Theory of superconductivity in a three-orbital model of Sr$_2$RuO$_4$

Q. H. Wang$^1$, C. Platt$^2$, Y. Yang$^1$, C. Honerkamp$^{3,4}$, F. C. Zhang$^{5,6}$, W. Hanke$^2$, T. M. Rice$^{5,7}$, and R. Thomale$^{2,8}$

$^1$National Laboratory of Solid State Microstructures, Nanjing University - Nanjing, 210093, China
$^2$Theoretical Physics, University of W¨urzburg - D-97074 W¨urzburg, Germany
$^3$Institute for Theoretical Solid State Physics, RWTH Aachen University - D-52056 Aachen, Germany
$^4$JARA - FIT Fundamentals of Future Information Technology - Germany
$^5$Department of Physics, and Center of Theoretical and Computational Physics, The University of Hong Kong - Hong Kong, China
$^6$Department of Physics, Zhejiang University - Hangzhou, China
$^7$Institute for Theoretical Physics, ETH Zurich - CH-8093 Z¨urich, Switzerland
$^8$Institut de th´eorie des ph´enom`enes physiques, ´Ecole Polytechnique F´ed´erale de Lausanne CH-1015 Lausanne, Switzerland

received 13 September 2013; accepted 13 October 2013
published online 6 November 2013

PACS 74.20.-z – Theories and models of superconducting state
PACS 74.20.Rp – Pairing symmetries (other than s-wave)
PACS 71.27.+a – Strongly correlated electron systems; heavy fermions

Abstract – In conventional and high transition temperature copper oxide and iron pnictide superconductors, the Cooper pairs all have even parity. As a rare exception, Sr$_2$RuO$_4$ is the first prime candidate for topological chiral p-wave superconductivity, which has time-reversal breaking odd-parity Cooper pairs known to exist before only in the neutral superfluid $^3$He. However, there are several key unresolved issues hampering the microscopic description of the unconventional superconductivity. Spin fluctuations at both large and small wave vectors are present in experiments, but how they arise and drive superconductivity is not yet clear. Spontaneous edge current is expected but not observed conclusively. Specific experiments point to highly band- and/or momentum-dependent energy gaps for quasiparticle excitations in the superconducting state. Here, by comprehensive functional renormalization group calculations with all relevant bands, we disentangle the various competing possibilities. In particular, we show the small wave vector spin fluctuations, driven by a single two-dimensional band, trigger p-wave superconductivity with quasi-nodal energy gaps.

Very soon after the discovery of superconductivity in Sr$_2$RuO$_4$ [1], it was proposed that the superconducting (SC) pairing is of unconventional nature [2,3]. Later experiments have provided evidence that the Cooper pair in the SC state is of odd parity [4] with total spin equal to one [5]. Further evidence indicates the superconductivity to be chiral, breaking time-reversal symmetry [6,7]. Sr$_2$RuO$_4$ is thus the first prime candidate for a chiral p-wave superconductor [8–11], an interesting analogue of the neutral superfluid $^3$He. It has recently received great interest as, by suitable manipulations, it may support zero-energy Majorana bound states in vortices [12], the building block for topological quantum computing [13]. However, there are a number of outstanding issues associated with the chiral p-wave superconductivity in Sr$_2$RuO$_4$. First, p-wave spin triplet pairing is expected to be associated with spin fluctuations at small wave vector. However, the spin density wave (SDW) fluctuation observed in Sr$_2$RuO$_4$ is dominated by a large wave vector at higher temperatures and coexist with a feature at small wave vector at lower temperatures [14]. A resolution of this puzzle is vital to understand the superconductivity. Second, one would expect a spontaneous electric current at the edge of the RuO$_2$ layers as a result of the chiral SC state. The edge current, however, has not been observed conclusively in experiments [15]. One possible reason is that the edge current is very fragile and difficult to establish against disorder. Another possibility is a topological cancellation from hole-like and electron-like bands [16], posing the question as whether the SC state is topologically nontrivial at all. Third, the specific measurement reveals abundance of low-energy quasiparticle excitations...
below the transition temperature \[17\]. This would point to multiple gaps of very different magnitudes and/or deep minima in strongly momentum-dependent gap functions. Previous theories have mostly focused on either the two-dimensional (2D) $\gamma$-band derived from the $xy$ orbital \[18-21\], or the quasi-one-dimensional (1D) $\alpha$ and $\beta$ bands derived from the $xz$ and $yz$ orbitals \[16,22\]. However, the evolution of wave vectors of the spin fluctuations is beyond such models, and in fact can only be accounted for by a complete three-band model. This in turn dictates the properties of the SC state mentioned above. A microscopic theory for ruthenate superconductivity should explain both SDW and SC fluctuations at different energy or temperature scales. A perturbative treatment of the three-band model led to the conclusion that spin triplet pairing is favorable but disturbed by the incommensurate spin fluctuations, questioning the role of spin fluctuations for triplet pairing \[23\]. In this paper, we apply the functional renormalization group theory (FRG) \[20,24\] to study a 3-band Hubbard model including both the 2D-$\gamma$ and the quasi-1D ($\alpha, \beta$) Fermi sheets as suggested by first-principle calculations and angle-resolved photoemission (ARPES) experiments \[9,25\]. The FRG is particularly promising to address the multi-scale energy issues in ruthenates. In recent years the FRG has been expanded to treat 2D multi-orbital systems such as the iron pnictides and selenides \[26-28\] and candidate models of topological superconductors with or without time-reversal symmetry \[29-33\]. In addition to addressing the pairing symmetry and energy scale, FRG gives information on the relative strength and wave vector of competing orders in the particle-hole channels. Our results show that the SDW fluctuations are mainly driven by the two 1D bands at large wave vectors, followed by the 2D band at small wave vectors as the energy scale is lowered in the RG flow. The latter triggers at an even lower-energy scale $p$-wave superconductivity, which is dominated by the 2D band and has a highly anisotropic gap and deep minima near the Brillouin zone boundary. Our theory predicts chiral edge modes and thus an edge current. However, the large-gap anisotropy indicates the fragility of the chiral edge modes against perturbations such as disorder, rendering the detection of edge current hard to accomplish. Our prediction on the strong SDW fluctuations at a small wave vector at low temperatures can be tested in further neutron scattering experiment, and the prediction on the strongly anisotropic gap function in momentum space should be tested in ARPES with high resolution at low temperatures.

The model we consider is described by the Hamiltonian

$$
H = \sum_{k,\sigma} \psi_{k\sigma}^\dagger \begin{bmatrix} a & b \end{bmatrix} \psi_{k\sigma} + U \sum_{i,a} n_{ia\uparrow} n_{ia\downarrow} + U' \sum_{i,a>b} n_{ia\uparrow} n_{ib\downarrow} \\
+ J \sum_{i,a>b,\sigma,\sigma'} \psi_{i a\sigma}^\dagger \psi_{i b\sigma'} \psi_{i a\sigma'}^\dagger \psi_{i b\sigma} \\
+ J' \sum_{i,a\neq b} \psi_{i a\uparrow}^\dagger \psi_{i a\downarrow} \psi_{i b\downarrow}^\dagger \psi_{i b\uparrow}^\dagger.
$$

Here, $k$ denotes the momentum, $\sigma$ the spin, $i$ the lattice site, and $a$ and $b$ the orbital labels, with $\psi_{i a\uparrow}$ annihilating an electron in $d_{xz}$, $d_{yz}$ and $d_{xy}$ orbitals, respectively. The local interaction parameters include intraorbital ($U$), interorbital ($U'$), and pair hopping ($J'$). The matrix dispersion function $\epsilon_{k}^{ab}$ has the following nonzero elements: $\epsilon_{k}^{11} = -2t_1 \cos k_x - \mu$, $\epsilon_{k}^{22} = -2t_1 \cos k_x - \mu$, $\epsilon_{k}^{12(21)} = -4t_2 \sin k_z \sin k_y$, and $\epsilon_{k}^{33} = -2t_1' (\cos k_x + \cos k_y) - 4t_2' \cos k_x \cos k_y + \Delta - \mu$, where in dimensionless units, $t_1 = 1$, $t_2 = 0.1$, $t_1' = 0.8$, $t_2' = 0.35$, $\Delta = -0.2$ is the crystal field splitting, and $\mu = 1.1$ is the chemical potential. This set of parameters produces the band structure shown in fig. 1(a). The inset shows the Fermi surface, which resembles closely what is observed experimentally \[9,25\]. The corresponding normal state density of states is shown in fig. 1(b). There are van Hove points at the $X$ points on the $\gamma$-band close to the Fermi level, while the band edge anomalies of the $\alpha$ and $\beta$ bands are far from the Fermi level.

As known for such a system with partial nesting and van Hove singularities near the Fermi level, there will be various competing and mutually interacting collective fluctuations in density wave and pairing channels. This physics can be investigated appropriately by FRG. It provides a coupled flow of wave vector resolved effective interactions in all particle-particle and particle-hole channels vs. a running energy scale $\Lambda$. We use the singular-mode FRG (SMFRG) \[28,30,34\] to gain benefit of resolving the interactions throughout the Brillouine zone in terms of form factors, and use the multi-patch FRG \[20,24\] to check, with increased angular resolution, that no important form factors have been left out. From the combination the dominant ordering tendencies can be most suitably addressed. In both schemes, the effective interaction at a given scale can always be decomposed as

$$
\Gamma^{ab;cd}(k, k', q) \rightarrow \sum_{m} S_{m}(q)\phi^{ab}_{m}(k, q)[\phi^{cd}_{m}(k', q)]^*,
$$

either in the SC, spin or charge channels. Here $a, b, c, d$ are orbital or band labels, $q$ is the associated collective wave vector, and $k$ (or $k'$) is an internal momentum of
The correlation between the emergence of small-$\mathbf{q}$ pairing modes. Arrows indicate level crossings associated with the evolution of the $q/\pi$ in the spin channel in (a) and the pairing symmetries in (b). The vertical dashed line highlights the correlation between the emergence of small-$\mathbf{q}$ spin feature and the $p$-wave pairing tendency.

In fig. 2 we show the SMFRG flow of the leading eigenvalues in the (a) spin and (b) SC channel, for bare interactions $(U, U', J, J') = (3.2, 1.3, 0.3, 0.3)$. In (a), changes in the dominant wave vector of the spin interaction are marked by arrows. At high scales, the spin channel dominates over the SC channel. The dominant spin-fluctuation wave vector evolves from $\mathbf{q} = (1,1)\pi$ to $\mathbf{q} \sim \mathbf{q}_1 = (0.625, 0.625)\pi$ as $\Lambda$ decreases. For $\Lambda < 5 \times 10^{-3}$, a further level crossing to $\mathbf{q} \sim \mathbf{q}_2 = (0.188, 0.188)\pi$ occurs. We checked the form factors to find that the $\mathbf{q}_2$-feature comes dominantly from the $\gamma$-band, while the $\alpha$ and $\beta$ bands mainly contribute to the $\mathbf{q}_1$-feature. The spin response at $\mathbf{q}_2$ is due to the proximity to the van Hove singularity in the $\gamma$-band mentioned previously. This spots an effect that cannot be detected in an analysis for vanishingly small interactions [16], as a finite interaction scale is needed for the proximate van Hove points to come into play. The evolution of the spin-fluctuation peak from larger to small $\mathbf{q}$ with decreasing energy scale is in good qualitative agreement with neutron scattering experiments [14] where a similar change is observed as a function of temperature. The charge channel (not shown) is screened down in the flow and only re-enhanced weakly as $\Lambda$ decreases. At lowest scales, both spin and charge channels saturate due to imperfect nesting.

In the inset of fig. 2, we plot the leading spin-channel eigenvalues $S_{\text{SDW}}(\mathbf{q})$ vs. $\mathbf{q}$ at the final stage of the RG flow. We see that the interaction in the spin channel peaks at $\mathbf{q}_2$, but the amplitude at $\mathbf{q}_1$ is also sizable. In both cases, the spin bilinears correspond to onsite spins. The attractive pairing interaction is induced at intermediate scales via the spin channel. As the dominant spin fluctuation vector changes during the flow, the dominant pairing fluctuations also undergo changes as a function of $\Lambda$. In fig. 2(b), we show the 10 leading attractive eigenvalues of the pairing channel. At low scales, the strongest growing eigenvalue belongs to a $p$-wave mode which is twofold degenerate due to the underlying $C_{4v}$ symmetry. By comparing the flow in the spin channel in (a), we see that this pairing mode is already seeded and enhanced as the $\mathbf{q}_2$-feature shows up (the correlation is shown by the vertical dashed line), supporting the interpretation that close-to-ferromagnetic spin fluctuations drive triplet $p$-wave pairing in this case. The final portion of the SC flow is log-linear in $\Lambda$, consistent with the fact that the spin and charge channels saturate and decouple from the SC channel in the lowest energy range.

We now analyze the detailed pairing function of the $p$-wave state. Figure 3(a) shows the form factors of the two degenerate $p$-wave functions. For the Fermiology and interaction regime considered, the gap function in SMFRG turns out to be much smaller on the $\alpha$ and $\beta$ bands than on the $\gamma$-band, in general agreement with multi-patch FRG. On the $\gamma$-band, the degenerate gap form factors from the SMFRG can be written approximately as $p_\alpha = p_1 \sin k_x + p_2 \cos k_z \sin k_x$, and $p_\beta = p_1 \sin k_x + p_2 \cos k_z \sin k_x$, where $p_1/p_2 = -0.4375$. Thus, it is worth noting that pairing on the next-nearest bond is important. In the ordered state, as confirmed by a mean field calculation using the
renormalized pairing interaction, the favorable state resulting from the $p$-wave instability is the chiral $p \pm ip'$ state, as the system maximizes condensation energy by breaking time-reversal symmetry. The gap amplitude $|\Delta(k)|$ in this case is shown in fig. 3(b) on the Fermi surface. For better quantitative clarity, in fig. 3(c) we plot the Fermi angle $\theta$-dependence of $|\Delta(\theta)|$ on the Fermi pockets $\alpha$ (blue), $\beta$ (red) and $\gamma$ (green). Since the amplitudes on the $\alpha$ and $\beta$ pockets are very small, they are enlarged (by a factor of 20) for better visibility. Near $X/Y$, the $\gamma$-band gap amplitude shows deep minima. This is understood as follows. In general, $p$-wave pairing is stabilized by attractive (repulsive) interactions upon forward (backward) scattering. The umklapp contribution to backward scattering, however, involves only a small momentum transfer, leading to a destructive interference. In this sense, the $p$-wave pairing for such a Fermi surface cannot benefit from the enhanced density of states near $X/Y$, and correspondingly, the energy scale for it is small (we get $\sim 0.1\text{meV}$). The low critical scale is also consistent with the late emergence of the small-$q$ spin fluctuations shown in fig. 2(a). The depth of the gap minima is enhanced by the second-nearest-neighbor pairing $p_2$ with opposite sign to $p_1$. (We checked that the other sign would lead to much less deep minimum. The sign selection is apparently not due to symmetry, but the correct sign is more favorable energetically by squeezing the region of small gap values.) The deep minimum feature is likewise found in multi-patch FRG (fig. 3(d)). There, the angular variation is found to be slightly stronger than for SMFRG, while the general behavior is the same. Similarly, also for the multi-patch FRG, the gap on the $\alpha$ and $\beta$ pockets is rather small, even below the minimum on the $\gamma$-band. While the qualitative behavior is similar to that of ref. [21], the anisotropy and band-selectiveness of the pairing is even stronger in our infinite-order approach.

The deep gap minima on the $\gamma$-band define a small gap scale of roughly a tenth of the gap maximum. We discuss two consequences of this small energy scale: fig. 4 shows the energy spectrum of the $p + ip'$ SC meanfield Hamiltonian on an infinite ribbon along the $y$-direction, with open boundary conditions along $x$. The energy eigenvalues are plotted vs. the transverse momentum $k_y$. The circles denote the amplitude of the wavefunctions on one of the two edges. We see that there are subgap edge modes that are unidirectional, i.e. chiral. The gapless bands localized on either edge cross at $k_y = 0$. There are two additional energy minima of the edge states near $k_y = \pm \pi$ that appear to be connected to the large second-nearest-neighbor pairing component $p_2$ which in turn enhances bulk gap minima. Note that the gapless chiral edge states are protected only up to the deep bulk gap minima, beyond which impurity scattering between the edge modes and the bulk continuum is allowed. The edge modes within the bulk gap minimum may be robust but such low-energy Bogoliubov-de Gennes quasiparticles are almost charge neutral. This severely reduces the robustness of the chiral edge state with respect to, e.g., edge disorder. Experimentally the detection of the edge current is not conclusive [15]. The subject on the cause of the missing edge current is controversial [35], and will be deferred to a later stage.

In thermodynamic quantities, the $\gamma$-band contributes significantly once the temperature is about the SC gap scale, in addition to the $\alpha$ and $\beta$ bands. In particular, while there are still open questions about the role of ($\alpha, \beta$) bands, the deep gap minima might contribute to explaining the power-law behavior in the specific heat at temperatures above the small gap scale [21]. This argument is similar in spirit to a recent discussion of possible anisotropic chiral $d$-wave superconductivity in sodium cobaltates [32].

Before closing we emphasize that the results for these qualitative features of the gap structure discussed so far are quite generic. (A detailed discussion of the general phase diagram of the given Fermiology beyond the specific ruthenate setting will be given elsewhere.) We obtain rather similar results for a considerable range of interactions, e.g., $(U, U', J, J') = (3.3, 1.1 \pm 0.1, 0.115, 0.115)$, where the only varying aspect we find from the data is the absolute instability scale of superconductivity. A general trend we find is that Hund’s rule coupling $J$ and pair hopping $J'$ favor large-q SDW

\[\begin{align*}
\text{(a)} & \quad \text{Energy spectrum vs. the transverse momentum $k_y$ in the $p + ip'$ SC state of a ribbon under open boundary conditions along $x$. Only the $d_{xy}$ orbitals are considered. The size of the circles denotes the wave function amplitude of the low-lying edge states integrated over two sites nearest to the respective edge.}
\text{(b) Relative pairing eigenvalues vs. spin-orbit coupling $\lambda$. The legend shows the five types of $d$-vector in the triplet pairing function. $S_{\text{min}}$ is the most attractive eigenvalue.}
\end{align*}\]
interactions, and if sufficiently large, would destabilize the \(p\)-wave pairing. Atomic spin-orbital coupling and interlayer hopping mixes \(d_{x^2}/d_{y^2}(n = 1, 2)\) and \(d_{xy}(n = 3)\) orbitals at the one-particle level, leading to inter-band proximity effect between the active and passive bands which may contribute to the pairing amplitude on the \(\alpha\) and \(\beta\) bands [36]. (At the level of interactions, the pair hopping \(J'\) is the only coupling which would allow for a proximity effect induced by the \(\gamma\)-band, and the effect is weak since this coupling is initially orthogonal to the \(p\)-wave channel.) To a good approximation, such effects can be included by using the above renormalized pairing interaction at a suitable energy scale \(\Lambda\) and continue the flow in the pairing channel alone, since the particle-hole channel is essentially saturated and decoupled from the pairing channel at this stage. The behavior of the eigenvalues of the pairing channel as a function of spin-orbit coupling strength \(\alpha_{\nu} = 2\epsilon_{\nu} a_{\nu} a_{\nu}^\dagger \psi_{\nu}^\dagger \psi_{\nu} \psi_{\nu} \psi_{\nu}^\dagger \) (where \(\epsilon\) is the antisymmetric tensor and \(\sigma\) is the Pauli matrix, and repeated orbital and spin labels are implicitly summed over) is shown in fig. 4(b). This leads us to the conclusion that the most favorable triplet pairing \(d\)-vector is \(d_{y^2} = (p_k \pm i\tau_3 k_z)\hat{z}\), as already found in previous works based on qualitative arguments. However, given the small splitting in the eigenvalues, an applied magnetic field could weaken the effect of spin-orbital coupling in favor of Majorana zero modes in vortices [12,13].

To conclude, we studied the pairing mechanism in \(Sr_2RuO_4\) within a three-orbital model by extensive FRG calculations that go beyond previous approaches for this system in that they take into account all three relevant bands and the competition between various interaction channels to arbitrary order in the bare couplings. Different FRG approaches we employ show the same trends in the dominant \(p\)-wave pairing and SDW channels: We find that the momentum \(Q\) of the dominant SDW interaction evolves from \(Q \sim (2/3, 2/3)\pi\) at high-energy scales to \(Q \sim (1/5, 1/5)\pi\) at low scales. The small-\(Q\) SDW fluctuations drive \(p\)-wave Cooper pairing predominantly on the \(\gamma\)-band derived from the \(d_{xy}\) orbital. The pairing receives contributions from first and second nearest neighbors on the \(Ru\) square lattice. The energetically most favorable combination of a \(p \pm ip'\) gap function has deep minima in amplitude on the \(\gamma\)-Fermi surface near \((\pi, 0)\) and \((0, \pi)\). This makes the chiral edge modes fragile already against a moderate amount of impurities.

***

We thank S. A. Kivelson, J. X. Li, M. Sigrist and S. Raghu for interesting discussions. This work was supported by NSFC (under grant No. 10974086, No. 11274269 and No. 11023002), the Ministry of Science and Technology of China (under grant No. 2011CBA00108 and No. 2011CB922101), DFG FOR 723, 912 and SPP 1458, and by the Swiss Nationalfonds, and the ERC starting grant TOPOELECTRICS ERC-StG-Thomale-StG-2013-336012.

REFERENCES

[1] Maeno Y. et al., Nature, 372 (1994) 532.
[2] Rice T. M. and Sigrist M., J. Phys.: Condens. Matter, 7 (1995) L643.
[3] Baskaran G., Physica B, 224 (1996) 490.
[4] Nelson K. D. et al., Science, 306 (2004) 1151.
[5] Ishida K. et al., Nature, 396 (1998) 658.
[6] Luke G. M. et al., Nature, 394 (1998) 558.
[7] Kapitulnik A., Xia J., Schiem E. and Palevski A., New J. Phys., 11 (2009) 055060.
[8] Mackenzie A. P. and Maeno Y., Rev. Mod. Phys., 75 (2003) 657.
[9] Bergemann C. et al., Adv. Phys., 52 (2003) 639.
[10] Maeno Y. et al., J. Phys. Soc. Jpn., 81 (2012) 011009.
[11] Callin C., Rep. Prog. Phys., 75 (2012) 042501.
[12] Ivanov D. A., Phys. Rev. Lett., 86 (2001) 268; Read N. and Green D., Phys. Rev. B, 61 (2000) 10267.
[13] Nayak C. et al., Rev. Mod. Phys., 80 (2008) 1083.
[14] Braden M. et al., Phys. Rev. B, 66 (2002) 064522.
[15] Kirtley J. R. et al., Phys. Rev. B, 76 (2007) 014526.
[16] Raghu S., Kapitulnik A. and Kivelson S. A., Phys. Rev. Lett., 105 (2010) 136401; Chong S. B., Raghu S., Kapitulnik A. and Kivelson S. A., Phys. Rev. B, 86 (2012) 064525; Raghu S., Chong S. B. and Lederer S., arXiv:1208.6344.
[17] Agterberg D. F., Rice T. M. and Sigrist M., Phys. Rev. Lett., 78 (1997) 3374.
[18] Hlubina R., Phys. Rev. B, 59 (1999) 9600.
[19] Nomura T. and Yamada K., J. Phys. Soc. Jpn., 69 (2000) 3675.
[20] Honerkamp C. and Salmoiffer M., Phys. Rev. Lett., 87 (2001) 187004; Honerkamp C. and Rice T. M., J. Low Temp. Phys., 131 (2003) 159.
[21] Nomura T. and Yamada K., J. Phys. Soc. Jpn., 71 (2002) 404.
[22] Huo J., Rice T. M. and Ziang F. C., Phys. Rev. Lett., 110 (2013) 167003.
[23] Nomura T. and Yamada K., J. Phys. Soc. Jpn., 71 (2002) 1993.
[24] Metzner W. et al., Rev. Mod. Phys., 84 (2012) 299.
[25] Damascelli A. et al., Phys. Rev. Lett., 85 (2000) 5194.
[26] Wang F., Zhai H., Ran Y., Vishwanath A. and Lee D. H., Phys. Rev. Lett., 102 (2009) 047005.
[27] Thomale R., Platt C., Hanke W. and Bernevig B. A., Phys. Rev. Lett., 106 (2011) 187003; Thomale R. et al., Phys. Rev. Lett., 107 (2011) 117001.
[28] Xiang Y. Y. et al., Phys. Rev. B, 86 (2012) 134508.
[29] Kiesel M. L. et al., Phys. Rev. B, 86 (2012) 020507.
[30] Wang W. S. et al., Phys. Rev. B, 85 (2012) 035144.
[31] Wang W. S. et al., Phys. Rev. B, 87 (2013) 115135.
[32] Kiesel M. L., Platt C., Hanke W. and Thomale R., Phys. Rev. Lett., 111 (2013) 097001.
[33] Xiang Y. Y., Wang W. S., Wang Q. H. and Lee D. H., Phys. Rev. B, 86 (2012) 024523.
[34] Husemann C. and Salmoiffer M., Phys. Rev. B, 79 (2009) 195125; Giering K. U. and Salmoiffer M., Phys. Rev. B, 86 (2012) 245122.
[35] See, e.g., Imai Y., Wakerayashi K. and Sigrist M., Phys. Rev. B, 85 (2012) 174532.
[36] Zhitomirsky M. E. and Rice T. M., Phys. Rev. Lett., 87 (2001) 057001.