Superconducting gap structure of heavy-Fermion compound URu$_2$Si$_2$ determined by angle-resolved thermal conductivity

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Abstract. In heavy-Fermion compound URu$_2$Si$_2$, unusual superconductivity is embedded in an enigmatic ‘hidden order’ phase. Recently, it has been shown that URu$_2$Si$_2$ is essentially a multiband superconductor associated with the semimetallic compensated electronic structure. Here, to pin down the detailed superconducting gap structure, we have performed thermal transport measurements on ultraclean URu$_2$Si$_2$ single crystals in magnetic fields rotating various directions relative to the crystal axes. By changing the amplitude of magnetic fields, we determined the nodal topology of electron and hole band separately. The results indicate a new type of unconventional superconductivity with two distinct gaps, in which horizontal line nodes lie within the basal $ab$ plane of the light-hole band with small gap and point nodes along the $c$-axis in the heavy electron band with large gap. This gap structure is consistent with ‘chiral’ d-wave symmetry with a form $\hat{k}_z(\hat{k}_x + i\hat{k}_y)$.

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Superconductivity in heavy-Fermion systems continues to be a central focus of investigations into strongly correlated electron systems. Among them, URu$_2$Si$_2$ has mystified researchers since the unconventional superconducting state ($T_c = 1.45$ K) is embedded within the ‘hidden order’ state ($T_h = 17.5$ K) [1]–[3]. Tiny magnetic moment appears ($M_0 \sim 0.02 \mu_B$) below $T_h$ [4], but $M_0$ is by far too small to explain the large entropy released during the transition. Although several exotic order parameters have been proposed for the hidden order phase [5]–[10], the genuine hidden order parameter is still an open question.

The superconductivity of URu$_2$Si$_2$ is intimately related to the hidden order phase. Thermodynamic and transport measurements revealed that the electronic excitation gap is formed at a large portion of the Fermi surface below $T_h$ and most of the carriers ($\sim 90\%$) disappear, resulting in a semimetallic ground state [11]–[14]. Superconductivity with such a low carrier density is remarkable since the superfluid density is very low, in some ways reminiscent of underdoped cuprates. The hidden order is affected by external parameters, such as pressure and magnetic fields. Figure 1(a) displays the field–temperature ($H-T$) phase diagram [15]–[19]. The superconductivity occurs deep inside the hidden order state, occupying a tiny portion in the diagram. Figure 1(b) displays the pressure–temperature ($p-T$) phase diagram [20]–[24]. Above the critical pressure of $p_c \sim 0.7$ GPa, a true bulk antiferromagnetic ordered state with large moments ($\sim 0.4 \mu_B$ U$^{-1}$) emerges. Although no dramatic modification of the Fermi surface occurs between the hidden order and antiferromagnetic phases, the superconductivity disappears at $p_c$. This indicates that superconductivity cannot coexist with the antiferromagnetic order, in contrast with other heavy-Fermion superconductors.

Thus, the hidden order provides an intriguing stage on which a new type of unconventional superconductivity appears. However, the experimental investigation for the superconducting state has seen little progress, mainly because the superconducting state is highly sensitive to disorder. Very recent studies using ultraclean single crystals have enabled us to develop an advanced understanding about the highly unusual superconducting state of URu$_2$Si$_2$ [14]. A number of unprecedented superconducting properties have been reported, including superconductor–insulator-like first-order transition at upper critical fields [14], formation of the quasiparticle Bloch state [25] and quantum transport of quasiparticles [26].
Figure 1. Phase diagram of URu$_2$Si$_2$. SC, HO and AF stand for the superconducting, hidden order and antiferromagnetic phase, respectively. (a) $H$–$T$ phase diagram. The hidden order phase is destroyed by the strong magnetic fields accompanying a radical reconstruction of the Fermi surface. Well below $T_c$, field-induced destruction of the hidden order occurs at $\approx$36 T, followed by a cascade of phase transitions between the consecutive field-induced phases. Superconductivity occupies a tiny portion in the diagram. (b) $p$–$T$ phase diagram. Above the critical pressure of $\sim$0.7 GPa, a true antiferromagnetic ordered state with large moments ($\sim$0.4$\mu_B$) emerges. Superconductivity coexists with the hidden order but does not coexist with the antiferromagnetic order.

The most important subject for understanding the unusual superconductivity may be the elucidation of the superconducting gap structure. The specific heat $C(T) \propto T^2$ at low temperatures [27] and nuclear magnetic resonance relaxation rate $T_1^{-1} \propto T^3$ [28] have suggested a line node in the superconducting gap or more accurately that the density of states (DOS) at low energy $N(E) \propto |E|$. However, recent transport experiments revealed the nearly perfect compensated electronic structure having the same number of electrons and holes in the hidden order state [29]. This indicates that URu$_2$Si$_2$ is an essentially multiband superconductor. Therefore, the determination of the gap structure of ‘each band’ is essentially important for understanding the exotic superconductivity of URu$_2$Si$_2$.

Very recently, the thermal transport [14] and heat capacity [30] measurements have pointed out the presence of point nodes along the c-axis. Moreover, thermal transport studies have suggested two distinct gap structures having different nodal topology. Based on these results, ‘chiral’ d-wave symmetry with a form $\hat{k}_x (\hat{k}_x + i \hat{k}_y)$ has been proposed. A more detailed clarification of the nodal structure is strongly required. Recently, it has been demonstrated that the thermal conductivity and the heat capacity measurements under magnetic field rotated relative to the crystal axes are a powerful method for determining the detailed gap structure including the nodal directions in the bulk [31]. In this paper, we have performed the angular variation measurements of the thermal conductivity in the superconducting state of ultraclean URu$_2$Si$_2$ single crystals, which provide a stringent test for the proposed gap structure.
1. Experimental

The single crystals of URu$_2$Si$_2$ were grown by the Czochralski pulling method in a tetra-arc furnace, using an electrotransport-purified uranium metal as a starting material [32]. The extremely low residual resistivity $\rho_0 = 0.48 \, \mu\Omega \text{cm}$ and large residual resistivity ratio of 670 attest the very high crystal quality. The upper critical fields in the magnetic field $H$ parallel to the $ab$ plane and $c$-axes, $H_{c2}^{ab}$ and $H_{c2}^c$, are 12.5 and 2.8 T, respectively.

The thermal conductivity $\kappa$ was measured down to 30 mK ($\simeq T_c / 50$) by the standard four-wire steady-state method along the $a$-axis (heat current $q \parallel a$). The contact resistance is less than 10 m$\Omega$ at low temperatures. In order to apply $H$ with high accuracy relative to the crystal axes, we used a system with two superconducting magnets generating $H$ in two mutually orthogonal directions and a dilution refrigerator equipped on a mechanical rotating stage at the top of a Dewar. By computer controlling the two superconducting magnets and the rotating stage, we were able to rotate $H$ with a misalignment of less than 0.02° from each axis.

2. Superconducting state

2.1. Multiband superconductivity

According to recent magneto-transport studies in the normal state by using ultraclean single crystals, the nearly perfect compensation having the same number of electrons and holes is realized in URu$_2$Si$_2$ [14, 33]. It has been pointed out that the transport properties are mainly governed by the hole band with light mass, which is nearly spherical, whereas thermodynamic properties are governed by the electron band with heavy mass, which is anisotropic (cigar-like). The light hole band, most likely to be $\alpha$ band with the largest volume, has been confirmed by the experiments of de Haas–van Alphen [33, 34] and Shuvnikov–de Haas [17, 35]. The heavy electron band has been confirmed by very recent cyclotron resonance experiments [36]. In what follows, we will discuss the superconducting gap structure by these two bands for simplicity, although there are additional small minor bands.

The inset of figure 2 shows the temperature dependence of the thermal conductivity divided by temperature, $\kappa / T$, in zero field. The Wiedemann–Franz ratio $L = (\kappa / T) \rho$ immediately above $T_c$, which is close to the Lorentz number $L_0 = 2.44 \times 10^{-8} \, \Omega \text{W K}^{-2}$, $L = 1.00 \pm 0.02 L_0$ and small phonon contribution shown by the dashed line in the inset of figure 2 [13] indicates dominance of electronic contribution in the superconducting state. The main panel of figure 2 shows the temperature dependence of $\kappa / T$ at low temperatures in zero field, plotted as a function of $T^2$. A finite residual term at $T \to 0 \, \text{K}$ limit, $\kappa_{00} / T$, is clearly resolved. In s-wave superconductors, $\kappa / T$ tends to zero at $T \to 0 \, \text{K}$ limit, because of the absence of quasiparticle state. Therefore the present results indicate the presence of normal electronic fluid even at $T \to 0 \, \text{K}$ limit, which is expected in nodal superconductors. As discussed in [14] the magnitude of $\kappa_{00} / T$ is close to the value expected in the superconductor with line nodes [37, 38].

Figure 3 shows $\kappa (H) / T$ obtained by extrapolating $T \to 0 \, \text{K}$ at each field value as a function of magnetic field normalized by upper critical fields. There are three characteristic field regimes denoted (I), (II) and (III). In low field (I)-regime, $\kappa (H) / T$ shows a steep increase for both $H \parallel c$ and $H \parallel a$. In the higher field (II)-regime, $\kappa (H) / T$ is nearly field independent for $H \parallel c$, while $\kappa (H) / T$ increases with $H$ with a convex shape after showing a plateau behavior for $H \parallel a$. At $H_{c2}$, for both field directions, $\kappa (H) / T$ exhibits a distinct jump to a very low value.
Figure 2. Main panel: $\kappa/T$ plotted as a function of $T^2$ at very low temperatures. Inset: $T$ dependence of $\kappa/T$ at high temperatures. Dashed line is the phonon contribution obtained from [13].

Figure 3. (a) The field dependence of $\kappa/T$ for $H \parallel c$ (open blue circles) and $H \parallel a$ (filled red circles) at $T \rightarrow 0$ K as a function of $H/H_{c2}$ and $H/H_{c2}$, respectively. There are three characteristic field regimes, (I), (II) and (III).

In the normal state regime (III), $\kappa(H)/T$ decreases with increasing magnetic field, associated with the magnetoresistance in the normal state. This characteristic field dependence of the thermal conductivity provides several pieces of important information of the gap structure.

The initial steep increase and subsequent plateau behavior of the thermal conductivity indicate that a substantial portion of quasiparticle is already restored in regime (I), much below...
the upper critical fields. This indicates the presence of small superconducting gap, which is typical behavior of the multiband superconductivity. Such a field dependence bears a strong resemblance to those of MgB$_2$ [39, 40] and PrOs$_4$Sb$_{12}$ [41]. The multiband superconductivity is compatible with the compensated electronic structure with two main bands having different masses. Then it is natural to consider that the thermal conductivity in the low field regime (I) is governed by light hole band and is governed by heavy electron band in the higher field regime (II).

2.2. Parity

Before discussing the superconducting gap structure, we discuss the parity of the superconducting wave function. The remarkable step-like reduction of $\kappa(H)/T$ at upper critical fields is observed for both $H \parallel a$ and $H \parallel c$. Since the thermal conductivity is related to an entropy flow, its jump (irrespective of its directions) immediately indicates the occurrence of the first-order phase transition, as reported in CeCoIn$_5$ [42]. Since the Maki parameter of URu$_2$Si$_2$ is more than unity for both field directions, the origin of the first-order transition is most likely due to the Pauli paramagnetic limiting. This suggests an even parity (spin singlet) pairing state in URu$_2$Si$_2$.

Usually, the heat conduction just below $H_{c2}$ is smaller than that in the normal state owing to the reduction of the number of heat carriers and the enhancement of the scattering rate by vortices. Therefore the observed field-induced transition at $H_{c2}$ of URu$_2$Si$_2$ with an enhanced entropy flow in the superconducting state compared with the normal state is highly unusual, which is unprecedented in type-II superconductors. Recently, this phenomenon has been discussed in terms of the quantum transport of the quasiparticles in ultraclean superconductors [26].

3. Superconducting gap structure

3.1. How do we determine the gap structure?

Since the thermal conductivity is governed by the hole band with small gap in the low-field regime (I), while it is governed by the electron band with large gap in the high-field regime (II), we can distinguish the nodal topology in each band by analyzing the thermal conductivity in regimes (I) and (II) separately.

In fully gapped superconductors in magnetic fields, all the quasiparticle states are bound to vortex cores. Therefore, the quasiparticle DOS $N(0)$ increases linearly with $H$, $N(0) \propto H$, in the vortex state. In sharp contrast, the DOS in nodal superconductors is dominated by delocalized near-nodal quasiparticles [43]. Applied field creates a circulating supercurrent flow $v_s(r)$ associated with vortices. The Doppler shift of the energy of a quasiparticle with momentum $p$, $E(p) \rightarrow E(p) - v_s \cdot p$, is important near the nodes, where the local energy gap is small, $\Delta(k) < |v_s \cdot p|$. The Doppler shift induces a remarkable field dependence of $N(0)$. For line nodes, $N(0)$ increases as $N(0) \propto \sqrt{H}$. A remarkable difference between the fully gapped and nodal superconductors appears in the field dependence of the thermal conduction [31]. In the fully gapped superconductors, the applied magnetic field hardly affects the thermal conduction except in the vicinity of $H_{c2}$, because the bound quasiparticles cannot carry the heat. On the other hand, in nodal superconductors, the thermal conductivity increases steeply with $H$ even at low field due to the delocalized quasiparticles induced by the Doppler shift.

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Another remarkable feature is that the Doppler shifted quasiparticle DOS depends sensitively on the angle between the magnetic field and nodal directions, as displayed in figure 4(a). By using this property, recent measurements of thermal conductivity and the heat capacity with $H$ applied at varying orientation relative to the crystal axes have established determining the superconducting gap structure in $k$-space in several systems [31], [44]–[47]. Figure 4(b) illustrates the angular variation of the DOS for three types of nodes when $H$ is rotated crossing the nodes: (i) horizontal line node, (ii) vertical line node and (iii) point node which are relevant to the present study. The DOS shows a minimum when $H$ is applied in
the nodal direction. Clear twofold or fourfold oscillations of thermal conductivity and heat capacity associated with the nodes have been reported in YBa$_2$Cu$_3$O$_7$ [48], CeCoIn$_5$ [42, 49, 50], CeRhIn$_5$ [51], CeIrIn$_5$ [52], Y(Lu)Ni$_2$B$_2$C [53, 54], $\kappa$(BEDT-TTF)$_2$ Cu(NCS)$_2$ [55], PrOs$_4$Sb$_{12}$ [56] and UPd$_2$Al$_3$ [57], when $H$ is rotated relative to the crystal axes.

3.2. Gap structure of spherical light hole band

We first discuss the thermal conductivity in regime (I), which is governed by the spherical light hole band. In this regime, $\kappa(H)/T$ is nearly isotropic with respect to the field direction in spite of the large anisotropic $H_{c2}$. It should be stressed that this is consistent with the spherical structure of the hole band. The steep increase of $\kappa(H)/T$ at low fields is in strong contrast to the behavior of fully gapped superconductors. For both field directions, $\kappa(H)/T \propto N(0)$ is nearly proportional to $\sqrt{H}$ in regime (I). The steep increase of $\kappa/T$, along with the finite $\kappa_{00}/T$ term, leads us to conclude the presence of line nodes in the spherical light hole band.

The next important question is the nodal topology; what is the position of the line nodes, parallel (type (i) in figure 4(b)) or perpendicular (type (ii)) to the $ab$ plane? Figure 5(a) shows the angular variation of the thermal conductivity $\Delta\kappa(\phi)/\kappa_n$ with $H$ rotated within the basal $ab$ plane at 0.05 T well below the boundary of regimes (I) and (II), plotted as a function of the azimuthal angle $\phi \equiv (a, H)$. Here, $\kappa_n$ is the normal state thermal conductivity. $\kappa(\phi)$ exhibits a

![Figure 5](http://www.njp.org/)
maximum at $\phi = 0^\circ$ and minima at $\phi = \pm 90^\circ$. $\kappa (\phi)$ can be decomposed into three terms, $\kappa (\phi) = \kappa_0 + \kappa_2 \phi + \kappa_4 \phi$, where $\kappa_0$ is a $\phi$-independent term, and $\kappa_2 = C_2 \phi \cos 2\phi$ and $\kappa_4 = C_4 \phi \cos 4\phi$ are terms with twofold and fourfold symmetry, respectively, as shown by the solid lines. In the present geometry ($q \parallel a$ and $H \parallel ab$ plane), the $\kappa_2 \phi$ term arises from the difference between the quasiparticle transport parallel and perpendicular to the vortices.

As indicated by the dashed line in figure 5(a), the angular variation of $\kappa (\phi)$ cannot be fitted solely by the twofold term, indicating the presence of the fourfold term, which is shown in figure 5(b). The amplitude of $\kappa_4 \phi$ is strongly temperature dependent. Figure 6 shows the temperature dependence of the amplitude of $C_4 \phi / \kappa_n$. $C_4 \phi$ is zero above $T_c$. $C_4 \phi$ is strongly enhanced below $T_c$, decreases with decreasing $T$ after showing a maximum at $T \sim 0.5$ K and vanishes at very low temperatures. We stress that the vanishing of $C_4 \phi$ at low temperatures is in sharp contrast to that expected for superconductors with vertical nodes and appears to indicate a small fourfold modulation of the gap around the Fermi surface [58]. Thus we conclude that the horizontal line nodes are located within the $ab$ plane in the light hole bands.

3.3. Gap structure of heavy electron band

Next, we discuss the gap structure of heavy electron band, which governs the thermal conductivity in regime (II). The plateau-like behavior of $\kappa (H) / T$ up to the upper critical field for $H \parallel c$ indicates that these quasiparticles do not experience the Doppler shift. On the other hand, the steep increase of $\kappa (H) / T$ with convex curvature for $H \parallel a$ can be interpreted as the Doppler shift effect for the quasiparticles due to nodes. The fact that the Doppler shift only occurs for fields parallel to the $ab$ plane suggests that it originates from point nodes in the gap along the $c$-axis as these do not yield the Doppler shift for fields along the $c$-axis. In contrast, the line nodes give rise to the Doppler shift for any field direction [14, 31]. The point node along the $c$-axis is also pointed out by the magnetic field dependence of the heat capacity, in which $C(H)$ increases linearly for $H \parallel c$-axis, while $C(H)$ increases with convex curvature [30].

To obtain further insight into the nodal structure, we examine the thermal conductivity in $H$ rotating within the $ab$ plane and $ac$ plane in regime (II). Figure 7(a) shows $\kappa (\phi)$ at
Figure 7. (a) Thermal conductivity $\kappa(\phi)$ at 0.25 K in $H$ rotating within the $ab$ plane in regime (II) of figure 3, where the thermal conduction is governed by the heavy electron band. The dashed line at 3 T is the result of the fitting obtained by $\kappa(\phi) = \kappa_0 + \kappa_{2\phi} + \kappa_v$, where $\kappa_v$ is given in equation (1). (b) The angular variation of the thermal conductivity after the subtraction of $\kappa_{2\phi}$ and $\kappa_v$. 0.25 K in $H$ rotating within the $ab$ plane. $\kappa(\phi)$ exhibits twofold symmetry with a maximum at $\phi = 0^\circ$. At high magnetic fields, $\kappa(\phi)$ exhibits a distinct cusp at $\phi = 0$, where $H$ is parallel to the thermal current $q$. This peculiar angular dependence of $\kappa(\phi)$ can be attributed to the quasiparticle scattering by vortex lattice. In the present ultraclean URu$_2$Si$_2$ single crystals, the quasiparticle mean free path $\ell_v$ significantly exceeds the intervortex distance $a_0 \simeq \sqrt{\Phi_0/B}$. In this case, vortex scattering gives rise to the twofold oscillation with non-sinusoidal modulation. Assuming that the motion of quasiparticles is restricted to the direction of $q$, the quasiparticle mean free path due to vortex scattering $\ell_v$ scales to the effective intervortex distance along the heat current as $\ell_v(\phi) = \ell_{v0}/|\sin \phi|$. Then the total mean free path $\ell$ is estimated to be $\ell^{-1}(\phi) = \ell_0^{-1} + \ell_v^{-1}(\phi)$, where $\ell_0$ is the quasiparticle mean free path due to impurities and
\[ \ell_{v0} = \ell_v(\phi = 90^\circ). \] This contribution to the thermal conductivity is given as

\[ \kappa_v(\phi) = A \left( \frac{\ell_0 \ell_{v0}}{\ell_{v0} + \ell_0 |\sin \phi|} \right), \]  

where \( A \) is a constant. We try to fit the angular variation by \( \kappa(\phi) = \kappa_0 + \kappa_{2\phi} + \kappa_v \), assuming that \( \ell_0 = 1.2 \mu m \) and \( \ell_{v0} \) is a fitting parameter. Note that \( \kappa_{2\phi} \propto \cos \phi \) arises from the geometrical effect between the vortex lattice and thermal current. As shown by the dashed line at 3 T in figure 7(a), the fitting curve well reproduces the data. To check the validity of this analysis, the field dependence of \( \ell_{v0} \) obtained by the fitting is shown in figure 8. \( \ell_{v0} \) exhibits a \( \sqrt{H} \) dependence, indicating that \( \ell_{v0} \) is scaled well by \( a_0 \). \( \ell_{v0} \) is one order longer than \( a_0 \), indicating a weak scattering by vortices. This is a contrasting result with that of the heavy fermion \( \text{CeCoIn}_5 \), in which the quasiparticle mean free path becomes comparable to \( a_0 \) immediately after entering the mixed state [59]. We note that according to the recent experiments on ultraclean \( \text{URu}_2\text{Si}_2 \), the quasiparticles are little scattered by the vortex lattice due to the formation of the quasiparticle Bloch states [25].

Figure 7(b) shows the angular variation of the thermal conductivity \( \delta \kappa/\kappa_n \) obtained by subtracting the twofold oscillations, \( \kappa_{2\phi} + \kappa_v \), from the data. Within the resolution of 0.2%, \( \delta \kappa/\kappa_n \) shows no oscillation with respect to \( \phi \). Consistently, no fourfold oscillation is observed in the recent specific heat measurements under rotated \( H \) [30]. These results definitely indicate that neither line node nor point node are located within the \( ab \) plane.

Next, we move on to the measurements in \( H \) rotated within the \( bc \)-plane. Figure 9 depicts the angular variation of the thermal conductivity \( \Delta \kappa(\theta)/\kappa_n \), plotted as a function of the polar angle measured from the \( c \)-axis \( \theta \equiv (c,H) \) at \( T = 0.25 \) K. In this geometry, since \( H \) is always perpendicular to \( q \), twofold oscillation arising from the geometrical effect between vortex lattice and thermal current is absent. \( \Delta \kappa(\theta)/\kappa_n \) exhibits a characteristic angular variation with cusp-like peaks at \( \theta = \pm 90^\circ \) in addition to a broad peak at \( \theta = 0^\circ \). The broad peak at \( \theta = 0^\circ \) is
pronounced at higher fields and its amplitude decreases rapidly with decreasing \( H \). Such a broad peak has been reported in the heat capacity measurements in rotated \( H \) and is attributed to the large anisotropy of upper critical fields parallel and perpendicular to the \( ab \) plane. The cusp-like peak at \( \theta = \pm 90^\circ \), which is observed even at \( H = 0.05 \) T \((\sim 0.4\% \) of upper critical field for \( H \parallel ab \)), is consistent with the point node along the \( c \)-axis (see type (iii) in figure 4(b)). Based on these results we arrive at the conclusion that the point nodes are along the \( c \)-axis in the heavy electron bands.

4. Discussion

Summarizing the salient features of the superconducting properties in ultraclean URu$_2$Si$_2$ single crystals:

1. Superconductivity in semimetallic state: superconductivity with extremely low carrier number is realized. Electronic structure is characterized by nearly perfect compensation having the same number of electrons and holes.

2. Spin-singlet superconductivity: the first-order phase transition takes place at \( H_{c2} \), implying Pauli paramagnetically limited superconducting state.

3. Quantum quasiparticle transport: the enhanced entropy flow in the superconducting state indicates that the quasiparticle transport in the vortex state cannot be described by classical transport.

4. Multi-gap superconductivity: twofold steep increase with plateau-like behavior in the field dependences of thermal conductivity indicates the multigap superconductivity.
Table 1. The gap functions of even parity representations in the tetragonal crystal with point group D\textsubscript{4h}. The nodal structure in the simplest form allowed by symmetry is shown in the third column, where V and H denote vertical line node and horizontal line node, respectively.

| Representation | Basis function | Node          |
|----------------|----------------|---------------|
| A\textsubscript{1g} | $1, k_x^2 + k_y^2, k_z^2$ | Full gap     |
| A\textsubscript{2g} | $k_x k_y (k_x^2 - k_y^2)$ | Line (V)     |
| B\textsubscript{1g} | $k_x^2 - k_y^2$ | Line (V)     |
| B\textsubscript{2g} | $k_x, k_y$ | Line (V)     |
| E\textsubscript{g}(1, 0) | $k_z, k_x$ | Line (V+H)   |
| E\textsubscript{g}(1, 1) | $k_z, (k_x + k_y)$ | Line (V+H)   |
| E\textsubscript{g}(1, i) | $k_z (k_x + i k_y)$ | Line (H)+Point |

(5) Two distinct gap structures having different nodal topology: the temperature, field and field-angle dependence of the thermal conductivity reveal that the nodal topology is different in distinct bands; horizontal line nodes within the $ab$ plane in the spherical light hole band, and point nodes along the $c$-axis in the cigar-like heavy electron band.

The above results provide strong constraints on the possible gap function in URu\textsubscript{2}Si\textsubscript{2}. Table 1 shows the classification scheme for the gap functions of even parity representations in the tetragonal crystal with point group D\textsubscript{4h}. We show the nodal structure in each symmetry. The classification scheme allows one type of state having both line and point nodes simultaneously; an even parity state belonging to two-dimensional E\textsubscript{g} representation with the basic d-wave form, $\hat{k}_z (\hat{k}_x + i \hat{k}_y)$, which is called ‘chiral’ d-wave symmetry. This gap function is in some way reminiscent of that suggested in the spin triplet Sr\textsubscript{2}RuO\textsubscript{4}, $\hat{z} (\hat{k}_x + i \hat{k}_y)$ [60]. Due to the breaking of the time reversal symmetry, twofold degeneracy appears in the in-plane component. Considering the body-centered tetragonal crystal lattice structure of URu\textsubscript{2}Si\textsubscript{2}, the basic form of the gap function with the correct periodicity in the momentum space is given by

$$\Delta = \Delta_0 \sin \frac{k_z}{2} c \left( \sin \frac{k_x + k_y}{2} a + i \sin \frac{k_x - k_y}{2} a \right).$$

The gap function possesses line nodes at the center and at the boundary of the Brillouin zone ($k_z = 0, \pm 2\pi/c$), and point nodes along the symmetry axis ($k_x = k_y = 0$).

Although the detailed Fermi surface topology has not been established\textsuperscript{6}, we attempt to speculate a schematic form of the Fermi surface pockets and gap structure, which is displayed in figure 10. Horizontal line nodes lie on the hole band, which is consistent with the experimental observations, whereas both Fermi surfaces have point nodes along the $c$-axis. In avoiding an electron pocket centered at the $\Gamma$ point, the line node for $k_z = 0$ does not play any role. The suggested gap function may give rise to interesting superconducting phenomena due to the time

\textsuperscript{6} Very recently, the band structure calculation has been reported in the large-moment antiferromagnetic state under pressure [61]. In this study, however, we use the Brillouin zone in the paramagnetic phase for simplicity.
reversal symmetry breaking. Then spontaneous static magnetic fields arising from the orbital current around the impurity or at the surface may appear below $T_c$. Such a phenomenon has never been reported in the previous studies and search using an ultraclean system is a future issue.

Finally, we discuss the pairing interaction inferred from the suggested gap structure. Considering the magnetic fluctuations of the spin structure in the antiferromagnetic phase neighboring the superconducting phase under pressure (see figure 1(b)), it would naturally give rise to the suggested gap function. An interlayer singlet pairing is induced by an antiferromagnetic correlation between the U atoms in the neighboring basal planes. A ferromagnetic correlation within the basal plane suppresses the in-plane singlet pairing, giving rise to an in-plane degeneracy of the gap functions. Although the order parameter of the hidden order remains elusive, the determined gap function implies the importance of antiferromagnetic correlations in the pairing mechanism.

5. Conclusion

In summary, by performing the angle-resolved thermal transport measurements by using ultraclean single crystals of URu$_2$Si$_2$, we provide several pieces of evidence for a new type of unconventional superconductivity with two distinct gaps having different nodal topology. The results are consistent with a ‘chiral’ d-wave state with a form $\hat{k}_z(\hat{k}_x + i\hat{k}_y)$. The hidden order of URu$_2$Si$_2$ appears to provide an intriguing stage of a highly unusual superconductivity, which adds a unique example to the list of unconventional superconductors.
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