i b_i, b_i f_i)$. The extra degrees of freedom after introducing these boson operators are restricted by the constraint $\bar{Q} = b_i^\dagger b_i + \sum_i f_i f_i^\dagger = 1$. In the mean-field approximation, $b$ and $b^\dagger$ are replaced by their expectation.
Figure 1. The Fermi surfaces and the anisotropic superconducting gap on the Fermi surfaces for different spin–orbit couplings. (a) and (d): \( v_0 = 0 \), (b) and (e): \( v_0 = 12 \) meV \( \text{Å} \), (c) and (f): \( v_0 = 80 \) meV \( \text{Å} \).

value, \( b_0 = \langle b \rangle = \langle b^\dagger \rangle \), and the constraint is approximated by adding a term \( \lambda_0 \langle b^\dagger b + \sum_i f_i^\dagger f_i - 1 \rangle \) to the Hamiltonian, where \( \lambda_0 \) is a Lagrangian multiplier, it renormalizes the impurity energy. Both \( b_0 \) and \( \lambda_0 \) need to be determined self-consistently by minimizing the free energy. The mean-field Hamiltonian is given by,

\[
H_{\text{MF}} = \hat{H}_{\text{imp}} + \hat{H}_{\text{hyb}} + H_{\text{sc}} + \lambda_0 (b_0^2 - 1),
\]

where \( \varepsilon_d = \varepsilon_d + \lambda_0 \) is the renormalized impurity energy, \( \hat{V}_0 = b_0 V_0 \) is the renormalized hybridization. In the Bogoliubov–de Gennes (BdG) formalism, the mean-field Hamiltonian can be recast as,

\[
\hat{H}_{\text{MF}} = \lambda_0 (b_0^2 - 1) + \varepsilon_d + \frac{1}{2} \Phi^\dagger \Lambda \Phi + \frac{1}{2} \sum_k \Psi_k^\dagger h_{\text{BdG}} \Psi_k + \frac{1}{2 \sqrt{N}} \sum_k \left( \Psi_k^\dagger \hat{V} \Phi + \Phi^\dagger \hat{V}^\dagger \Psi_k \right),
\]

where \( \Phi = (f_1, f_2, f_1^\dagger, f_2^\dagger)^\dagger \) and \( \Psi_k = \left[ \psi_k, (\psi_k^\dagger) \right]^\dagger \) are the Nambu spinors. The matrices \( \Lambda \) and \( h_{\text{BdG}} \) are given by

\[
\Lambda = \xi_d \varsigma^\dagger \tau_0 \sigma_0,
\]

\[
h_{\text{BdG}} = \varepsilon_0(k) \varsigma^\dagger \tau_0 \sigma_0 + \gamma_0(k) \varsigma^\dagger \tau_x \sigma_0 + \gamma_0(k) \varsigma^\dagger \tau_x \sigma_x + \gamma_0(k) \varsigma^\dagger \varsigma \sigma_y - \varsigma^\dagger (\Delta_d \tau_0 + \Delta_c \tau_2) \varsigma,
\]

where \( \varsigma_0 \) and \( \varsigma_{xy,z} \) are the identity matrix and Pauli matrices in the Nambu spinor space, respectively. \( V \) is a \( 8 \times 8 \) matrix representation of the renormalized hybridization, whose elements are given by, \( V_{0} = V_0 (\delta_2z - i \delta_2y) \).

Using the standard functional integral techniques, we find that the free energy of impurity electrons is given by,

\[
F = \lambda_0 (b_0^2 - 1) + \varepsilon_d + \frac{1}{2} \beta \sum_{\omega_n} \text{tr} \ln |G_f(i\omega_n)|,
\]

where \( \beta = 1/k_B T \) is the inverse temperature, \( k_B \) is the Boltzmann constant and \( T \) represents temperature, \( \omega_n = \pi (2n + 1)/2 \) is the Matsubara frequency, \( G_f(i\omega_n) = [i\omega_n - \Lambda - \Sigma_f(i\omega_n)]^{-1} \) is the Green’s function of the impurity states expressed in imaginary-frequency representation, \( \Sigma_f(i\omega_n) = \frac{1}{2} \sum_k \hat{V}^\dagger G_c^{(0)}(i\omega_n, k) \hat{V} \) is the self-energy, and \( G_c^{(0)}(i\omega_n, k) \) is the unperturbed Green’s function of the electrons in single-layer FeSe. Minimizing the free energy, equation (13), we can find the self-consistent equations of \( b_0 \) and \( \lambda_0 \),

\[
b_0^2 + \frac{1}{2 \beta} \sum_{\omega_n} \text{tr} [G_f(i\omega_n) \varsigma_0 \sigma_0] = 0,
\]

\[
\lambda_0 b_0^2 + \frac{1}{2 \beta} \sum_{\omega_n} \text{tr} [G_f(i\omega_n) \Sigma_f(i\omega_n)] = 0.
\]

The LDOS near the magnetic impurity is obtained by using the analytic continuation, \( i\omega_n \rightarrow E + i0^+ \), according to

\[
N_c(E; R) = -\frac{1}{\pi} \text{Im} \left[ G_c(E; R, R)^+ \frac{1 + \varsigma_0 \tau_0 \sigma_0}{2} \right],
\]

where the Green’s function of the conduction electrons are determined by,

\[
G_c(E; R, R) = \frac{1}{N} \sum_{k, k'} e^{(k-k') R} G_c(E; k, k'),
\]

\[
G_c(E; k, k') = G_c^{(0)}(E, k) \left[ \delta_{kk'} + T(E) G_c^{(0)}(E, k') \right].
\]

Here \( T(E) = \hat{V} G_f(E) \hat{V}^\dagger / N \) is the \( T \)-matrix, \( G_c^{(0)}(E, k) \) is the analytic continuation of \( G_c^{(0)}(i\omega_n, k) \). By solving equations (14) and (15), we can get the values of \( \varepsilon_d \) and \( V_0 \), the Green’s functions \( G_f(E) \) and \( G_c(E; k, k') \), and the LDOS \( N_c(E; R) \).

In addition to the LDOS, the QPI is another important quantity to identify the symmetry of the resonance states. Generally, QPI is calculated as follows,

\[
\delta N_c(E; q) = \frac{1}{N} \sum_{R} e^{-i q \cdot R} \delta N_c(E; R),
\]

\[
\delta N_c(E; R) = N_c(E; R) - N_c^{(0)}(E; R),
\]
where $N_c^{(0)}(E, R)$ is the LDOS of the system without impurity, it takes the form of equation (16) with $G_c(E; R, R)$ being replaced by the free Green’s function $G_c^{(0)}(E; R, R)$. In the practical calculation, a mesh size of $4096 \times 4096$ in the momentum space is chosen, the fast Fourier transformation is used to evaluate the QPI. The infinitesimal imaginary part of the self-energy, it takes the form of equation (16) with $\Sigma = \frac{1}{2} \omega_0$, which is about one thousandth of the superconducting gap $\Delta_0$.

### 3. Numerical results

It is difficult to find the analytic solutions of $\lambda_0$ and $b_0$ due to the complex band structure. Here we show the numerical results. Figure 2(a) shows $b_0^2$ versus $\varepsilon_d$ for different hybridizations and different spin–orbit couplings. One can find that, for each line, when $|\varepsilon_d|$ is greater than a threshold value $|\varepsilon_d^{\text{th}}|$, the self-consistent equations (14) and (15) do not have a solution. In other words, $\varepsilon_d^{\text{th}}$ identifies the boundary between two different phases: when $|\varepsilon_d| < |\varepsilon_d^{\text{th}}|$, the magnetic impurity is coupled to the single-layer FeSe superconductor (strong coupling regime); when $|\varepsilon_d| > |\varepsilon_d^{\text{th}}|$, the magnetic impurity and the host material are decoupled (isolated magnetic moment regime). See, e.g, [15, 49, 50] and [51–54] for similar results in $d$-wave superconductors and marginal Fermi liquids.

Figures 2(b)–(d) show the phase diagram for three different spin–orbit couplings, $V_{so} = 0$, 12 meV Å and 80 meV Å. We find that the phase boundary (black lines) for the two cases with finite superconducting gaps, $V_{so} = 0$ and $V_{so} = 12$ meV Å, are very close to each other. The other case, $d$-wave pairing with nodes for $V_{so} = 80$ meV Å, has a larger strong coupling regime. In the following studies, we set the parameters located in regime (II) for all the three cases, such that the resonance states appears.

Now we analysis the LDOS, $N_c(E, R)$. The first column in figure 3 shows the LDOS close to the impurity, $N_c(E, R = 0)$, for the three typical different spin–orbit couplings we considered. The second and third columns show the corresponding spatial distributions of the resonance states in these cases at $E = \Omega_r$ and $E = -\Omega_r$, respectively (i.e. $\Omega_r \approx -0.388$ meV in figure 3(a), $\Omega_r \approx -0.462$ meV in figure 3(d) and $\Omega_r \approx -0.628$ meV in figure 3(g)). By comparing both the positions and the intensities of the resonance peaks in the first column of figure 3, we find that for the two cases with finite superconducting gaps, the resonance peaks are very close to each other (the red lines), though the LDOSs for the host materials (the gray lines) are different. Furthermore, the spatial distributions of the first resonance states at $E = \Omega_r$ for these two cases shown in figure 3(b) and (e) are similar to each other. They both look rotation symmetric and consistent.
with the experimental results given in [43]. More detailed analysis show that the radii of the localized resonance states in figures 3(b) and (e) are about 10 Å, which is the same magnitude of the resonance states around a Cr adatom in single-layer FeSe superconductor as shown in [43]. In addition, the spatial distributions of the second resonance states for $v_{so} = 0$ and 12 meV Å shown in figures 3(c) and (f) are also similar to each other. This result demonstrates that the LDOS is insensitive to the strength of the spin–orbit coupling as long as the superconducting gap is not closed. For the third case, $v_{so} = 80$ meV Å, the superconducting gap has eight nodes located in the regime between the two Fermi surfaces shown in figure 1(c). The LDOS of the quasiparticles for this case near the Fermi surface reveals linear behavior, $N_c(E) \propto |E|$ (See the gray line in figure 3(g) for more details). As shown in figures 3(h) and (i), the spatial distributions of LDOS corresponding to the resonance states near $E = \Omega_r$ and $E = -\Omega_r$ are fourfold rotation symmetric, which reveals the typical characteristic of $d$-wave pairing potentials.

Figure 4 shows the QPIs corresponding to the two resonance peaks, $E = \Omega_r$ for QPIs in the left column and $E = -\Omega_r$ for QPIs in the right column, respectively. These patterns reveal the distribution of the resonance states in the momentum space. One can find that the QPIs are rotation symmetric for the first two rows, figures 4(a)–(d), where the spin–orbit coupling strengths are chosen to be $v_{so} = 0$ and 12 meV Å for the first and second row, respectively. These results reveals the characteristic of the plain $s$-wave pairing as demonstrated in the experiment [43]. When the spin–orbit coupling is sufficient large to close the superconducting gap,
ie., $v_{so} = 80$ meV Å as shown in the third row of figure 4, the QPI patterns break the rotation symmetry evidently. As shown in figure 4(f), when the energy is close to the positive resonance energy, $E = \Omega_r$, the QPI exhibits fourfold rotation symmetry. When the energy is close to the negative resonance energy, $E = -\Omega_r$, the QPI exhibits even lower symmetry, it is twofold symmetric as shown in figure 4(e).

In addition, figure 5 shows different QPIs for fixed spin–orbit coupling $v_{so} = 12$ meV Å and different energies away from the two resonance peaks. $E = -22.5$ meV, $-16.5$ meV, $-9.0$ meV, $9.0$ meV, $16.5$ meV and $22.5$ meV are chosen. A remarkable feature in this plot is that the fourfold rotation symmetry is slightly broken when the energy is close to the superconducting gap, $E = \pm 9.0$ meV. One can find in figures 5(c) and (f) that there are two blue circles along the line $k_x = -k_y$ near the points $(k_x, k_y) \approx (\pi/3, -\pi/3)$ and $(-\pi/3, \pi/3)$. The expected fourfold rotation symmetric pattern near $(\pi/3, \pi/3)$ and $(-\pi/3, -\pi/3)$ is vanishing. When the energy is far away from the superconducting gap, i.e. $E = \pm 22.5$ meV as shown in figures 5(a) and (d), the plot behaves fourfold rotation symmetric. Theoretically, it is difficult to trace the origin of the symmetry broken in the QPI mappings. We have calculated the QPIs for vanishing spin–orbit coupling $v_{so} = 0$ and different energies $E = -22.5$ meV, $-16.5$ meV, $-9.0$ meV, $9.0$ meV, $16.5$ meV and $22.5$ meV (not plotted here). We find no evidence of fourfold rotation symmetry broken in this case, so we suspect that the fourfold rotation symmetry broken shown in figures 5(c) and (f) is induced by spin–orbit coupling.
4. Conclusion

We have investigated the LDOS and QPI of the resonance states near a quantum magnetic impurity in single-layer FeSe superconductor using the Anderson impurity model coupled to the BdG Hamiltonian proposed in [24]. In the strong coupling regime, the LDOS spectrum have two resonance peaks symmetrically located at the two sides of the Fermi energy. However, the intensities of the resonance peaks break the particle-hole symmetry. Three typical strengths of spin–orbit coupling are considered, i.e. (1) the vanishing spin–orbit coupling, (2) finite spin–orbit coupling keeping the superconducting gap nodeless ($v_{so} = 12 \text{ meV Å}$), and (3) strong spin–orbit coupling leading to the $d$-wave pairing with nodes ($v_{so} = 80 \text{ meV Å}$). The spatial distributions of the LDOS at the resonance energy $E = \Omega_r$ for the finite-gapped cases are spatial rotation symmetric, which behave like the traditional plain $s$-wave pairing symmetry and consistent with the experimental results. Especially, for the second case with $v_{so} = 12 \text{ meV Å}$, which has been named resilient nodeless $d$-wave pairing and has been used to explain the superconductivities in single-layer FeSe, our investigations give positive evidences. The third case with gap nodes displays typical behavior of the transitional $d$-wave pairing. The spatial distributions of the LDOS appear fourfold rotation symmetry.

The QPI for the first two cases, $v_{so} = 0$ and $v_{so} = 12 \text{ meV Å}$, behaves $s$-wave symmetric. For the third case, $v_{so} = 80 \text{ meV Å}$, the QPI behaves fourfold rotation symmetric for the positive resonance energy, $E = -\Omega_r$. It behaves two-fold rotation symmetric for the negative resonance energy, $E = \Omega_r$. For the resilient nodeless $d$-wave pairing, which has been considered [46] as the candidate pairing potential for monolayer FeSe, our calculation show that, when the energy is close to the superconducting gap, $E = \pm 9.0 \text{ meV}$ as shown in figures 5(c) and (f), the QPI of magnetic impurity scattering breaks the fourfold rotation symmetry. This can be

![Figure 5. QPI mappings for different energies away from the resonance peaks, $E = -22.5 \text{ meV}, -16.5 \text{ meV}, -9.0 \text{ meV}, 16.5 \text{ meV}$ and $22.5 \text{ meV}. The spin–orbit coupling is chosen to be $v_{so} = 12 \text{ meV Å}$. The other parameters are set to be as the same as in figure 3.](image-url)
considered as an experimental accessible approach to distinguish the resilient nodeless $d$-wave pairing proposed in [46] and the plain $s$-wave pairing without spin–orbit coupling, e.g. the QPI of magnetic impurity scattering in superconductors with plain $s$-wave pairing is expected to preserve the fourfold rotation symmetry.

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References

[1] Hewson A C 1997 The Kondo Problem to Heavy Fermions (Cambridge: Cambridge University Press)
[2] Balatsky A V, Vekhter I and Zhu J X 2006 Rev. Mod. Phys. 78 373
[3] Yu L 1965 Acta Phys. Sin. 21 75
[4] Shiba H 1968 Prog. Theor. Phys. 40 435
[5] Rusinov A I 1969 JETP Lett. 9 85
[6] Tsuei C C, Kirtley J R, Chi C C, Yu-Jahnes L S, Gupta A, Shaw T, Sun J Z and Ketchen M B 1994 Phys. Rev. Lett. 73 593
[7] Tsuei C C, Kirtley J R, Ren Z F, Wang J H, Raffy H and Li Z Z 1997 Nature 387 481
[8] Hoffman J E, McElroy K, Lee D H, Lang K M, Eisaki H, Uchida S and Davis J C 2002 Science 297 1148
[9] Wang Q H and Lee D H 2003 Phys. Rev. B 67 020511
[10] Hanaguri T, Kohsaka Y, Davis J C, Lupien C, Yamada I, Azuma M, Takano M, Ohishi K, Ono M and Takagi H 2007 Nat. Phys. 3 865
[11] Hanaguri T, Kohsaka Y, Ono M, Maltseva M, Coleman P, Yamada I, Azuma M, Takano M, Ohishi K and Takagi H 2009 Science 323 923
[12] Pan S H, Hudson E W, Lang K M, Eisaki H, Uchida S and Davis J C 2000 Nature 403 746
[13] Hudson E W, Lang K M, Madhavan V, Pan S H, Eisaki H, Uchida S and Davis J C 2001 Nature 411 920
[14] Bobroff J, Alloul H, MacFarlane W A, Mendels P, Blanchard N, Collin G and Marucco J F 2001 Phys. Rev. Lett. 86 4116
[15] Zhang G M, Hu H and Yu L 2001 Phys. Rev. Lett. 86 704
[16] Polkovnikov A, Sachdev S and Vojta M 2001 Phys. Rev. Lett. 86 296
[17] Zhu J X and Ting C S 2000 Phys. Rev. B 63 020506
[18] Vojta M, Zitlzer R, Bulla R and Pruschke T 2002 Phys. Rev. B 66 134527
[19] Polkovnikov A 2002 Phys. Rev. B 65 064503
[20] Dai X and Wang Z 2003 Phys. Rev. B 67 180507
[21] Baar S et al 2016 J. Supercond. Nov. Magn. 29 659
[22] Tsai W F, Zhang Y Y, Fang C and Hu J 2009 Phys. Rev. B 80 064513
[23] Bang Y, Choi H Y and Won H 2009 Phys. Rev. B 79 054529
[24] Akbari A, Eremen I and Thalmeier P 2010 Phys. Rev. B 81 014524
[25] Sau J D and Demler E 2013 Phys. Rev. B 88 205402
[26] Fu Z G, Zhang P, Wang Z and Li S S 2012 J. Phys.: Condens. Matter 24 145502
[27] Zha G Q and Jin Y Y 2017 Europhys. Lett. 120 27002
[28] Guo Y W, Li W and Chen Y 2017 Front. Phys. 12 127403
[29] Hsu F C et al 2008 Proc. Natl Acad. Sci. USA 105 14262
[30] Wang Q et al 2015 Nat. Mater. 117 059
[31] Wang Q Y et al 2012 Chin. Phys. Lett. 29 037402
[32] Liu D et al 2012 Nat. Commun. 3 931
[33] He S et al 2013 Nat. Mater. 12 605
[34] Tan S et al 2014 Nat. Mater. 14 285
[35] Lee J et al 2014 Nature 515 245
[36] Ge J F, Liu Z L, Liu C, Cao C L, Qian D, Xue Q K, Liu Y and Jia J F 2014 Nat. Mater. 14 285
[37] Fang C, Wu Y L, Thomale R, Bernevig B A and Hu J 2011 Phys. Rev. X 1 011009
[38] Zhou Y, Xu D H, Zhang F C and Chen W Q 2011 Europhys. Lett. 95 17003
[39] Yang F, Wang F and Lee D H 2013 Phys. Rev. B 88 100504
[40] Maier T A, Graser S, Hirschfeld P J and Scalapino D J 2011 Phys. Rev. B 83 100515
[41] Wang F, Yang F, Guo M, Lu Z Y, Xiang T and Lee D H 2011 Europhys. Lett. 93 57003
[42] Mazin I I 2011 Phys. Rev. B 84 024529
[43] Fan Q et al 2015 Nat. Phys. 11 946
[44] Kang J and Fernandes R M 2016 Phys. Rev. Lett. 117 217003
[45] Zhang Y, Lee J J, Moore R G, Li W, Yi M, Hashimoto M, Lu D H, Devereaux T P, Lee D H and Shen Z X 2016 Phys. Rev. Lett. 117 177001
[46] Agterberg D F, Shishidou T, O’Halloran J, Brydon P M R and Weinert M 2017 Phys. Rev. Lett. 119 267001
[47] Korshunov M M and Eremin I 2008 Phys. Rev. B 78 140509
[48] Cao H Y, Tan S, Xiang H, Feng D L and Gong X G 2014 Phys. Rev. B 89 014501
[49] Cassanello C R and Fradkin E 1996 Phys. Rev. B 53 15079
[50] Cassanello C R and Fradkin E 1997 Phys. Rev. B 56 11246
[51] Withoff D and Fradkin E 1990 Phys. Rev. Lett. 64 1835
[52] Gonzalez-Buxton C and Ingersent K 1998 Phys. Rev. B 57 14254
[53] Zhuang H B, Sun Q F and Xie X C 2009 Europhys. Lett. 86 58004
[54] Uchoa B, Rappoport T G and Castro Neto A H 2011 Phys. Rev. Lett. 106 016801