**Distinct reduction of Knight shift in superconducting state of Sr$_2$RuO$_4$ under uniaxial strain**

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(Rceived 4 December 2019; revised 11 May 2020; accepted 22 July 2020; published 31 August 2020)

 Shortly after the discovery of superconductivity in Sr$_2$RuO$_4$, spin-triplet pairing was proposed and further corroborated by a constant Knight shift ($K$) across the transition temperature ($T_c$). However, a recent experiment observed a drop in $K$ at $T_c$ which becomes larger under uniaxial strain, ruling out several spin-triplet scenarios. Here we show that even parity interorbital spin-triplet pairing can feature a $d$ vector that rotates when uniaxial strain is applied, leading to a larger drop in the spin polarization perpendicular to the strain direction, distinct from spin-singlet pairing. We propose that anisotropic spin polarization under strain will ultimately differentiate triplet versus singlet pairing.

DOI: 10.1103/PhysRevResearch.2.032055

**Introduction.** The discovery of superconductivity in Sr$_2$RuO$_4$ [1] has had great attention over the past two decades. It has been considered the best solid-state system which exhibits a time-reversal symmetry breaking $p$-wave spin-triplet pairing analog of the A phase in $^3$He [2]. The microscopic Hamiltonian of nuclear magnetic resonance (NMR) and spin-polarization drop should be smaller. For a singlet, the two spin polarizations parallel versus perpendicular to the strain direction. When the strain is applied along the $a$ axis, the $d$ vector rotates towards the $b$ axis as shown by the red arrows in Fig. 1, leading to a larger drop in the spin-polarization along the $b$ axis than that of the unstrained case, as reported in Ref. [12]. With the same strain condition, the $a$-axis polarization drop should be smaller. For a singlet, the two spin polarizations are the same. Thus we propose a NMR Knight shift experiment with the reasonably large field (but below the 1.5 T upper critical field) along the $a$ axis under the $a$-axis strain, to be compared with the $b$-axis polarization. This will ultimately differentiate spin-triplet versus -singlet pairings.

Below we formulate the proposed idea using a model which consists of a Kanamori interaction and a $t_{2g}$ tight binding model with SOC. While the atomic SOC leading to an s-wave gap is used for clarity, it can be generalized by including momentum dependent SOC terms leading to any even-parity OSST pairing (such as $d$-wave or $g$-wave).

**Microscopic Hamiltonian.** Sr$_2$RuO$_4$ is a multiorbital system with non-negligible SOC. The orbital degrees of freedom allow for four distinct pairings which satisfy the antisymmetric fermion wave function requirement, i.e., $\Delta(k) = -\Delta(\varepsilon,-k)$. The four types are: (i) even-parity intraorbital (or interorbital-triplet) spin singlet ($\phi_x$, $\phi_y$), (ii) odd-parity

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FIG. 1. The red arrows at representative momenta show the $d$ vector of inter-orbital spin triplet pairing for (a) unstrained and (b) uniaxial-strain along $a$ axis. This transforms to intraband pseudospin-singlet pairing on the FS and inter-band pairings (see the main text for details). The $d$ vector rotation occurs the most in the diagonal direction of the Brillouin zone. The length of each arrow represents the in-plane component; the shorter the arrow, the bigger the $c$-axis component. Note that the arrows with an inverted tail correspond to a vector primarily along the $c$ axis. The blue color on the FS denotes the size of gap. The red arrow at the bottom corner of each panel represents the averaged $d$ vector direction projected onto the $ab$-plane denoted by $\theta$; $\theta = 45^\circ$ and $63^\circ$ in (a) and (b), respectively.

interorbital-singlet spin singlet, (iii) odd-parity intraorbital (or interorbital-triplet) spin triplet ($\tilde{\text{D}}_i$), and (iv) even-parity OSST ($\tilde{\text{D}}_s$), where $\nu$ represents interorbital, and $\alpha$, intraorbital pairings among $t_{2g}$ [13].

A generic Hamiltonian $H = H_{\text{kin}} + H_{\text{SOC}} + H_{\text{int}}$ consisting of a tight binding model, SOC, and Kanamori interaction is considered. The tight binding and SOC terms are used to reproduce the Fermi surface (FS) reported in Ref. [14], and are listed in Ref. [15]. The underlying FS of three bands, $\alpha$, $\beta$, and $\gamma$ is reported earlier [16–18], and was further refined in Ref. [14] shown as the solid lines in Fig. 1. The interaction term is given by

$$H_{\text{int}} = \sum_{i,a} U^{i} \epsilon_{ia} \epsilon_{ia}^{\dagger} + \sum_{i,a \neq b} V^{i} \epsilon_{ia}^{\dagger} \epsilon_{ia} \epsilon_{ib}^{\dagger} \epsilon_{ib},$$

$$+ \frac{J_H}{2} \sum_{i,a \neq b} \epsilon_{ia}^{\dagger} \epsilon_{ib}^{\dagger} \epsilon_{ia} \epsilon_{ib} + \frac{J_H}{2} \sum_{i,a \neq b} \epsilon_{ia}^{\dagger} \epsilon_{ia} \epsilon_{ib}^{\dagger} \epsilon_{ib},$$

(1)

with Hubbard interaction, $U$, and Hund’s coupling, $J_H$, where $V = U - 2J_H$, and where $a$ and $b$ represent the $t_{2g}$ orbitals ($yz, xz, xy$). This can be expressed in terms of pairing order parameters, including the OSST parameters, which appear as

$$\frac{H_{\text{eff}}}{2N} = (V - J_H) \sum_{\nu} \tilde{\text{D}}^{\nu}_{i}(q) \cdot \tilde{\text{D}}_{\nu}(q),$$

(2)

where $\tilde{\text{D}}^{\nu}_{i}(q)$ is given by

$$\tilde{\text{D}}^{\nu}_{i}(q) = \frac{1}{4N} \sum_{k} \epsilon_{i\alpha}^{\dagger} [i\sigma^i \sigma^l]_{\nu \nu} \left[ \hat{\lambda}^{\nu}_{i\alpha} \right]_{ab} c_{-k+q \nu \alpha} \epsilon_{-k+q \nu \alpha},$$

(3)

with $l = x, y, z$, \(\hat{\lambda}^{\nu}_{i\alpha}\) are $3 \times 3$ antisymmetric matrices in the orbital basis under the exchange of the $t_{2g}$ orbitals for three different interorbital matrices denoted with $\nu = X$ (between $xz$ and $xy$ orbitals), $Y$ ($yz$ and $xy$), and $Z$ ($xz$ and $yz$). Their expressions are given in Ref. [15]. The full form of the interaction written in terms of pairing order parameters, including induced intra orbital spin singlets $\phi_\nu$, and interorbital-triplet spin singlets, both of which appear with repulsive interactions, are also given in Ref. [15]. The OSST channel has an attractive interaction for $3J_H > U$, and while this is larger than most values of Hund’s coupling in $4d$ transition metals, where $J_H$ is about $20\%–30\%$ of $U$ [19], recent studies going beyond mean-field theory support OSST pairing originating from Hund’s coupling without the strict condition of $3J_H > U$ [20–23].

The direction of the $d$ vector is determined by the SOC [13,24], with order parameters belonging to the $A_{1g}$ representation for atomic SOC [25–27]. The importance of the SOC in Sr$_2$RuO$_4$ was addressed earlier [13,28–31], and recently re-emphasized [14].

**Pairing gap under strain.** Since the OSST pairing corresponds to pairing between orbitals with different energies at $\mathbf{k}$ and $-\mathbf{k}$, we consider the possibility of finite momentum pairing, i.e., FFLO state. Using a self-consistent mean-field theory, we find the zero-momentum $\mathbf{q} = 0$ state is always the lowest state despite the pairing between different orbitals. However, the pairing amplitude appears to be extremely small as shown in Fig. 2 with the magnitudes of the $\text{D}_i$ and induced intraband spin singlets, $\phi_\nu$. They are thousands of times smaller than the $t_{2g}$ bandwidth, even though the attractive interaction is reasonably large. We set $3J_H - U = 0.5$ for the current results, and the mean-field theory in general overestimates the gap size. The interorbital pairing would appear to require a finite $\mathbf{q}$ value to produce a gap on the FS without orbital hybridization or SOC. However, when the atomic SOC is finite the OSST pairing probed onto the band basis transforms into *intraband* pseudospin-singlet pairing on the FS, denoted by $\tilde{\text{D}}_i$, where $i = \alpha, \beta$, and $\gamma$ in the quasiparticle dispersion shown in Fig. S1. The quasiparticle dispersion represents strongly anisotropic gaps, which are very small in size, both at and below the FS. This suggests that when the bandwidth is renormalized by electronic correlations, and becomes narrower, the OSST is further favoured. A recent dynamical mean-field theory reported a strong mass renormalization of the bands [32], which would also enhance the OSST pairing.

To study the uniaxial strain effects, we change the ratio of the hopping integrals along the $a$ and $b$ axes such that
\( t_{ij} = (1 - \delta) t_{ij} \) and \( t_{ij} = (1 + \delta) t_{ij} \) for \( j = 1, 2, 3 \). Uniaxial strain along the \( a \) axis corresponds to \( \delta < 0 \). The change of different order parameters as a function of \( \delta \) is shown in Fig. 2. The pairing gap is roughly quadratic in \( \delta \) as expected from the even parity pairing. While \( D_X^\alpha \) and \( D_Y^\alpha \) exhibit opposite behavior under strain, these \( A_{1g} \) solutions do not exhibit a split transition under strain [25]. When the \( \gamma \) band touches the vHS around \( \delta = \pm 0.07 \), the pairing amplitude is peaked. Since mean-field theory causes the gap to be proportional to the transition temperature, \( T_c \) is also peaked as reported in Refs. [33,34]. The overall gap size is minuscule in comparison to the energy scale of the kinetic and potential terms as discussed above.

**Rotation of the \( d \) vector under uniaxial strain.** For spin-triplet pairing, the \( d \) vector represents the direction along which the spin projection of the condensed pair has eigenvalue zero [3]. When SOC is finite, the mean-field solutions find the pinning of the \( d \) vector depending on the interorbital composition via SOC. For the pairing between \( xz \) and \( xy \) orbitals, the \( d \) vector points along the \( x \) direction (represented by \( D_X^\alpha \)), \( yz \) and \( xy \) along the \( y \) direction (\( D_Y^\alpha \)), and \( xz \) and \( yz \) along the \( z \) axis (\( D_Z^\alpha \)). The \( x, y, \) and \( z \) axes are the same as the crystallographic axes of \( a, b, \) and \( c, \) as \( \text{Sr}_3\text{RuO}_4 \) is a tetragonal lattice. The \( d \) vector changes in momentum space as shown in Fig. 1(a), as the orbital composition changes along the FS. The red arrows represent the \( d \) vector directions. The shorter the length of arrow, the bigger the \( c \)-axis component of the \( d \) vector. There is a finite \( d \) vector at every momentum point, and on average it is finite in all directions leading to a reduction of the spin polarization in all directions.

In the absence of strain, due to the tetragonal symmetry, there is a \( \pi \) rotational symmetry between \( \hat{D}^\alpha \) and \( \hat{D}^\beta \). This leads to the same reduction of the spin polarization along the \( a \) and \( b \) axes (and any other directions related to the symmetry of the tetragonal lattice). However, when the uniaxial strain is applied, the orbital composition changes mainly around \( X \) and \( Y \) regions of the Brillouin Zone (BZ) as shown by the underlying FS in Fig. 1(b). Most importantly, the \( yz \) orbital contribution to all bands increases, causing the \( d \) vector at every momentum to rotate towards the \( b \)-axis, with the most change occurring around the diagonal direction of the BZ. This will then affect the magnitude of the spin polarization in the superconducting state, and generates a directional dependence, which we show below.

**Spin polarization under strain.** The magnetic susceptibility \( \chi_{ij} \) measured by the NMR Knight shift is given by \( \delta M_j / \delta B_j \) where \( M \) is the magnetization, \( B \) is an external magnetic field, and \( j = x, y, z \). Using a Zeeman coupling \( H_{\text{Zeeman}} = \sum \left( L_i + gS_i \right) \cdot B \), we compute the contribution from the spin polarization at a site \( i \), in the \( j \) direction, \( \langle S_i \rangle_j \), assuming the orbital contribution, which has been suggested to be small [35], can be separated. We also compute the contribution from the orbital magnetization \( \langle L_i \rangle \), and there is a slight drop in the superconducting state as shown in Fig. S2 in Ref. [15]. The results are shown in Fig. 3, which shows the spin magnetization along the \( x \) and \( y \) directions as the strain changes. Here we plot the ratio between the strained values, and the normal state unstrained cases.

A conventional spin triplet will feature a Knight shift which appears the same as a singlet for the field parallel to the \( d \) vector and shows no change from the normal state for the field perpendicular to the \( d \) vector. On the other hand, OSST pairing leads to intraband pseudospin-singlet pairing occurring near the FS, and interband spin triplet away from the FS. Thus, the low field response behavior is due primarily to the intraband pairing [36], which causes a large drop in the approximately isotropic Knight shift as shown in Fig. 3(a). However, by increasing the field such that it is a significant fraction of the gap size \( (B \sim 0.2 \times |D_0|) \), the interband pairing with \( d \) vector rotation is observable, and such rotation results in an anisotropic Knight shift under strain as shown in Fig. 3(b). Thus, for OSST, the Knight shift is more affected by intraband pseudospin-singlet pairing at low fields, and interband pairing at higher fields. As expected from the \( d \) vector rotation under the \( a \)-axis strain, we find a greater drop in the magnetization from the normal to superconducting state in the \( y \) direction compared with the \( x \) direction, with a difference of about 20% for the larger field value. The magnetization in the \( x \) direction also drops under strain due to the strain bringing the sample deeper into the superconducting state. The value of the drop from the normal to superconducting state depends on the value of the SOC, and by decreasing the SOC, the Knight shift drop and the anisotropy under strain enhance further.

**Extending to three-dimensional bands.** \( \text{Sr}_3\text{RuO}_4 \) has a layered structure, and one expects to see more \( k_z \) dispersion of the bands originating from \( xz \) and \( yz \) orbitals due to their shape, and less dispersion from the \( xy \) orbital. The momentum dependent \( t_{2g} \)-orbital projection of the wave function for the \( \alpha, \beta \) and \( \gamma \) bands on the three-dimensional FS was reported [31], which is consistent with the three dimensional (3D) tight binding model constructed in Ref. [37]. The \( \beta \) and \( \gamma \) bands still have significant overlap of \( xy \) and one dimensional (1D) orbitals, even though detailed composition depends on \( k_z \) as shown in Ref. [31], while the \( \alpha \) band is made of 1D orbitals. Thus the above analysis done in the two-dimensional (2D) system can be generalized to a layered three-dimensional system. The qualitative uniaxial strain effect, i.e., the relative directional dependence of the spin polarization under a uniaxial strain, is independent of the details of \( c \)-axis hopping.
parameters, even though the quantitative drop may depend on the strength of the hopping parameters. Using the tight binding parameters in Ref. [37], we found the $d$ vector directions are similar to the 2D case. The angle $\phi$ represents the tilting from the $ab$ plane, which is about 17–19° depending on $k_z$.

A clear rotation of the averaged $d$ vector is shown as a red arrow in a top corner in Fig. 4, and the main conclusion of the $d$ vector rotation can be generalized to the 3D model including the layer coupling.

**Discussion and summary.** In multiorbital systems, orbital degrees of freedom extend the types of superconducting pairings. Even-parity spin-triplet pairings are allowed when the pairing occurs between different orbitals with the antisymmetric fermionic wavefunction condition, i.e., orbital singlets. In the band basis, this maps to interband, and intraband pairings when SOC is finite [13]. SOC in the OSST pairing determines the intraband gap on the FS. The idea of OSST pairing is not limited to the atomic ($s$-wave) SOC leading to an $s$-wave SOC terms as stated in Ref. [12], and further experimental analysis is required to determine the $d$ vector rotation. For unstrained samples, the impact of pulse energy is stronger than in the case of strained samples as stated in Ref. [12], and further experimental analysis is required to determine if the reduction is more than 50% in the unstrained case.

Another consequence of SOC is a complex order parameter. Generally, the order parameter with the SOC induced spin-singlet components can be written with a phase factor, $\mathbf{D} + e^{i\theta} \hat{\mathbf{D}}$ (where $\mathbf{D}$ is defined to be imaginary such that it is even under TR), where the relative phase between the two is determined by the atomic SOC [13], and $\theta = 0$ for uniform SOC. Despite not breaking time-reversal symmetry, as the time-reversal operator maps the order parameter to itself, the order parameter near impurities may change its relative phase from $\theta = 0$ leading to nontrivial effects. Thus, the multicomponent order parameter may be important to understand the $\mu$SR [6] and Josephson junction [44] results. This is an open topic for future study.

In summary, we showed OSST pairing with SOC leads to a significant reduction of the Knight shift and an anisotropic Knight shift response under uniaxial strain, which can ultimately be used to differentiate spin-triplet from spin-singlet pairing in Sr$_2$RuO$_4$. When the strain is applied along the $a$ axis, interorbital pairing involving $d_{xy}$ and $d_{yz}$ is further enhanced leading to a rotation of the $d$ vector towards the $b$ axis. As a consequence of the $d$ vector rotation, the Knight shift becomes anisotropic relative to the strain axis. It has more drop in the magnetization when the magnetic field is perpendicular to the strain and less when the field is parallel to the strain. Such anisotropy is not expected in the spin singlet, thus we propose the Knight shift measurement with the field along the $a$ axis, which can be compared with the data presented in Ref. [12]. This will ultimately determine a long-standing debate of a possible spin-triplet pairing in Sr$_2$RuO$_4$. This idea can also be extended to other multiorbital systems with significant Hund’s coupling and SOC.

**Acknowledgments.** This work was supported by the Natural Sciences and Engineering Research Council of Canada Discovery Grant 2016-06089, and the Center for Quantum Materials at the University of Toronto.

[1] Y. Maeno, H. Hashimoto, K. Yoshida, S. Nishizaki, T. Fujita, J. G. Bednorz, and F. Lichtenberg, *Nature* (London) **372**, 532 (1994).

[2] A. P. Mackenzie and Y. Maeno, *Rev. Mod. Phys.* **75**, 657 (2003).

[3] A. J. Leggett, *Rev. Mod. Phys.* **47**, 331 (1975).
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[4] T. M. Rice and M. Sigrist, J. Phys.: Condens. Matter 7, L643 (1995).
[5] K. Ishida, H. Mukuda, Y. Kitaoka, K. Asayama, Z. Q. Mao, Y. Mori, and Y. Maeno, Nature (London) 396, 658 (1998).
[6] G. M. Luke, Y. Fudamoto, K. M. Kojima, M. I. Larkin, J. Merrin, B. Nachumi, Y. J. Uemura, Y. Maeno, Z. Q. Mao, Y. Mori, H. Nakamura, and M. Sigrist, Nature (London) 394, 558 (1998).
[7] P. G. Björnsson, Y. Maeno, M. E. Huber, and K. A. Moler, Phys. Rev. B 72, 012504 (2005).
[8] K. D. Nelson, Z. Mao, Y. Maeno, and Y. Liu, Science 306, 1151 (2004).
[9] J. Jang, D. G. Ferguson, V. Vakaryuk, R. Budakian, S. B. Chung, P. M. Goldbart, and Y. Maeno, Science 331, 186 (2011).
[10] C. Kallin, Rep. Prog. Phys. 75, 042501 (2012).
[11] A. P. Mackenzie, T. Scaffidi, C. W. Hicks, and Y. Maeno, npj Quantum Mater. 2, 40 (2017).
[12] A. Pustogow, Y. Luo, A. Chronister, Y.-S. Su, D. A. Sokolov, F. Jerzembeck, A. P. Mackenzie, C. W. Hicks, N. Kikugawa, S. Bagh et al., Nature (London) 574, 72 (2019).
[13] C. M. Puetter and H.-Y. Kee, Europhys. Lett. 98, 27010 (2012).
[14] A. Tamai, M. Zingl, E. Rozbicki, E. Cappelli, S. Riccò, A. de la Torre, S. McKeown Walker, F. Y. Bruno, P. D. C. King, W. Meevasana, M. Shi, M. Radovic, N. C. Plumb, A. S. Gibbs, A. P. Mackenzie, C. Berthod, H. U. R. Strand, M. Kim, A. Georges, and F. Baumberger, Phys. Rev. X 9, 021048 (2019).
[15] See Supplemental Material at http://link.aps.org/supplemental/10.1103/PhysRevResearch.2.032055 for details of the full Hamiltonian, the quasiparticle dispersion, and the magnetization calculations, as well as additional calculations for the orbital magnetization and odd-parity pairing.
[16] A. P. Mackenzie, S. R. Julian, A. J. Diver, G. J. McMullan, M. P. Ray, G. G. Lonzarich, Y. Maeno, S. Nishizaki, and T. Fujita, Phys. Rev. Lett. 76, 3786 (1996).
[17] C. Bergemann, S. R. Julian, A. P. Mackenzie, S. Nishizaki, and Y. Maeno, Phys. Rev. Lett. 84, 2662 (2000).
[18] A. Damascelli, D. H. Lu, K. M. Shen, N. P. Armitage, F. Ronning, D. L. Feng, C. Kim, Z.-X. Shen, T. Kimura, Y. Tokura, Z. Q. Mao, and Y. Maeno, Phys. Rev. Lett. 85, 5194 (2000).
[19] A. Georges, L. d. Medici, and J. Mravlje, Ann. Rev. Condens. Matter Phys. 4, 137 (2013).
[20] S. Hoshino and P. Werner, Phys. Rev. Lett. 115, 247001 (2015).
[21] S. Hoshino and P. Werner, Phys. Rev. B 93, 155161 (2016).
[22] O. Gingras, R. Nourafkan, A.-M. S. Tremblay, and M. Côté, Phys. Rev. Lett. 123, 217005 (2019).
[23] F. B. Kugler, M. Zingl, H. U. R. Strand, S.-S. B. Lee, J. von Delft, and A. Georges, Phys. Rev. Lett. 124, 016401 (2020).
[24] H. G. Suh, H. Menke, P. M. R. Brydon, C. Timm, A. Ramires, and D. F. Agterberg, Phys. Rev. Research 2, 032023 (2020).
[25] A. Ramires and M. Sigrist, Phys. Rev. B 100, 104501 (2019).
[26] W. Huang, Y. Zhou, and H. Yao, Phys. Rev. B 100, 134506 (2019).
[27] A. K. C. Cheung and D. F. Agterberg, Phys. Rev. B 99, 024516 (2019).
[28] E. Pavarini and I. I. Mazin, Phys. Rev. B 74, 035115 (2006).
[29] M. W. Haverkort, I. S. Efthimov, L. H. Tjeng, G. A. Sawatzky, and A. Damascelli, Phys. Rev. Lett. 101, 026406 (2008).
[30] E. J. Rozbicki, J. F. Annett, J.-R. Souquet, and A. P. Mackenzie, J. Phys.: Condens. Matter 23, 094201 (2011).
[31] C. N. Veenstra, Z.-H. Zhu, M. Raichle, B. M. Ludbrook, A. Nicolaou, B. Slomski, G. Landolt, S. Kittaka, Y. Maeno, J. H. Dil, I. S. Efthimov, M. W. Haverkort, and A. Damascelli, Phys. Rev. Lett. 112, 27002 (2014).
[32] M. Kim, J. Mravlje, M. Ferrero, O. Parcollet, and A. Georges, Phys. Rev. Lett. 120, 126401 (2018).
[33] C. W. Hicks, D. O. Brodsky, E. A. Yelland, A. S. Gibbs, J. a. N. Bruin, M. E. Barber, S. D. Edkins, K. Nishimura, S. Yonezawa, Y. Maeno, and A. P. Mackenzie, Science 344, 283 (2014).
[34] A. Steppke, L. Zhao, M. E. Barber, T. Scaffidi, F. Jerzembeck, H. Rosner, A. S. Gibbs, Y. Maeno, S. H. Simon, A. P. Mackenzie, and C. W. Hicks, Science 355, eaaf9398 (2017).
[35] K. Ishida, M. Manago, K. Kinjo, and Y. Maeno, J. Phys. Soc. Jpn. 89, 034712 (2020).
[36] Y. Yu, A. K. C. Cheung, S. Raghu, and D. F. Agterberg, Phys. Rev. B 98, 184507 (2018).
[37] H. S. Ritsing, T. Scaffidi, F. Flicker, G. F. Lange, and S. H. Simon, Phys. Rev. Research 1, 033108 (2019).
[38] S. Benhabib, C. Lupien, I. Paul, L. Berges, M. Dion, M. Nardone, A. Zitouni, Z. Q. Mao, Y. Maeno, A. Georges, L. Taillefer, and C. Proust, arXiv:2002.05916.
[39] S. Ghosh, A. Shekhter, F. Jerzembeck, N. Kikugawa, D. A. Sokolov, M. Brando, A. P. Mackenzie, C. W. Hicks, and B. J. Ramshaw, arXiv:2002.06130.
[40] S. A. Kivelson, A. C. Yuan, B. Ramshaw, and R. Thomale,npj Quantum Mater. 5, 43 (2020).
[41] K. K. Ng and M. Sigrist, Europhys. Lett. 49, 473 (2000).
[42] P. Steffens, Y. Sidis, J. Kulda, Z. Q. Mao, Y. Maeno, I. I. Mazin, and M. Braden, Phys. Rev. Lett. 122, 047004 (2019).
[43] H. U. R. Strand, M. Zingl, N. Wentzell, O. Parcollet, and A. Georges, Phys. Rev. B 100, 125120 (2019).
[44] S. Kashiyawa, K. Saitoh, H. Kashiyawa, M. Koyanagi, M. Sato, K. Yada, Y. Tanaka, and Y. Maeno, Phys. Rev. B 100, 094530 (2019).