Polar Charge Fluctuation and Superconductivity in Organic Conductor

Akihiko Sekine1*, Joji Nasu1, and Sumio Ishihara1,2
1Department of Physics, Tohoku University, Sendai 980-8587, Japan and
2Core Research for Evolutional Science and Technology (CREST), Sendai 980-8587, Japan
(Dated: May 5, 2014)

Superconductivity and polar charge fluctuation are studied in an organic conductor with the dimer-molecule degree of freedom. The extended Hubbard models, where the intra-dimer electronic structure and the inter-dimer Coulomb interactions are taken into account, are analyzed by the random-phase approximation and the fluctuation-exchange approximation. Superconductivity appears in a vicinity of the charge-density wave (CDW) phase where the electronic charge distributions are polarized inside of dimers. The extended s-wave type pairing is favored and its competitive relations with the superconductivity due to the spin fluctuation depends on the triangular lattice geometry. Comparison between two superconductivities realized near the polar and non-polar CDW phases are also presented.

PACS numbers: 74.70.Kn, 77.80.-e, 74.20.Mn

Unconventional superconductivity is one of the central issues in modern solid state physics. Beyond the standard Bardeen-Cooper-Schrieffer theory, non-phononic mechanisms and anisotropic pairing symmetries have been proposed theoretically and examined experimentally in a wide variety of materials, such as copper oxides, iron pnictides, heavy fermion compounds and so on. Organic conductors are one of the families in which exotic superconductivities have been examined intensively and extensively since the discovery of superconductivity in the TMTSF salts.

Recently much attention has been attracted in layered quasi two-dimensional κ-type BEDT-TTF salts, κ-(BEDT-TTF)$_2$X (X: a monovalent anion). The maximum superconducting (SC) transition temperature ($T_c$) of 13.2K is observed in κ-(BEDT-TTF)$_2$Cu[N(CN)$_2$]Cl under high pressure. In the conducting BEDT-TTF layers, two BEDT-TTF molecules are regarded as a molecular dimer which builds a two-dimensional triangular lattice. A geometrical anisotropy in the triangular lattice is adjustable by changing an anion X and/or applying pressure. One hole-carrier per dimer occupies the lowest-unoccupied molecular-orbital in a dimer. When the intra-dimer Coulomb interaction is much stronger than the band width, the system is regarded as a Mott insulator, which is termed a so-called dimer-Mott insulator. An antiferromagnetic ordered state as well as a spin-liquid state are served in the insulating state.

In the conducting BEDT-TTF layers, two BEDT-TTF molecules are regarded as a molecular dimer which builds a two-dimensional triangular lattice. A geometrical anisotropy in the triangular lattice is adjustable by changing an anion X and/or applying pressure. One hole-carrier per dimer occupies the lowest-unoccupied molecular-orbital in a dimer. When the intra-dimer Coulomb interaction is much stronger than the band width, the system is regarded as a Mott insulator, which is termed a so-called dimer-Mott insulator. An antiferromagnetic ordered state as well as a spin-liquid state are served in the insulating state. The polar charge fluctuation is realized cooperatively with the dimer-Mott insulating systems, and a mechanism of the superconductivity.

In this paper, we study superconductivity in low-dimensional organic conductors, in particular, focus on a role of the polar charge fluctuation. We introduce the extended Hubbard model where, in contrast to the previous works, we take into account both the electronic degree of freedom inside of a dimer and the inter-dimer Coulomb interactions, which are essential for the polar charge order and fluctuation. Two types of the molecule configurations where two and four molecules are included in a unit cell are analyzed by the random-phase approximation (RPA) and the fluctuation-exchange (FLEX) methods. The SC phase appears near the charge-density wave (CDW) phase where the electronic charge distribution is polarized inside of dimers, as well as near the spin-density wave (SDW) phase. The extended s-wave symmetry pairing is favored near this CDW phase and is realized cooperatively with the d$_x$-$d_y$-type pairing due to the spin fluctuation. Contribution of the polar charge fluctuation is remarkable around the so-called frustration point where the dimer molecules are arranged on an equilateral triangular lattice. A reentrant feature of the phase boundary is observed for the superconductivity induced by the polar charge fluctuation.

We introduce the two types of the molecule configurations, termed the two-band model and the four-band model, where inequivalent two molecules and four molecules are introduced in a unit cell, respectively. Schematic views are shown in Fig. I. First we present the Hamiltonian for the two-band model:

$$\mathcal{H} = -\sum_{i\sigma} t_{i\sigma} \big(e_{i\sigma}^\dagger c_{i\sigma} + \text{H.c.}\big) + U \sum_{i\gamma} n_{i\gamma\uparrow} n_{i\gamma\downarrow}$$
$$+ U' \sum_{i} n_{i\sigma} n_{i\sigma} - \sum_{\langle ij \rangle \gamma\gamma'} t_{ij}^{\gamma\gamma'} \left( e_{i\gamma\sigma}^\dagger c_{j\gamma'\sigma} + \text{H.c.} \right)$$
$$+ \sum_{\langle ij \rangle \gamma\gamma'} V_{ij}^{\gamma\gamma'} n_{i\gamma} n_{j\gamma'},$$

where $c_{i\sigma}$ is an electron annihilation operator at the i-th dimer with molecular orbital $\gamma(= a,b)$, and spin $\sigma(=\uparrow,\downarrow)$, and $n_{i\gamma}(=\sum_{\sigma}\sum_{\gamma} c_{i\gamma\sigma}^\dagger c_{i\gamma\sigma})$ is the number operator. The first term represents the intra-dimer electron transfer, and the second and third terms represent the intra-molecule electron transfer. For the four-band model, the second and third terms are replaced by

$$\sum_{\langle ij \rangle \gamma\gamma'} V_{ij}^{\gamma\gamma'} n_{i\gamma} n_{j\gamma'}.$$
The inter-dimer Coulomb interactions, and the bubble-type diagrams where the Coulomb interactions are considered, are in Fig. 1(a). Spatial difference between the two molecules in the same dimer is neglected, for simplicity. The average electron density per dimer is fixed to be 3, corresponding to the geometrical anisotropy in a triangular lattice is represented by the two kinds of $t_{ij}$, $t'$, and those of $V_{ij}$, $V'$, as shown in Fig. 1(a). Spatial difference between the two molecules in the two band model, and $t_b$, $t_p$, and $t_q$ are given.

In the four band model, the inter-dimer electron transfers and electron-electron interactions, respectively. In the two-band model, the inter-dimer transfers and Coulomb interactions are assumed to be diagonal as $t_{ij}^c = \delta_{ij} - t_{ij}$ and $V_{ij}^{c\gamma} = \delta_{ij} - V_{ij}$, for simplicity. The geometrical anisotropy in a triangular lattice is represented by the terms $t_{ij}$, $t'$, and those of $V_{ij}$, $V'$, as shown in Fig. 1(a).

The self energy is given by

$$
\Sigma_k(\epsilon_n) = \frac{T}{N} \sum_{k' \epsilon_m} V_{k-k'}(\epsilon_m) (\epsilon_n - \epsilon_m)
$$

where $V_{ij}$ is the Fourier transform of $V_{ij}$, and $V_{ij}^{c\gamma}$ is the interaction $V_{ij}^{c\gamma}$ in the RPA scheme are given by

$$
\chi_q^c(\omega_n) = \left[ 1 + \frac{\chi_0^c(\omega_n)}{\omega_n} \right]^{-1} \chi_q^c(\omega_n),
$$

where $-\mu$ and $+\mu$ signs are for the spin and charge susceptibilities, respectively. "Hat" implies the matrix, and $\chi_q^c(\omega_n)$ is the bare susceptibility with the Matsubara frequency $\omega_n$ defined by

$$
\chi_q^c(\omega_n) = \frac{T}{N} \sum_{k \epsilon_m} G^0_q(\chi_q)(\epsilon_m) 
$$

We define $U_q = W_q^{\uparrow\downarrow} - W_q^{\downarrow\uparrow}$, and the bare Green’s function $G^0_q(\epsilon_n)$. In the FLEX method, the spin and charge susceptibilities are calculated in the equations given in Eqs. (3) and (4) where the bare Green’s functions are replaced by the full Green’s functions defined by the Dyson equation, $G_k(\epsilon_n) = G^0_k(\epsilon_n) - \Sigma_k(\epsilon_n)$. The self energy is given by

$$
\Sigma_k(\epsilon_n) = \frac{T}{N} \sum_{k' \epsilon_m} V_{k-k'}(\epsilon_m) (\epsilon_n - \epsilon_m)
$$

where we introduce the following static susceptibilities: the static spin susceptibility

$$
\chi_q^{\text{spin}} = \int_0^\beta d\tau (S_q^z(\tau) S_{-q}^z),
$$

the static charge susceptibility

$$
\chi_q^{\text{charge}} = \int_0^\beta d\tau (Q_q(\tau) Q_{-q}),
$$

and the static polarization susceptibility

$$
\chi_q^{\text{polar}} = \int_0^\beta d\tau (P_q(\tau) P_{-q}),
$$

where $Q(\tau) = e^{iH\tau} Q e^{-iH\tau}$ and $S_q^z$, $Q_q$, and $P_q$ are the Fourier transforms of the spin, charge and polarization operators, respectively. These are defined by $S_q^z = (1/2) \sum_{\sigma} \langle S_q^{z\sigma} \rangle$, $Q_q = (1/2) \sum_{\sigma} \langle S_q^{\sigma} \rangle$, and $P_q = (1/2) \sum_{\sigma} \langle P_q^{\sigma} \rangle$ for the two-band in Fig. 1(b).
model, and \( S^\gamma_{1,2} = (1/2) \sum_{\xi=(1,2,3,4)} n_i \xi \{ n_i \xi + n_i \bar{\xi} \} \),
\( Q_i = (1/2) \sum_{\alpha=1,2} n_{i \alpha} \{ n_{i \alpha} + n_{i \bar{\alpha}} \} \) and \( P_i = (1/2) \sum_{\xi=(1,2,3,4)} n_{i \xi} \{ n_{i \xi} + n_{i \bar{\xi}} \} \) for the four-band model. The static susceptibilities are represented by the spin- and charge-susceptibility matrices previously introduced. The explicit representations, for example for the two band model, are given as

\[
\chi^{\text{spin}}_q = \frac{1}{2} \sum_{\gamma_1,\gamma_2} \chi^S_{(\gamma_1\gamma_2)}(0),
\]

\[
\chi^{\text{charge}}_q = \frac{1}{2} \sum_{\gamma_1,\gamma_2} \chi^C_{(\gamma_1\gamma_2)}(0),
\]

and

\[
\chi^{\text{polar}}_q = \frac{1}{2} \sum_{\gamma_1,\gamma_2} \chi^C_{(\gamma_1\gamma_2)}(0) \epsilon_{\gamma_1\gamma_2},
\]

where \( \epsilon_{\gamma_1\gamma_2} = 1 \) for \( \gamma_1 = \gamma_2 \) and \( -1 \) for \( \gamma_1 \neq \gamma_2 \). The formulas for the four-band model are given in the similar ways. On the analogy of the definitions of the static-susceptibilities, we introduce the "bare" static charge-susceptibility \( \chi^0_\text{charge} \), and the "bare" static polarization susceptibility \( \chi^0_\text{polar} \). These are defined in Eqs. (11) and (12), where the expectations \( \langle \cdot \rangle \) are replaced by those without the Coulomb interactions, and are calculated by Eqs. (11) and (12), where \( \chi^C_q(0) \) is replaced by \( \chi^0_q(0) \) in the two-band model.

The SC transition temperature and the gap function are calculated by the Eliashberg equation. The linearized equation is given by

\[
\lambda \Delta^{(l)}(\mathbf{k}, \alpha, \beta, s) = -\frac{T}{N} \sum_{\mathbf{k'}, \mu, \mu'} \Gamma^{(l)}_{\mathbf{k}-\mathbf{k'}}(\alpha, \beta, \gamma, \delta, s, \mu, \mu') \epsilon_{\alpha s} \epsilon_{\beta s} \sum_{\mathbf{k}, \nu, \nu'} G_{\mathbf{k}', \mu, \mu'}(\epsilon_{\nu s}) G_{\mathbf{k}, \nu, \nu'}(\epsilon_{\nu s}) \lambda \Delta^{(l)}_{\mathbf{k}, \mu, \mu', \nu, \nu'}, \quad \lambda \Delta^{(l)}(\mathbf{k}, \alpha, \beta, s) = -\frac{T}{N} \sum_{\mathbf{k'}, \mu, \mu'} \Gamma^{(l)}_{\mathbf{k}-\mathbf{k'}}(\alpha, \beta, \gamma, \delta, s, \mu, \mu') \epsilon_{\alpha s} \epsilon_{\beta s} \sum_{\mathbf{k}, \nu, \nu'} G_{\mathbf{k}', \mu, \mu'}(\epsilon_{\nu s}) G_{\mathbf{k}, \nu, \nu'}(\epsilon_{\nu s}) \lambda \Delta^{(l)}_{\mathbf{k}, \mu, \mu', \nu, \nu'},
\]

where \( l = s(t) \) for a spin singlet (triplet) pairing. A parameter \( \lambda \) is introduced and \( T_c \) is determined by the condition \( \lambda = 1 \). We define the interactions

\[
\Gamma^{(s)}_q = \frac{3}{4} U^s q^s q^s - \frac{1}{2} \varepsilon^s q^s q^s + \frac{1}{2} \left( \bar{U}^s + U^s \right),
\]

and

\[
\Gamma^{(t)}_q = -\frac{1}{4} U^t q^t q^t - \frac{1}{2} \varepsilon^t q^t q^t + \frac{1}{2} \left( \bar{U}^t + U^t \right). \quad (15)
\]

We take \( 32 \times 32 \times 64 \) momentum-point meshes in the Brillouin zone and up to \( 8,192 \times 16,384 \) Matsubara frequencies in the numerical calculations by RPA (FLEX).

First we show the numerical results in the two-band model. We take \( t'/t = V'/V > 1 \) that correspond to the geometrical anisotropy in a triangular lattice for \( \kappa-(\text{BEDT-TTF})_2\text{Cu}_2(\text{CN})_3 \). The phase diagram at \( T/t_0 = 0.01 \) calculated by RPA is presented in Fig. 2a. The topologically similar phase diagrams are reported in Refs. [29] and [30] where one molecule without the dimer degree of freedom is introduced at each site in triangle and square lattices.

The boundary for the SDW (CDW) phase is identified by the condition \( \det[1 - (\alpha)\bar{U}^c_\alpha(0) \bar{U}^c_\alpha(0)] = 0 \). The SDW and CDW phases appear in small \( t/t_0 \) and small \( V/t_0 \) regions, respectively. The phase boundary for SDW does not depend on \( V \), because the inter-dimer Coulomb interactions are not included in \( \bar{U}^c_\alpha \). The momenta where the static spin susceptibility \( \chi^{(s)}_\alpha \) takes the maxima are \( q = (0.75 \pi, 0.375 \pi) \) and \( (-0.75 \pi, -0.375 \pi) \) at \( (t/t_0, V/t_0) = (0.65, 0.5) \). To examine the CDW phase in more detail, we calculate the static charge and polar susceptibilities. At the boundary for the CDW phase, \( \chi^{(s)}_\alpha \) diverges and \( \chi^{(t)}_\alpha \) does not. That is, the charge disproportionation occurs inside of dimers and the electron number per dimer is equal for all dimers. A schematic charge configuration is shown in Fig. 2. The momenta where \( \chi^{(p)}_\alpha \) takes the maxima are \( q = (2 \pi/3, 2 \pi/3) \) and \( (-2 \pi/3, -2 \pi/3) \) at \( (t/t_0, V/t_0) = (0.75, 0.8) \), implying a three-fold periodicity of the electronic dipoles moment and no macroscopic polarization.

The SC phase boundary and the SC gap functions are shown in Fig. 2. The SC phase volume near the SDW phase is larger than that near the CDW phase. In the present lattice structure with the \( D_{2h} \) point group, a spin-singlet SC-gap symmetry is classified by the \( \Gamma_{1g} \) and \( \Gamma_{1g} \) irreducible representations. The \( d_{z\bar{z}} \)-type \( A_{1g} \) SC gap is realized near the SDW phase at \( (t/t_0, V/t_0) = (0.7, 0.5) \) and \( (0.75, 0.8) \), and is gradually changed into the extended \( s \)-type \( A_{1g} \) gap near the CDW phase. This extended \( s \)-type gap is also obtained by the calculation where we set \( \bar{U}^s q^s \bar{U}^s = 0 \) in Eqs. (14) and (15).
anisotropy in the electron transfer, i.e. \( r \) ties at the transition temperatures at \( T/t_0 = 1.5 \) to those at \( V = 0 \), i.e. \( r_T \equiv T_c(V/t_0 = 1.5)/T_c(V = 0) \) as functions of the anisotropy in the electron transfer, \( t'/t \). The transition temperatures at \( V/t_0 = 1.5 \) are also shown. The charge and spin susceptibilities at \( T/t_0 = 0.0025 \) for \( t'/t = 0.8 \) and those for \( t'/t = 1.05 \) are presented in (c) and (d), respectively. All data are obtained in the two-band model analyzed by FLEX. Other parameter values are chosen to be \( U/t_0 = 4, U'/t_0 = 2.8, \) and \( t'/t_0 = 0.5 \).

It is concluded that the polar charge fluctuation induces the extended \( s \)-type SC state.

The two-band model is further analyzed by the FLEX method. The CDW phase with the electric dipole moments inside of dimers, obtained by RPA [see Fig. 2a], is reproduced, but the SDW phase is not realized in finite temperatures due to the Mermin-Wagner’s theorem satisfied in FLEX. The SC transition temperatures near the CDW phase are shown in Fig. 2c. The SC transition temperatures are determined by fitting the numerical data of \( \lambda \) calculated down to 0.002\( t_0 \) by a function \( \lambda = -a \log T + b \). The transition temperature monotonically increases, when the system approaches to the phase boundary. This is evidence that the SC state around this region is caused by the polar charge fluctuation. A numerical value of \( T_c \) for \( t'/t = 1.05 \) (\( t'/t = 0.8 \)) is estimated to be about 0.5K (5K) at \( V/t_0 = 1.5 \), when \( t_0 = 0.2eV \) is assumed. It is worth to note that, in comparison with the results by RPA [see Fig. 2a], the polar CDW phase and the SC phase near the CDW phase are realized until at least \( t/t_0 = 0.3 \). This originates from a suppression of the SDW instability due to the Mermin-Wagner’s theorem and is consistent with \( \kappa \)-type BEDT-TTF salts.

Effects of the triangular-lattice geometry on the SC transition temperature is examined by calculating a ratio of the transition temperature \( \frac{T_c}{V/t_0} = 1.5 \) to that at \( V = 0 \), i.e. \( r_T \equiv T_c(V/t_0 = 1.5)/T_c(V = 0) \) as functions of the anisotropy in the electron transfer, \( t'/t \). We also plot the \( t'/t \) dependence of \( T_c \) at \( V/t_0 = 1.5 \), which is consistent with the results in Ref. 13. It is clearly shown that \( r_T \) takes its maximum around the so-called frustration point \( t'/t = 1 \). That is, a relative contribution of the polar charge fluctuation to the spin fluctuation increases around \( t'/t = 1 \).  

In Figs. 3c and 3d, the spin susceptibility and the polarization susceptibility at \( T/t_0 = 0.0025 \) are shown for \( t'/t = 0.8 \) and 1.05, respectively. We define \( \chi^{\text{spin}}_q \) and \( \chi^{\text{polar}}_q \) as the maxima of \( \chi^{\text{spin}}_q \) and \( \chi^{\text{polar}}_q \), respectively, when the momenta \( q \) are varied. It is shown that \( \chi^{\text{spin}}_q \) at \( t'/t = 1.05 \) is reduced from that in \( t'/t = 0.8 \), on the other hand, \( \chi^{\text{polar}}_q \) are almost independent of \( t'/t \). This is because the spin fluctuation is almost governed by \( \chi^{\text{spin}}_q \), which is sensitive to the anisotropy of the electron trans-
The SC phase induced by the polar charge fluctuation does not depend on a detailed lattice structure.

Finite temperature phase diagram near the polar and nonpolar CDW phases calculated by RPA is presented in Fig. 5 as a function of $V_q$. A reentrant feature is observed in the phase boundary between the polar CDW phase and the metallic or SC phase, but not in the boundary for the non-polar CDW phase. This is interpreted by the competitive momentum dependences of the charge susceptibility and the effective interaction as follows. In the RPA scheme, the momentum dependences of the susceptibilities are governed by those of the bare susceptibilities and the long-range Coulomb interactions. Contour maps of $V^{\rm max}(q)$, which is defined as the largest eigenvalues of the Fourier transforms of the inter-dimer interaction $V_{ij}$, are calculated. Results near the polar CDW phase ($V_q/t_0 = 0.35$) and the non-polar CDW phase ($V_q/t_0 = 0.85$) are plotted in Figs. 5(b) and (c), respectively. We also show the contour maps of the “bare” static polar- and charge-susceptibilities $\chi^0_{\text{polar}}$ and $\chi^0_{\text{charge}}$, in Fig. 5(d) and (e), respectively. It is noticed that momentum dependences of $\chi_q^{0\text{polar}}$ and $V^{\text{max}}(q)$ at $V_q = 0.35$ [see Figs. 5(b) and (d)] are competitive with each other; $\chi^0_q$ (polar) $V^{\text{max}}(q)$ takes its maxima (minima) along $q_y = 0$. This fact brings about a competition between the susceptibility and the inter-dimer interaction, and suppresses the instability toward the polar-CDW at low temperature. On the other hand, such competitive momentum dependences are weak for $\chi^0_q$ (charge) and $V^{\text{max}}(q)$ near the nonpolar CDW phase [see Figs. 5(c) and (e)]. A similar reentrant feature and the competitions are suggested in Refs. 29 and 33.

![Figure 6](image-url)
Finally, we mention the SC properties in the two-band model with \( t'/t < 1 \), that corresponds to \( \kappa-(\text{BEDT-TTF})_2X \) for X=Cu(NCS)_2 and Cu[N(CN)_2]Cl. Phase diagram and gap functions calculated by RPA are presented in Fig. 6. In contrast to the results in Fig. 2 where we take \( t'/t > 1 \), the SC phase with the \( d_{x^2-y^2}\)-type \( B_{1g} \) symmetry gap appears near the SDW phase. This change in the gap symmetry by changing \( t'/t \) is consistent with the previous results. The extended \( s\)-type \( A_{1g} \) gap is also realized near the CDW phase. As a consequence of the competition between the two SC phases with different symmetry gaps, a phase volume of the extended \( s\)-type SC is shrunk. We propose that the SC state induced by the polar charge fluctuation is favored in the case of \( t'/t > 1 \).

In summary, we study the superconductivity in the low-dimensional organic conductor where electronic structure inside of dimer is taken into account. We focus on roles of the polar charge fluctuation on the superconductivity. The two types of the extended Hubbard models are analyzed by the RPA and FLEX methods. We find that the extended \( s\)-type SC state is realized near the CDW phase where the charge distributions inside of dimers are polarized. A monotonic increase of \( T_c \) when the system approaches to the polar CDW phase is direct evidence of the SC state induced by the polar charge fluctuation. The obtained \( T_c \) values are comparable to those in the \( \kappa-(\text{BEDT-TTF})_2X \) systems. Competitive relationship between this SC state and that induced by the spin fluctuation depends on the geometrical anisotropy in a triangle lattice; two SC states are cooperative (competitive) with each other in \( t'/t > 1 \) (\( t'/t < 1 \)). The reentrant structures emerge in the boundary between the SC state and the polar-CDW state. The present study opens a new direction for research of the superconductivity and ferroelectricity.

Acknowledgments

Authors would like to thank M. Naka, S. Yamazaki, H. Takashima, and T. Watanabe for their valuable discussions. This work was supported by KAKENHI from MEXT and Tohoku University “Evolution” program. JN is supported by the global COE program “Weaving Science Web beyond Particle-Matter Hierarchy” of MEXT, Japan. Parts of the numerical calculations are performed in the supercomputing systems in ISSP, the University of Tokyo.

*Present address: Institute for Materials Research, Tohoku University, Sendai 980-8577, Japan.

---

1. D. Jérome, A. Mazaud, M. Ribault and K. Bechgaard, J. Phys. Lett. (Paris) 41, 95 (1980).
2. A. Ardavan, S. Brown, S. Kagoshima, K. Kanoda, K. Kuroki, H. Mori, M. Ogata, S. Uji, and J. Wosnitza, J. Phys. Soc. Jpn. 81, 011004 (2012).
3. K. Kanoda, J. Phys. Soc. Jpn. 75, 051007 (2006).
4. K. Miyagawa, K. Kanoda, and A. Kawamoto, Chem. Rev. 104, 5653 (2004).
5. J. M. Williams, A. J. Schultz, U. Geiser, K. D. Carlson, A. M. Kini, H. H. Wang, W.-K. Kwok, M.-H. Whangbo, J. E. Schirber, Science, 252, 1501 (1991).
6. H. Kino and H. Fukuyama, J. Phys. Soc. Jpn. 64, 1877 (1995).
7. H. Kino and H. Fukuyama, J. Phys. Soc. Jpn. 65, 2158 (1996).
8. Y. Shimizu, K. Miyagawa, K. Kanoda, M. Maesato, and G. Saito, Phys. Rev. Lett. 91, 107001 (2003).
9. S. Yamashita, Y. Nakazawa, M. Oguni, Y. Oshima, H. Nojiri, Y. Shimizu, K. Miyagawa, and K. Kanoda, Nat. Phys. 4, 459 (2008).
10. M. Yamashita, N. Nakata, Y. Kasahara, T. Sasaki, N. Yoneyama, N. Kobayashi, S. Fujimoto, T. Shibauchi, and Y. Matsuda, Nat. Phys. 5, 44 (2009).
11. H. Kino and H. Kontani, J. Phys. Soc. Jpn. 67, 3691 (1998).
12. H. Kondo and T. Moriya, J. Phys. Soc. Jpn. 67, 3695 (1998).
13. H. Kondo and T. Moriya, J. Phys. Soc. Jpn. 70, 2800 (2001).
14. K. Kuroki, T. Kimura, R. Arita, Y. Tanaka and Y. Matsuda, Phys. Rev. B 65, 100516(R) (2002).
15. H. Kondo, and T. Moriya, J. Phys. Soc. Jpn. 73, 812 (2004).
16. T. Watanabe, H. Yokoyama, Y. Tanaka, and J. Inoue, J. Phys. Soc. Jpn. 75, 074707 (2006).
17. M. Abdel-Jawad, I. Terasaki, T. Sasaki, N. Yoneyama, N. Kobayashi, Y. Uesu and C. Hotta, Phys. Rev. B 82, 125119 (2010).
18. M. Naka and S. Ishihara, J. Phys. Soc. Jpn., 79, 063707 (2010).
19. C. Hotta, Phys. Rev. B 82, 241104(R) (2010).
20. H. Li, R. T. Clay, and S. Mazumdar, J. Phys. Cond. Mat. 22, 272201 (2010).
21. P. Lunkenheimer, J. Muller, S. Krohns, F. Schrettle, A. Loidl, B. Hartmann, R. Rommel, M. de Souza, C. Hotta, J. A. Schlueter, and M. Lang, Nat. Mat. 11, 755 (2012).
22. S. Iguchi, S. Sasaki, N. Yoneyama, H. Taniguchi, T. Nishizaki, and T. Sasaki, arXiv:1208.0633.
23. T. Takimoto, T. Hotta, and K. Ueda, Phys. Rev. B 69, 104504 (2004).
24. T. Nakano and K. Kuroki, J. Phys. Soc. Jpn. 75, 034706 (2006).
25. G. Esirgen, H.-B. Schuttler, and N. E. Bickers, Phys. Rev. Lett. 82, 1217 (1999).
26. A. Kobayashi, S. Katayama, K. Noguchi, and Y. Suzumura, J. Phys. Soc. Jpn. 73, 2053 (2004).
27. K. Yoshimi, H. Seo and H. Maebashi, Meeting Abst. Phys. Soc. Jpn. 68, 876 (2012).
28. S. Onari, R. Arita, K. Kuroki and H. Aoki, Phys. Rev. B 73, 014526 (2006).
29. A. Kobayashi, Y. Tanaka, M. Ogata, and Y. Suzumura, J. Phys. Soc. Jpn. 73, 1115 (2004).
30. Y. Tanaka, Y. Yanase, and M. Ogata, J. Phys. Soc. Jpn. 73, 2053 (2004).
31. H. Kontani and M. Ohno, Phys. Rev. B 74, 014406 (2006).
32. T. Komatsu, N. Matsukawa, T. Inoue, and G. Saito, J. Phys. Soc. Jpn. 65, 1340 (1996).
33. J. Merino and R. H. McKenzie, Phys. Rev. Lett. 87, 237002 (2001).