Pairing symmetry of heavy fermion superconductivity in the two-dimensional Kondo-Heisenberg lattice model

Yu Liu\textsuperscript{1}, Guang-Ming Zhang\textsuperscript{2,3,*}, and Lu Yu\textsuperscript{1,3}

\textsuperscript{1}Institute of Physics, Chinese Academy of Sciences, Beijing 100190, China.
\textsuperscript{2}State Key Laboratory of Low-Dimensional Quantum Physics and Department of Physics, Tsinghua University, Beijing 100084, China;
\textsuperscript{3}Collaborative Innovation Center of Quantum Matter, Beijing, China and

\textsuperscript{*}Correspondence author. Email: gmzhang@tsinghua.edu.cn

(Dated: February 16, 2016)

In the two-dimensional Kondo-Heisenberg lattice model away from half-filled, the local antiferromagnetic exchange coupling can provide the pairing mechanism of quasiparticles via the Kondo screening effect, leading to the heavy fermion superconductivity. We find that the pairing symmetry strongly depends on the Fermi surface (FS) structure in the normal metallic state. When $J_H/J_K$ is very small, the FS is a small hole-like circle around the corner of the Brillouin zone, and the s-wave pairing symmetry has a lower ground state energy. For the intermediate coupling values of $J_H/J_K$, the extended s-wave pairing symmetry gives the favored ground state. However, when $J_H/J_K$ is larger than a critical value, the FS transforms into four small hole pockets crossing the boundary of the magnetic Brillouin zone, and the d-wave pairing symmetry becomes more favorable. In that regime, the resulting superconducting state is characterized by either nodal d-wave or nodeless d-wave state, depending on the conduction electron filling factor as well. A continuous phase transition exists between these two states. This result may be related to the phase transition of the nodal d-wave state to a fully gapped state, which is recently observed in Yb doped CeCoIn$_5$.

The study of heavy fermion intermetallic compounds has played an important role in our understanding of strongly correlated electron systems\cite{1, 2}, and the Kondo lattice model is believed to capture the low temperature physics\cite{3}. It has been established that the coherent superposition of individual Kondo screening clouds gives rise to a huge mass enhancement of quasiparticles, leading to a heavy Fermi liquid with a large Fermi surface comprising conduction electrons as well as local moments\cite{4, 5}. Competing with the Kondo singlet formation, the local moments indirectly interact with each other via the Ruderman-Kittel-Kasuya-Yosida (RKKY) interaction, and the antiferromagnetic (AFM) long-range ordered phase emerges in the small Kondo exchange coupling regime\cite{6–8}.

In addition, there has been growing evidence that local AFM correlation of the local moments can also induce the formation of Cooper pairs of heavy quasiparticles, leading to unconventional heavy fermion superconductivity. Due to the tiny energy gap, the direct experimental measurement of superconducting gap function is extremely difficult. Only recently, inelastic neutron scattering experiments on CeCu$_2$Si$_2$ have revealed a spin resonance peak, an indirect evidence of the nodal $d$-wave superconductivity\cite{9}. Thermal conductivity and heat capacity measurements on CeCoIn$_5$ also supported a superconducting gap with nodes along the diagonal directions of the Brillouin zone\cite{10, 11}. More recently, in scanning tunneling spectroscopy experiments the nodal points in the superconducting gap of CeCoIn$_5$ were found, as a more direct evidence of $d_{x^2-y^2}$ symmetry\cite{12, 13}.

At the same time, an s-wave superconductivity has also been confirmed in CeRu$_2$ and CeCo$_2$ in nuclear quadrupole resonance measurements\cite{14, 15}, where the spin-lattice relaxation rate exhibits an exponential decay at low temperatures and the Hebel-Slichter peak is observed. Furthermore, the laser photoemission spectroscopy measurements on CeRu$_2$ have demonstrated a clear anisotropic s-wave superconducting gap at the Fermi level\cite{16}. Since the Coulomb repulsion is dominant in heavy fermion systems, these results cast doubts on whether this fully gapped superconducting state can be understood within the pairing mechanism of the local AFM correlation.

From the theoretical side, to simplify the RKKY interaction, one can explicitly introduce the local AFM Heisenberg superexchange interaction $J_H$ among the local moments into the Kondo lattice system\cite{20–26}. When the local moments are expressed in terms of fermionic spinons, the large-$N$ fermionic approach\cite{17–19} can be used to treat the Kondo-Heisenberg lattice model very efficiently in the limit of $J_K > J_H$, after a spinon hopping order parameter is introduced\cite{20, 21, 24, 26}. Then an effective hybridization between the conduction electron and spinon bands leads to a paramagnetic heavy Fermi liquid. Based on the reconstructed large Fermi surface (FS), the instabilities of AFM order and unconventional superconductivity can be further analyzed. As long as the AFM long range order is suppressed, the spinon singlet pairings can be also promoted from the local AFM spin exchange, further reducing the ground state energy. Via the Kondo screening effect, the Cooper pairs of the conduction electrons can be induced, leading to heavy fermion superconductivity. Some numerical evidence on robust d-wave pairings has been observed in the two-dimensional Kondo-Heisenberg lattice systems\cite{27–29}.

In this paper, we develop an effective mean field (MF) theory on heavy fermion superconductivity in the Kondo-Heisenberg lattice model on a two-dimensional square lattice. In the paramagnetic Fermi liquid phase ($J_K > J_H$), we first notice that the FS topology undergoes a dramatic change as the ratio of $x = J_H/J_K$ is gradually increased. Since the
spinon singlet pairing from the local AFM exchange coupling has a form factor, which is either an extended s-wave or d-wave symmetry, we find that the superconducting pairing symmetry depends on the FS structure. In the presence of the spinon pairings, we have to introduce the local pairing order parameter between the conduction electrons and spinons in the Kondo spin exchange interaction. When $x$ is very small, the FS is a small hole-like circle around the corner of the first Brillouin zone. We find that the s-wave superconducting state has a lower ground state energy[30]. For the intermediate coupling values of $J_H/J_K$, the extended s-wave pairing symmetry gives the favored ground state. However, as $x$ is larger than a critical value, the corresponding FS consists of four hole pockets crossing the boundary of the magnetic Brillouin zone, and the d-wave pairing symmetry is more favorable. In this regime, the resulting superconducting state can be a nodal d-wave or nodeless d-wave state, depending on the conduction electron filling factor as well. This result may be used to explain the recent experimental observation in the Yb doped CeCoIn$_5$ (Ref.[31]).

The model Hamiltonian of the Kondo-Heisenberg lattice model is defined by:

$$ H = \sum_{k,\sigma} \epsilon_k c_{k\sigma}^\dagger c_{k\sigma} + J_K \sum_i \mathbf{S}_i \cdot \mathbf{s}_i + J_H \sum_{(ij)} \mathbf{S}_i \cdot \mathbf{S}_j, \quad (1) $$

where $\epsilon_k$ denotes the conduction electron band, the local moments have the fermionic representation $\mathbf{S}_i = \frac{1}{2} \sum_{\sigma,\sigma'} f_{i\sigma}^\dagger \tau_{\sigma\sigma'} f_{i\sigma}$ and $\tau$ is the Pauli matrices. There is a local constraint: $\sum_{\sigma} f_{i\sigma} f_{i\sigma} = 1$ to restrict any charge fluctuations, and the $f$-fermions only describe the spin degrees of freedom of the local moments, which will be referred to as spinons.

Following the large-$N$ fermionic approach[20, 22], the Kondo spin exchange and Heisenberg superexchange terms can be simply expressed as

$$ \mathbf{S}_i \cdot \mathbf{s}_i = -\frac{1}{2} \left( f_{i\uparrow} c_{i\uparrow}^\dagger + f_{i\downarrow} c_{i\downarrow}^\dagger \right) \left( c_{i\uparrow} f_{i\uparrow}^\dagger + c_{i\downarrow} f_{i\downarrow}^\dagger \right), $$

$$ \mathbf{S}_i \cdot \mathbf{S}_j = -\frac{1}{2} \left( f_{i\uparrow} f_{j\uparrow}^\dagger + f_{i\downarrow} f_{j\downarrow}^\dagger \right) \left( f_{i\uparrow}^\dagger f_{j\uparrow} + f_{i\downarrow}^\dagger f_{j\downarrow} \right), \quad (2) $$

where constant terms have been neglected. Then a Kondo screening and a uniform short-range AFM order parameters can be introduced as

$$ V = -\sum_{\sigma} \left\langle c_{i\sigma}^\dagger f_{i\sigma} \right\rangle, \chi = -\sum_{\sigma} \left\langle f_{i\sigma}^\dagger f_{i\sigma} \right\rangle. \quad (3) $$

To avoid the accidental degeneracy of the conduction electrons on a square lattice, we have

$$ \epsilon_k = -2t \left( \cos k_x + \cos k_y \right) + 4t' \cos k_x \cos k_y - \mu, \quad (4) $$

where $t$ and $t'$ are the first and second nearest neighbor hopping parameters, and a chemical potential of the conduction electrons $\mu$ has been introduced as a Lagrangian multiplier to fix the density of the conduction electrons $n_c$. Under the uniform MF approximation, the spinons also form a very narrow band with the dispersion $\chi_k = J_H \chi \left( \cos k_x + \cos k_y \right) + \lambda$

where $\lambda$ is also a Lagrangian multiplier to impose the local constraint on average.

Then the MF Hamiltonian for the heavy Fermi liquid phase reads

$$ H = \sum_{k\sigma} \left( c_{k\sigma}^\dagger f_{k\sigma}^\dagger \right) \left( \frac{1}{2} J_K V \chi \right) \left( c_{k\sigma} f_{k\sigma}^\dagger \right) + E_0, \quad (5) $$

with $E_0 = N \left( J_H \chi^2 + J_K V^2/2 - \lambda + \mu n_c \right)$. The quasiparticle excitation spectra can be easily obtained

$$ E_k^{(\pm)} = \frac{1}{2} \left[ (\epsilon_k + \chi_k) \pm \sqrt{(\epsilon_k - \chi_k)^2 + (J_K V)^2} \right]. \quad (6) $$

Then the self-consistent equations for $\chi$, $V$, $\lambda$ and $\mu$ can be obtained by minimizing the ground state energy. We first perform a numerical calculation, choosing the following parameters $t' = 0.3t$, $J_K = 2t$, and $n_c = 0.8$. The following results are obtained: as the local AFM interaction $J_H$ gradually increases, the lower branch of the heavy quasiparticle spectrum $E_k^{(-)}$ is calculated and the corresponding FS topology obtained and displayed in Fig.1.

For a very small local AFM Heisenberg spin exchange, i.e. $J_H/t \leq 0.355$, we can clearly see that the FS is a hole-like circle around the corner of the first Brillouin zone, which corresponds to the large electronic FS of the heavy quasiparticles. At $J_H/t = 0.356$, the topology of the FS starts to change: a small circle emerges in the center of the deformed hole FS. As $J_H/t$ is further increased, both circles expand and the small one is deformed into a rotated square. Up to $J_H/t = 0.481$, the two deformed circles intersect each other and then decompose into four kidney-like Fermi pockets. When $J_H/t$ continues to increase, the resulting FS is shifted inward along the diagonal direction. The detailed discussion had been presented in our previous publication[26]. At $J_H/t = 1.0556$, the Fermi pockets are close to the momentum $(\pi/2, \pi/2)$ and its equivalent points. Actually such a FS structure starts to cross the boundary of the magnetic Brillouin zone. Quantum Monte Carlo cluster approach has been used to study the evolution of the FS close to the magnetic transition in the two-
dimensional Kondo lattice system[5], where the heavy quasiparticle bands drop below the FS giving rise to hole pockets around \( k = (\pi/2, \pi/2) \) and equivalent points. Our obtained FS structure is consistent with this numerical result.

To further consider the instability of the heavy Fermi liquid state, we should notice that the local AFM Heisenberg superexchange can also be written in terms of the spinon singlet pairs up to a constant

\[
S_i \cdot S_j = -\frac{1}{2} (f^\dagger_{ij} f^\dagger_{jk} - f^\dagger_{ki} f^\dagger_{kj}) (f_{ij} f_k - f_{jk} f_{i}^\dagger).
\]

(7)

Actually, Coleman and Andrei[20] had emphasized that the local SU(2) gauge invariance of the local Heisenberg spin operator generally requires the consideration of both spinon hopping and pairing order parameters. Further arguments can be made by using the symplectic representation of the local magnetic spins[32]. Then the spinon pairing parameter is introduced as

\[
\Delta_{ij} = -\langle f^\dagger_{ij} f^\dagger_{jk} - f^\dagger_{ki} f^\dagger_{kj} \rangle.
\]

(8)

For a two-dimensional square lattice model, the spinon pairing order parameter \( \Delta_{ij} \) has a form factor with either an extended s-wave or the d-wave symmetry[33]. Both the extended s-wave and d-wave form factors in the momentum space have sign change in the Brillouin zone and are shown in Fig.2.

When the spinon pairings are present, the local pairing order parameter between the conduction electrons and spinons has to be introduced, because the Kondo spin exchange interaction can also be expressed as

\[
S_i \cdot s_i = -\frac{1}{2} (c^\dagger_{i} c^\dagger_{j} - c^\dagger_{j} c^\dagger_{i}) (f_{ij} c^\dagger_{i} - f^\dagger_{ji} c^\dagger_{i}).
\]

(9)

Then a local s-wave pairing order parameter is defined by

\[
\Delta_{c-f} = -\langle c^\dagger_{i} c^\dagger_{j} - c^\dagger_{j} c^\dagger_{i} \rangle,
\]

(10)

so the MF model Hamiltonian in momentum space can be written in a compact form

\[
\mathcal{H}_{mf} = \sum_k \psi^\dagger_k \begin{pmatrix}
\epsilon_k & 0 & J_K V & J_K \Delta_{cf} \\
0 & -\epsilon_k & J_K \Delta_{cf} & J_K V \\
J_K \Delta_{cf} & J_K \Delta_{cf} & \chi_k & J_H \Delta_k \\
J_K V & J_K V & J_H \Delta_k & -\chi_k
\end{pmatrix} \psi_k + \sum_k (\epsilon_k + \chi_k) + E_0,
\]

(11)

where a Nambu spinor has been defined as

\[
\psi^\dagger_k = \begin{pmatrix} c_{k\uparrow} & c_{k\downarrow} \end{pmatrix}, \quad E_0 = N \left[ J_K (V^2 + \Delta_{cf}^2) / 2 + J_H (\Delta_k^2 + \chi^2) - \lambda - \mu n_c \right].
\]

The spinon pairing gap function is chosen as

\[
\Delta_k = \Delta_0 (\cos k_x \pm \cos k_y),
\]

(12)

for extended s-wave and d-wave pairing, respectively. Diagonalizing this MF model Hamiltonian, two quasiparticle bands are derived

\[
E^\pm_k = \sqrt{E_{k1} + \sqrt{E_{k1}^2 - E_{k2}^2}},
\]

\[
E_{k1} \equiv \frac{1}{2} (\epsilon_k + \chi_k + J_H \Delta_k^2) + J_K^2 (V^2 + \Delta_{cf}^2) / 4,
\]

\[
E_{k2} \equiv \sqrt{\left( \epsilon_k \chi_k - J_K \left( V^2 - \Delta_{cf}^2 \right) / 4 \right)^2 + (\epsilon_k J_H \Delta_k - J_K^2 V \Delta_{cf} / 2)^2}.
\]

Due to the particle-hole symmetry of the superconducting quasiparticles, all negative energy states are filled up in the ground state, and the ground state energy density is thus obtained

\[
E_g = \frac{\sqrt{2}}{N} \sum_k \sqrt{E_{k1} + E_{k2}}
\]

\[+ \frac{J_K}{2} \left( V^2 + \Delta_{cf}^2 \right) + J_H \Delta_k^2 + J_H \chi^2 + \mu (\nu_c - 1). \]

The saddle point equations for the MF order parameters \( \chi, \Delta, \Delta_{cf} \) and \( \lambda \) can also be determined by minimizing the ground state energy. The chemical potential \( \mu \) is still determined by the conduction electron density \( \nu_c \). It should be emphasized that the obtained MF order parameters \( V, \chi \), and \( \lambda \) are different from those values in the normal paramagnetic phase. In particular, the difference of the chemical potential \( \mu \) modifies the position of the Fermi energy as well as the FS structure.

Although there are no direct attractions among the conduction electrons, the spinon singlet pairings and the pairings of the conduction electrons and spinons provide an indirect glue for the formation of the Cooper pairs of conduction electrons via the Kondo screening/hybridizing effect. So the resulting ground state represents a heavy fermion superconducting state. With the help of the double-time retarded Green function, the Cooper pairing order parameter of the conduction electrons can also be deduced

\[
\langle c^\dagger_{k\uparrow} c^\dagger_{-k\downarrow} \rangle = \frac{J_H J_K^2 \Delta_k \left( \Delta_{cf}^2 - V^2 \right) + 2 J_K^2 \chi_k V \Delta_{cf}}{8 E_{k1}^2 \sqrt{2 (E_{k1} + E_{k2})}}.
\]

(13)

We find that the coexisting cases \( (\Delta_{cf} \Delta_0 \neq 0) \) of the local electron-spinon and spinon-spinon pairing are unstable. Therefore, only three cases of the s-wave \( (\Delta_{cf} \neq 0) \), the extended s-wave and d-wave pairing symmetries are discussed in the following. When we choose \( t' = 0.3t, \nu_c = 0.8 \) and \( J_K = 2t \), the MF self-consistent equations are numerically solved, respectively. The obtained pairing strengths are displayed in Fig.3a. When \( 0 < J_H / t < 1.02 \), the s-wave pairing strength is the largest. For \( 1.02 < J_H / t < 1.36 \),
the system. However, when matches the FS structure of the heavy quasiparticles, and the 1 and d-wave symmetric superconducting states can further save the local AFM spin exchange interaction transition are expected in the superconducting phase. When 0 < J_H/t < 0.886, the extended s-wave pairing symmetry has a relatively lower energy in the range 0.886 < J_H/t < 1.056, and the d-wave pairing state becomes favorable for J_H/t > 1.056. The corresponding results are displayed in Fig.3b. As a reference, the ground state energy of the heavy Fermi liquid phase has been subtracted to obtain the condensation energies. As we expected, the s-wave, extended s-wave, and d-wave symmetric superconducting states can further save the ground state energy.

According to the above results, two discontinuous phase transition are expected in the superconducting phase. When the local AFM spin exchange interaction J_H/t < 0.886, the s-wave pairing state is the ground state. For 0.886 < J_H/t < 1.056, the pairing form factor with s-wave pairing symmetry matches the FS structure of the heavy quasiparticles, and the extended s-wave superconducting state is the ground state of the system. However, when J_H/t > 1.055, the pairing form factor with d-wave symmetry matches the FS, and the d-wave superconducting state becomes more favorable. Therefore, due to the presence of the FS deformation, the resulting superconducting pairing symmetry is very sensitive to the local AFM spin exchange.

Moreover, as the filling factor of the conduction electrons is increased, another continuous transition can be exhibited from the lower superconducting quasiparticle excitation spectrum. The d-wave symmetric superconducting state shows a phase transition from the nodal d-wave to the nodeless d-wave states. This is a very unusual phenomenon in the superconducting phase. When the Fermi level of the normal phase is very close to the top band edge, the d-wave pairing amplitude becomes larger than the bandwidth of the heavy quasiholes, and the superconducting pairing state actually belongs to the strong pairing regime. Then the nodes in the superconducting quasiparticle excitation spectrum are no longer protected by the d-wave symmetry. In Fig.4, we show that the lower superconducting quasiparticle excitation spectrum with t'/t = 0.3, J_H/t = 1.42 and J_K/t = 2.0 for different conduction electron filling factor. For the cases of n_c = 0.8 and 0.829, we can see the nodes around (π/2, π/2) and its equivalent points, while for n_c = 0.87 and 0.9, a small gap opens up. A general ground state phase diagram has been summarized in Fig.5.

In conclusion, we have presented an effective MF theory for heavy fermion superconductivity in the two-dimensional...
Kondo lattice model with the local AFM Heisenberg exchange coupling among the local moments. We would like to emphasize that, it is the local AFM short-range interaction that can deform the FS structure of the heavy quasiparticles and induce unconventional superconducting long-range ordered states. Due to the presence of spinon singlet pairing, the Cooper pairs among the conduction electrons are induced via the Kondo screening effect. However, the pairing symmetry depends on the FS topology of heavy quasiparticles in the normal state, which is determined by the strength of the local AFM spin exchange and the conduction electron filling factor. When the local AFM Heisenberg exchange coupling is weak, the heavy Fermi liquid phase has a small hole Fermi pocket around the corner of the Brillouin zone, and the s-wave and the extended s-wave form factor of the local AFM interaction matches the FS, so the s-wave superconductivity is more favorable. Such a situation occurs in the materials of CeRu$_2$ and CeCoIn$_5$, as supported by various experimental measurements\[14–16\]. Since these materials show a heavy Fermi liquid behavior in the normal state, the local AFM spin exchange interaction is expected to be very weak but can still provide an attractive glue among the fermionic spinons.

When the local AFM Heisenberg exchange coupling is stronger than a critical value and the conduction electron filling factor is not close to the half-filled limit, the corresponding FS of the heavy Fermi liquid changes into four small hole pockets around the momentum ($\pi/2$, $\pi/2$) and its equivalent points. Such a FS structure more favors the d-wave form factor of the local AFM correlation. So the nodal d-wave superconducting state can be obtained, and such a situation happens in the materials of CeCu$_2$Si$_2$ and CeCoIn$_5$, with support from various experimental measurements\[9, 10, 12\].

More recently, some interesting experimental results have been discovered upon Yb doping in CeCoIn$_5$. The superconductivity is extremely stable and the transition temperature linearly depends on the Yb doping concentration\[34\], and the temperature dependence of the London penetration depth indicates that the nodal d-wave superconductivity changes into a fully gapped state after a critical Yb doping\[31\]. The transport measurements\[35\] have indicated that the large amount of Yb doping introduces holes into the system. This is equivalent to increasing the conduction electron filling factor in our theory, so the transition observed experimentally just corresponds to the continuous phase transition from the nodal d-wave state to the nodeless d-wave state. In order to fully explain this transition, further theoretical work including the estimation of the fluctuations around the MF solution should be considered.

The authors would like to thank D. H. Lee and T. Xiang for stimulating discussions. This work was supported by the National Natural Science Foundation of China (Grant Nos. 20121302227, 11120101003, 11121063) and by China Postdoctoral Science Foundation (Grant No. 2013M541069).

[1] Stewart G 2001 Revs. Mod. Phys. 73 797
[2] Lohneysen H V, Rosch A, Vojta M and Wölfle P 2007 Revs. Mod. Phys. 79 1015
[3] Si Q, Rabello S, Ingersent K and Smith J 2001 Nature 413 804
[4] Watanabe H and Ogata M 2007 Phys. Rev. Lett. 99 136401
[5] Martin L C and Assaad F F 2008 Phys. Rev. Lett. 101 066404
[6] Domiach S 1977 Physica B & C 91 231
[7] Lacroix C and Cyrot M 1979 Phys. Rev. B 20 1969
[8] Zhang G M, Gu Q and Yu L 2000 Phys. Rev. B 62 69
[9] Zhong G M and Yu L 2000 Phys. Rev. B 62 67
[10] Stockert O et al 2011 Nature Phys. 7 119.
[11] Izawa K, Yamaguchi H, Matsuda Y, Shishido H, Settai R and Onuki Y 2001 Phys. Rev. Lett. 87 057002
[12] Aoki H et al 2004 J. Phys. Condensed Matter 16 L13
[13] Allan M P et al 2013 Nature Phys. 9 468
[14] Zhou B B et al 2011 Phys. Rev. Lett. 101 057001
[15] Read N and Newns D M 1983 J. Phys. C: Solid State Phys. 16 3273
[16] Auerbach A and Levin K, 1986 Phys. Rev. Lett. 57 877
[17] Mills A J and Lee P A 1987 Phys. Rev. B 35 3394
[18] Coleman P and Andrei N 1989 J. Phys.: Condens. Matter 1 4057
[19] Read N and Newns D M 1983 Phys. Rev. Lett. 62 595
[20] Grover T and Senthil T 2010 Phys. Rev. B 81 205102
[21] Zhuang G M, Su Y H and Yu L 2011 Phys. Rev. B 83 033102
[22] Sato N K et al 2001 Nature Phys. 1 410 340
[23] Xavier J C and Dagotto E 2008 Phys. Rev. Lett. 100 146403
[24] Bodensiek O, Zitko R, Vojta M, Jarrell M, and Pruschke T 2013
[25] Grover T and Senthil T 2010 Phys. Rev. B 81 205102
[26] Sato N K et al 2001 Phys. Rev. Lett. 87 057002
[27] Aoki H et al 2004 J. Phys. Condensed Matter 16 L13
[28] Allan M P et al 2013 Nature Phys. 9 468
[29] Zhou B B et al 2011 Phys. Rev. Lett. 101 057001
[30] Read N and Newns D M 1983 J. Phys. C: Solid State Phys. 16 3273
[31] Auerbach A and Levin K, 1986 Phys. Rev. Lett. 57 877
[32] Mills A J and Lee P A 1987 Phys. Rev. B 35 3394
[33] Coleman P and Andrei N 1989 J. Phys.: Condens. Matter 1 4057
[34] Read N and Newns D M 1983 Phys. Rev. Lett. 62 595
[35] Grover T and Senthil T 2010 Phys. Rev. B 81 205102
[36] Zhuang G M, Su Y H and Yu L 2011 Phys. Rev. B 83 033102
[37] Sato N K et al 2001 Nature Phys. 1 410 340
[38] Xavier J C and Dagotto E 2008 Phys. Rev. Lett. 100 146403
[39] Bodensiek O, Zitko R, Vojta M, Jarrell M, and Pruschke T 2013
[40] Read N and Newns D M 1983 J. Phys. C: Solid State Phys. 16 3273
[41] Auerbach A and Levin K, 1986 Phys. Rev. Lett. 57 877
[42] Mills A J and Lee P A 1987 Phys. Rev. B 35 3394
[43] Coleman P and Andrei N 1989 J. Phys.: Condens. Matter 1 4057
[44] Read N and Newns D M 1983 Phys. Rev. Lett. 62 595
[45] Grover T and Senthil T 2010 Phys. Rev. B 81 205102
[46] Zhuang G M, Su Y H and Yu L 2011 Phys. Rev. B 83 033102
[47] Sato N K et al 2001 Nature Phys. 1 410 340
[48] Xavier J C and Dagotto E 2008 Phys. Rev. Lett. 100 146403
[49] Bodensiek O, Zitko R, Vojta M, Jarrell M, and Pruschke T 2013
[50] Read N and Newns D M 1983 J. Phys. C: Solid State Phys. 16 3273
[51] Auerbach A and Levin K, 1986 Phys. Rev. Lett. 57 877
[52] Mills A J and Lee P A 1987 Phys. Rev. B 35 3394
[53] Coleman P and Andrei N 1989 J. Phys.: Condens. Matter 1 4057
[54] Read N and Newns D M 1983 Phys. Rev. Lett. 62 595
[55] Grover T and Senthil T 2010 Phys. Rev. B 81 205102
[56] Zhuang G M, Su Y H and Yu L 2011 Phys. Rev. B 83 033102
[57] Sato N K et al 2001 Nature Phys. 1 410 340
[58] Xavier J C and Dagotto E 2008 Phys. Rev. Lett. 100 146403
[59] Bodensiek O, Zitko R, Vojta M, Jarrell M, and Pruschke T 2013
[60] Read N and Newns D M 1983 J. Phys. C: Solid State Phys. 16 3273
[61] Auerbach A and Levin K, 1986 Phys. Rev. Lett. 57 877
[62] Mills A J and Lee P A 1987 Phys. Rev. B 35 3394
[63] Coleman P and Andrei N 1989 J. Phys.: Condens. Matter 1 4057