Nodal gap structure and order parameter symmetry of the unconventional superconductor UPt$_3$

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Abstract

Spanning a broad range of physical systems, complex symmetry breaking is widely recognized as a hallmark of competing interactions. This is exemplified in superfluid $^3$He which has multiple thermodynamic phases with spin and orbital quantum numbers $S = 1$ and $L = 1$, that emerge on cooling from a nearly ferromagnetic Fermi liquid. The heavy fermion compound UPt$_3$ exhibits similar behavior clearly manifest in its multiple superconducting phases. However, consensus as to its order parameter symmetry has remained elusive. Our small angle neutron scattering measurements indicate a linear temperature dependence of the London penetration depth characteristic of nodal structure of the order parameter. Our theoretical analysis is consistent with assignment of its symmetry to an $L = 3$ odd parity state for which one of the three thermodynamic phases in non-zero magnetic field is chiral.

1. Introduction

Recent interest in topological superconductors has focused attention on materials that exhibit chiral symmetry, or have been proposed to exhibit chiral symmetry, including Sr$_2$RuO$_4$, $^3$He, and UPt$_3$ [1]. In particular, the heavy fermion compound UPt$_3$ [2, 3] has attracted theoretical attention [4–6] in part as a consequence of conflicting experimental reports on the nature of its unconventional superconducting state. For example, the observation of Pauli limiting in the upper critical field [7] appears to be incompatible with temperature independence of the Knight shift [8, 9]. Josephson tunneling interference measurements [10] and measurements of the polar Kerr effect [11] provide evidence for an order parameter that is chiral in the $B$-phase. However, recent directional thermal conductivity experiments are interpreted otherwise [6]. Here we use small angle neutron scattering (SANS) from the vortex lattice (VL) to provide a bulk probe of the temperature dependence of the penetration depth, obtaining evidence for the nodal structure of the order parameter in the $B$-phase supporting its identification as an odd parity, chiral state, with $E_{2u}$ symmetry consistent with theory [12].

One of the most striking properties of UPt$_3$ is the fact that the $H – T$ superconducting phase diagram has three distinct superconducting vortex phases shown in figure 1(a), conventionally labeled A, B, and C. Experiments and theory demonstrate that this phase diagram can only be explained by an unconventional superconducting order parameter [3], a close parallel to superfluid $^3$He. However, a complete theoretical description of the superconducting state of UPt$_3$ has not been settled, and there are several candidate models that can account for the material’s unusual physical properties. The order parameter structure that is consistent with a number of experiments is an odd-parity, $f$-wave ($L = 3$) orbital state of $E_{2u}$ symmetry [4, 12]. However, with some success, comparisons with experiment have also been made for an even-parity, $d$-wave ($L = 2$) orbital state of $E_{1g}$ symmetry [15]. Both of these order parameters are chiral and break time reversal symmetry in the low
temperature $B$-phase in contrast to a recent proposal [6] for an odd-parity, $f$-wave ($L = 3$) model with $E_{1u}$ orbital symmetry which is non-chiral and time reversal symmetric in the $B$-phase.

All of these order parameters have nodes in the superconducting energy gap, each with different nodal structure in the three vortex phases. Consequently, it is of particular importance to explore physical properties that are directly linked to this nodal structure and that are sensitive to the gap dispersion at the nodes. Using SANS from the VL, we have measured the temperature dependence of the components of the London penetration depth, $\lambda_i (T)$, that probe the gap nodal structure along the principal directions of the crystal [14, 15] finding linear behavior in the low temperature limit. Our calculations using the quasiclassical Green's function approach and an ellipsoidal Fermi surface are consistent with the $\lambda_i (T)$ data over a wide range of temperature for an order parameter with $E_{2u}$ symmetry. We compare our results for $\lambda_i (T)$ with those from other experimental methods including ac-susceptibility [16, 17] and muon spin rotation ($\mu$SR) [18, 19].

2. Experimental methods

Our sample consists of a high-quality, 15 g single crystal ($\text{RRR} > 600$), cut into two pieces, and is described by Gannon et al [20]. The UPt$_3$ crystals were co-aligned, fixed with silver epoxy to a copper cold finger, and mounted to the mixing chamber of a dilution refrigerator with the crystal $a$-axis vertical and the $c$- and $a^*$-axes in the horizontal plane. Rotation of the dilution insert allowed easy reorientation of the $a^*$ or $c$-axes to be parallel to the magnetic field and neutron beam inside a horizontal superconducting magnet on the SANS-I and SANS-II beamlines at the Paul Scherrer Institut in Villigen, Switzerland. For measurements on SANS-I, the neutron wavelength was 6 Å with 11 m of collimation and detector to sample distance between 16 and 20 m. For measurements on SANS-II, 9 Å neutrons were used with 6 m of collimation and the detector to sample distance was 6 m.

3. Results

A typical result of a SANS diffraction pattern from UPt$_3$ is shown in figure 1(b) with magnetic field $H = 0.3$ T parallel to the $a^*$-axis. In the present work we have made measurements of many similar patterns as a function of temperature and magnetic field, figure 2. Only the first order Bragg reflections were observed as in figure 1(b), since UPt$_3$ has relatively long penetration depths. The diffraction patterns were constructed from a superposition of scattering images measured at different rocking angles $\phi$ as the sample and magnet are tipped about a horizontal axis perpendicular to the neutron beam. The image intensity at a given angle is shown in figure 3 for a typical experiment where the Bragg condition was satisfied above the beam center. By symmetry, there are two peaks below the beam center and 2 peaks on the horizontal, indicated by red circles. The UPt$_3$ crystal axes are shown as white arrows.

Figure 1. Phase diagram of UPt$_3$ and its vortex diffraction pattern. (a) A schematic of the phase diagram for the three vortex phases, $A$, $B$, and $C$, of UPt$_3$ for $H || a^*$. The normal to superconducting transition is $T_c$. The transitions between phases are $T_{AB}$ and $T_{BC}$. (b) An example of a diffraction pattern, measured at $\approx 50$ mK for $H = 0.3$ T with $H || a^*$. The Bragg condition was only satisfied for diffraction peaks above the beam center. By symmetry, there are two peaks below the beam center and 2 peaks on the horizontal, indicated by red circles. The UPt$_3$ crystal axes are shown as white arrows.
The symmetry of the diffraction pattern is the same as that of the real space VL, rotated by 90° with a rescaling of the axes. The distortion of the VL from a perfect hexagon is a result of penetration depth anisotropy in the plane perpendicular to $a^*$. When a 0.2 T field is applied parallel to the $c$-axis, a perfect hexagonal VL is seen within our resolution, in agreement with previous measurements for that orientation at a similar field [23].

Figure 2 shows the opening angle $2\alpha$ of the VL, defined as the angular separation of diffraction peaks displayed in the inset, as a function of applied magnetic field along the $a^*$-axis measured at $\approx 50$ mK. The $B$–$C$-phase transition is given by the vertical dashed line at $H = 0.6$ T. Red lines are guides to the eye.

Our $2\alpha$ data can be best described as having a linear field dependence in the $B$-phase that becomes field independent in the $C$-phase, where the $B$–$C$-transition occurs between $H = 0.5$ and 0.6 T in this field orientation, inferred from the phase diagram of Adenwalla et al [25]. The field dependence to the opening angle indicates that non-local corrections to the London theory are significant [12, 26, 27]. A change in the field dependence of the opening angle at the $B$–$C$-transition was also reported by Yaron et al [22].

The intensity in a diffraction peak is related to the Fourier transform of the local field variations from the real space VL. Systematic measurements as a function of rocking angle were made to produce rocking curves such as displayed in figure 3. We measured the first order Fourier component of the diffraction, called the form factor $|h_1|$, expressed as,

$$|h_1|^2 = R \frac{16 \Phi_0 q}{2 \pi^2 \lambda_n^2} \frac{1}{t},$$

The form factor is calculated from the reflectivity, $R$, equal to the integrated intensity of a rocking curve multiplied by $\cos \alpha$ (the Lorentz factor), divided by the incident neutron flux. In equation (1), $\Phi_0 = 2.07 \times 10^{-7}$ T · Å$^2$ is the magnetic flux quantum; $q$ is the magnitude of the scattering vector of the reflection being measured; the gyromagnetic ratio of the neutron is $\gamma = 1.91$; $\lambda_n$ is the incident neutron wavelength; and $t$ is
the effective sample thickness which we have taken to be 3.9 mm—the equivalent thickness of a uniform sample with the same width, height, and volume as our sample. The field dependence of our measurements of rocking curve widths do not show the sudden broadening at the B–C phase transition reported by Yaron et al [22]. All of our rocking curves are only ~20% broader than the resolution limit for our experiments. We also do not see a change in slope of the field dependence of |$h_1$| at the B–C transition as reported earlier [22]. It is likely that absence of these effects can be attributed to the higher quality of our crystal and the oscillatory field procedure which we have used to overcome flux pinning.

In the London theory the form factor is related to the material properties through the magnetic penetration depth $\lambda$. The form factor for an isotropic superconductor is given by,

$$|h_1| = \frac{B}{1 + \lambda^2q^2} e^{-\xi^2q^2},$$  

(2)

where $\xi$ is the superconducting coherence length and $c$ is a constant, typically taken to be $\frac{1}{2}$. The fractional part of equation (2) comes directly from the London equations [28]. The exponential factor is a correction to the London theory to account for the non-zero extent of the vortex cores [29, 30]. This simple gaussian model for the core correction with the constant $c = \frac{1}{2}$ has been found to be more accurate than more sophisticated models [30]. Nonetheless, the temperature dependence of the penetration depth is not sensitive to the choice of this correction and its exact value is immaterial to the conclusions in the present work. For an anisotropic superconductor the form factor can be expressed in terms of the principal values of the penetration depth $\lambda_i$ corresponding to currents flowing along each of the principal directions of the crystal, with $i = 3$ for currents along the $c$-axis.

Measuring the form factor therefore provides a direction-specific probe of the low lying excitations in the superconducting state sensitive to gap nodes [14].

For uniaxial anisotropy, as for UPt$_3$, $\lambda_1 = \lambda_2 \neq \lambda_3$ and the form factor for fields along the $a$ or $a^*$-axis becomes [28]

$$|h_1| = \frac{B}{1 + \lambda_i^2q^2 \sin^2 \alpha + \lambda_c^2q^2 \cos^2 \alpha} e^{-\xi^2q^2},$$  

(3)

where the $\lambda_i$ are related to the corresponding diagonal components of the quasi-particle mass tensor $m_0$, and the opening angle $2\alpha$, through the relation,

$$\tan^2 \alpha = \left(\frac{m_3}{3m_1}\right) = \left(\frac{\lambda_3^2}{3\lambda_1^2}\right).$$  

(4)

If there is no variation in the VL geometry as a function of temperature, as demonstrated in figure 5, then the temperature dependence of the form factor given by equation (3) reflects the temperature dependence of $\lambda_3$ where the denominator of equation (3) simplifies to $1 + \frac{\lambda_3^2}{3\lambda_1^2} q^2 \cos^2 \alpha$.

We have made measurements of the temperature dependence of the VL scattering for magnetic fields along both the crystal $c$ and $a^*$-axes with the field reduced from above $H_{c2}$, followed by damped field oscillations before measurement at each temperature. Rocking curves were obtained for each orientation at base temperature and at intermediate temperatures to determine that there was no broadening as temperature was varied. The magnet and sample were rotated to the center of the rocking curve and the scattered intensity $I(T)$ was measured ‘rocked-on’ ($\phi = \phi_0$) as a function of temperature, figure 4.

The opening angle and the scattering vector, taken directly from the diffraction pattern, are both temperature independent as shown in figures 5(a) and (b). The penetration depth anisotropy at low temperatures obtained directly from the opening angle is, $\lambda_1/\lambda_3 = 1.83 \pm 0.04$ at $H = 0.2$ T, giving a quasiparticle mass anisotropy of $m_3/m_1 = 3.34 \pm 0.13$. Using the average values of $\alpha$ and $q$, we calculated $|h_1|$ from equation (1) at $H = 0.2$ T for each temperature and field orientation. We determined $\lambda_3(T)$, shown in figure 6 from the simplified version of equation (3) using our $|h_1|$ values for $H || a^*$, the average values of $\alpha$ and $q$ for this field orientation, and $\xi = 110$ Å [7]. Since the penetration depth is isotropic in the plane perpendicular to the $c$-axis, we used equation (2) and our results for $|h_1|$ with $H || c$ to find $\lambda_3(T)$.

4. Penetration depth at low temperature

The nodal structure of the order parameter is evident from the VL scattering cross-section in its low temperature limiting behavior where the quasiparticle thermal excitation energies are much less than $k_B T_C$. From earlier work, notably thermal conductivity and sound attenuation [3] together with the theory [4, 12], this limiting low temperature region is $T/T_C \lesssim 0.4$ which we conservatively take to be $T/T_C \lesssim 0.3$. We have compared linear and quadratic fits to the temperature dependence of $\lambda_3$ over this temperature range shown in detail in figure 6(b).
Our data is consistent with a linear temperature dependence which provides a significantly better description than quadratic behavior as indicated by our chi-squared analyzes for the fits shown in this figure.

For $\lambda_1(T)$, i.e. $H \parallel c$, the accuracy of the data is less than for $\lambda_3(T)$ since the corresponding penetration depth is larger and the spatial variations of the local magnetic field from which the neutrons are scattered are much smaller. However, we have independent information from the diffraction pattern resident in our measurement of the opening angle $\alpha(T)$. Within the context of the London theory we can determine $\lambda_1(T)$ from equation (4). We plot this determination of $\lambda_1(T)$ in figure 6(a) as a blue line, which is also linear in temperature just as is $\lambda_3(T)$. Our extrapolations to zero temperature with linear fits to the data give: $\lambda_1(0) = 6, 800 \pm 210$ Å and $\lambda_3(0) = 3, 920 \pm 60$ Å.

5. Theoretical analysis

To interpret our data in terms of the pairing symmetry of UPt$_3$, we provide a brief discussion of the nodal structures of the superconducting gap. Gap profiles in the $B$-phase for various candidate models for the symmetry of the order parameter are shown on an ellipsoidal Fermi surface in the inset to figure 7(b). In the low field and low temperature $B$-phase, where all of the data shown in figure 6 were measured, the three predominate pairing models discussed earlier have three different nodal structures. For the $E_{2u}$ model [12], there are point nodes at the poles of the Fermi surface which open with quadratic wave-vector dispersion, and a line node around the equator of the Fermi surface that opens with linear dispersion. The $E_{1g}$ model [13] also has point
Figure 6. The temperature dependence of the penetration depth. (a) $\lambda_1$ calculated with equation (2) (green squares) and $\lambda_3$ (red circles, from equation (3)) over the whole temperature range with linear fits to each (green and red lines) from base temperature to $T/T_c = 0.3$ in a magnetic field $H = 0.2$ T. The blue line is derived from analysis of $\lambda_1(T)$ together with the opening angle $\alpha(T)$ in the context of the London theory, equation (4). The statistical accuracy of the measurements for $\lambda_3$ is approximately the size of the data points. (b) $\lambda_3$ in the low temperature region with linear (red) and quadratic (yellow) power law fits. The linear fit is significantly better with $\chi^2$ favoring linear temperature dependence by a factor of 2.23 as compared with a quadratic fit. Additionally, the temperature dependence over a broad range is consistent with our theoretical analysis, figure 7 for quadratically dispersed point nodes along the $c$-axis that leads to a linear temperature dependence at low temperatures.

Figure 7. Theoretical calculation of the penetration depth. Comparison is made with three models for the symmetry of the order parameter $E_{2u}$, $E_{1g}$, and $E_{1u}$. (a) The calculations of $\lambda_1(T)/\lambda_1(0)$ as a function of temperature are shown for each model order parameter symmetry and an ellipsoidal Fermi surface, (b) inset. The open circles in (a) are data from figure 6(a) labeled $\lambda_1$. An independent and more accurate data set for $\lambda_1$ (solid circles) was obtained from the data for $\lambda_3$ combined with measurements of the opening angle as described in the text, equation (4). (b) The calculations of $\lambda_3(T)/\lambda_3(0)$ give the temperature dependence of the penetration depth for currents along the $c$-axis. The linear behavior at low temperature for the $E_{2u}$ state is a consequence of the quadratic dispersion of the energy gap for the polar nodes. Solid curves are for $E_{2u}$ (blue), $E_{1g}$ (green), and $E_{1u}$ (yellow). Dashed lines show results for different nodal openings $\mu_1 (E_{1u})$ and $\mu_2 (E_{2u})$. (b Inset) Gap profiles for the three candidate order parameters in the B-phase on an ellipsoidal Fermi surface.
nodes at the poles, however these nodes open linearly. Similar to \( E_{2u} \), the \( E_{1u} \) model also has a line node around the equator that opens linearly. The \( E_{1u} \) model [6] has a somewhat more complicated gap structure in the \( B \)-phase, with point nodes at the poles that have linear dispersion and two line nodes in planes parallel to the equator where there is an antinode. Our measurements of \( \lambda \), test the nodal structure on parts of the Fermi surface having a significant basal plane component of the Fermi velocity, while measurements of \( \lambda \), are sensitive to the nodes where the Fermi velocity has a large \( c \)-axis component.

For a polar point node with quadratic dispersion, a linear temperature dependence of \( \lambda \) is expected in the low temperature limit, while for a point node with linear dispersion, there would be a \( T^2 \) temperature dependence. Extending analysis to a wider range of temperature and interpretation of our results in terms of order parameter symmetry requires a theoretical calculation including the effects of thermal excitations of quasiparticles averaged over the whole Fermi surface.

We performed calculations of the penetration depth within the framework of the quasiclassical theory [31, 32], and compared three models for the symmetry of the order parameter \( E_{2u}, E_{1u} \), and \( E_{1u} \) in figures 7(a) and (b). The calculations of superfluid density \( \rho_s(T) \) and penetration depth \( \lambda_{1,3}(T)/\lambda_{1,3}(0) = \sqrt{\rho_{1,3}(0)/\rho_{1,3}(T)} \) were performed for an ellipsoidal Fermi surface \( p_x^2 + p_y^2 + 3p_z^2 = p_0^2 \), with mass anisotropy \( m_1/m_3 = 3 \) to account for the observed anisotropy of the penetration depth \( \lambda_{1}(0)^2/\lambda_{3}(0)^2 \approx 3.3 \), and the normal state transport, \( \kappa_c(T)/\kappa_\alpha(T) \approx 2.8 \) [33]. The choice for the form of the order parameter is less obvious, and there are several approaches to model the gap structure using standard functions, such as spherical functions, or Allen Fermi surface harmonics [34]. We made the more natural choice of ellipsoidal harmonics since they are orthogonal on an ellipsoidal Fermi surface, and transform into spherical harmonics with proper rescaling of the Fermi surface. Substituting \( (\rho_s, \rho_t, \rho_b) = (k_x, k_y, k_z/\sqrt{3}) \) we write the gap profiles for the three models, displayed in the inset to figure 7(b). We used a single-component model for the order parameter, valid deep inside the \( B \)-phase and followed the approach from previous work [34], normalizing the temperature to the lower critical temperature \( T_{nl} \approx 0.88 T_c \).

For this choice of gap functions, and treating \( \lambda(0) \) as the only adjustable quantity, the theory with \( E_{2u} \), order parameter symmetry closely replicates the observed \( \lambda_3(T)/\lambda_3(0) \) over a wide range of temperature. The best fit is shown in figure 7(b) by a solid blue line, whereas the \( E_{1u} \) and \( E_{1u} \) models with ellipsoidal harmonics provide considerably worse fits. The \( E_{2u} \) model is also consistent with the data for \( \lambda_1 \) (solid circles) in figure 7(a).

It was pointed out that ‘pure’ ellipsoidal or spherical harmonics might not necessarily reflect the correct low-energy structure of the excitations, and do not replicate the observed anisotropy of the heat transport, \( \kappa_c(T)/\kappa_\alpha(T) \) [35]. A set of gap functions was suggested that are parametrized near line and point nodes with variable slope coefficients \( \mu \) for angles \( \delta \) with respect to the \( c \)-axis.

\[
\Delta E_{2u} = \Delta_0 \left| k_x \left( k_x + i k_y \right)^2 \right|
\]
\[
\Delta E_{1u} = \Delta_0 \left| k_z \left( k_x + i k_y \right) \right|
\]
\[
\Delta E_{1u} = \Delta_0 \left| 5k_x^2 - p_0^2 \right| \sqrt{k_x^2 + k_y^2}
\]

(5)

equation (5) corresponds to the opening nodal parameters: \( \mu = \mu_1 = \mu_2 = 1 \). However these authors [35] only found a good fit for the thermal conductivity data and sound attenuation with \( \mu = 1, \mu_1 = 1/3, \mu_2 = 2 \) [4, 35].

The latter parameter set results in dashed lines in figure 7(b). With this ansatz the \( E_{2u} \) model cannot fit the data over the entire temperature range, while \( E_{1u} \) follows the observed data fairly well within error bars but with significantly higher \( \chi^2 \) value at lower temperatures, as discussed in figure 6(b). This model is also a less likely candidate based on previous analysis of sound attenuation [4]. In summary, it is compelling that \( E_{2u} \) symmetry with a simple parameter set and elliptical harmonics is in excellent agreement with our measured penetration depths over a wide temperature range.

6. Discussion and summary

Signore et al [17] reported a linear temperature behavior which could not be associated with any specific component of the penetration depth. Interpretation of their ac-susceptibility measurements requires an analysis of the real and imaginary parts of the electromagnetic response from which extraction of the penetration depth is not trivial and is necessarily sensitive to surface quality [16, 17]. An early \( \mu \)SR investigation by Broholm et al [18]
found a penetration depth anisotropy much too small to be consistent with other observations of the superconducting state [3, 21, 22]. In a later μSR study, Yauaucn et al. [19] obtained \( \lambda_c(0) = 6, 040 \pm 130 \, \text{Å} \) and \( \lambda_s(0) = 4, 260 \pm 150 \, \text{Å} \) with \( H = 0.018 \, \text{T} \), qualitatively consistent with what we report here.

Evidence for gap nodes has been sought from the low temperature behavior of the thermal conductivity and attenuation of sound [4, 33–38]. The earliest reports [33], provided evidence for both a polar gap node along the \( c \)-axis and a line node in the basal plane. However, a conclusion in terms of a specific order parameter symmetry from nodal gap quasiparticle excitations was not possible [34, 35]. At high temperatures in the \( A \)-phase, measurements of transverse sound attenuation [38], VL structure [23], and directional tunneling [39] are consistent with \( E_{2u} \) symmetry. In contrast, a recent report of the directional dependence of the thermal conductivity in the \( B \)-phase was argued to support a \( E_{1u} \) state [6]. This theory requires weak spin–orbit coupling in order to maintain consistency with spin susceptibility measurements from the Knight shift. The latter is in conflict with observations of Pauli limiting anisotropy evidenced in the upper critical field [7] and it is in conflict with most other theoretical work [3, 4]. Our approach has been to use SANS to measure the vortex structure from which we have determined the penetration depth. These measurements are not compromised by imperfections at the sample surface since they are an average over the whole superconducting crystal and they provide absolute values for the penetration depth. The interpretation of transport measurements makes an assumption for the existence of a single order parameter domain that is not required for our measurements of the penetration depth from which we infer that superconductivity in UPt3 is an odd parity state with \( E_{2u} \) symmetry and that consequently, the \( B \)-phase is chiral.

Acknowledgments

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References

[1] Sauls J A and Eschrig M 2009 New J. Phys. 11 075008
[2] Stewart G R, Fisk Z, Willis J O and Smith J L 1984 Phys. Rev. Lett. 52 679
[3] Joynt R and Tallifer L 2002 Rev. Mod. Phys. 74 235
[4] Graf M J, Yip S K and Sauls J A 2000 Phys. Rev. B 62 14393
[5] Norman M R 2013 Unconventional superconductivity Novel Superfluids (International Series of Monographs on Physics vol 2) ed K H Bennemann and J B Ketterson (Oxford: Oxford University Press) ch 12
[6] Tsutsumi Y, Machida K, Ohmi T and Ozaki M 2012 J. Phys. Soc. Japan 81 074717
[7] Shivaram B S, Rosenbaum T F and Hinks D G 1986 Phys. Rev. Lett. 57 1259
[8] Tou H, Kitaoka Y, Asayama K, Kimura N, Onuki Y, Yamamoto E and Maezawa K 1996 Phys. Rev. Lett. 77 1374
[9] Tou H, Kitaoka Y, Ishida K, Asayama K, Kimura N, Onuki Y, Yamamoto E, Haga Y and Maezawa K 1998 Phys. Rev. Lett. 80 3129
[10] Strand J D, Harlingen D J V, Kycia J B and Halperin W P 2009 Phys. Rev. Lett. 103 197002
[11] Schemm E R, Gannon W J, Wishne C M, Halperin W P and Kapitulnik A 2014 Science 345 190
[12] Sauls J A 1994 Adv. Phys. 43 113
[13] Park K A and Joynt R 1996 Phys. Rev. B 53 12346
[14] Prozorov R and Giannetta R W 2006 Supercond. Sci. Technol. 19 R41
[15] Kawano-Furukawa H et al 2011 Phys. Rev. B 84 024507
[16] Gross–Alltag F, Chandrasekhar B S, Einzel D, Hirschfeld P J and Andres K 1991 Z. Phys. B 52 243
[17] Signore P J, Andraka B, Meisel M W, Brown S E, Fisk Z, Giorgi A L, Smith J L, Gross-Alltag F, Schubert E A and Menovsky A A 1995 Phys. Rev. B 52 4446
[18] Broholm C, Aeppli G, Kleiman R N, Harshman D R, Bishop D J, Bucher E, Williams D L, Ansaldo E J and Heffner R H 1990 Phys. Rev. Lett. 65 2062
[19] Yauaucn A, de Röétier P D, Huxley A, Flouquet J, Bonville P, Grubenn P C M and Mulders A M 1998 J. Phys.: Condens. Matter 10 9791
[20] Gannon W J, Halperin W P, Rastovski C, Eskildsen M R, Dai P and Stunault A 2012 Phys. Rev. B 86 104510
[21] Kleiman R N, Broholm C, Aeppli G, Bucher E, Štúchel N, Bishop D J, Clausen K N, Mortensen K, Pedersen J S and Howard B 1992 Phys. Rev. Lett. 69 1120
[22] Yaron U, Gammel P L, Boebinger G S, Aeppli G, Schiffer P, Bucher E, Bishop D J, Broholm C and Mortensen K 1997 Phys. Rev. Lett. 78 3185
[23] Huxley A, Rodière P, Paul D McK, VanDijk N, Cubitt R and Flouquet J 2000 Nature 406 160
[24] Das P, Rastovski C, O’Brien T R, Schlesinger K J, Dewhurst C D, DeBeer-Schmitt L, Zhigadlo N D, Karpinski J and Eskildsen M R 2012 Phys. Rev. Lett. 108 167001
[25] Adenwalla S, Lin S W, Ran Q Z, Zhao Z, Ketterson J B, Sauls J A, Taillefer L, Hinks D G, Levy M and Sarma B K 1990 Phys. Rev. Lett. 65 2298
[26] White J S et al 2011 Phys. Rev. B 84 104519
[27] Sauls J A 2013 private communication
[28] Kogan V G 1981 Phys. Lett. A 85 298
[29] Yaouanc A, de Réotier P D and Brandt E H 1997 Phys. Rev. B 55 11107
[30] Densmore J M et al 2009 Phys. Rev. B 79 174522
[31] Serene J W and Rainer D 1983 Phys. Rep. 101 221
[32] Xu D, Yip S K and Sauls J A 1995 Phys. Rev. B 51 16233
[33] Lussier B, Ellman B and Taillefer L 1994 Phys. Rev. Lett. 73 3294
[34] Norman M and Hirschfeld P 1996 Phys. Rev. B 53 5706
[35] Graf M, Yip S K and Sauls J 1996 J. Low. Temp. Phys. 102 367
[36] Lussier B, Ellman B and Taillefer L 1996 Phys. Rev. B 53 5145
[37] Suderow H, Brison J P, Huxley A and Floquet J 1997 J. Low Temp. Phys. 108 11
[38] Ellman B, Taillefer L and Poirier M 1996 Phys. Rev. B 54 9043
[39] Strand J D, Bahr D J, Harlingen D J V, Davis J P, Gannon W J and Halperin W P 2010 Science 328 1368