Low Field Scaling of the Flux-Flow Resistivity in the Unconventional Superconductor UPt_3

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Measurements of the flux-flow resistivity in the unconventional superconductor UPt_3 are reported for a large range of magnetic field. In agreement with a recent theory, ρ_f at low field is far larger than that found in conventional superconductors. Its field dependence at different temperatures shows a predicted scaling relation for clean superconductors. The crossover from localized to delocalized quasiparticle excitations around the vortex is also observed as the magnetic field increases.

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Since the flux-flow resistivity ρ_f in type-II superconductors is dominated by quasiparticle excitations in and around vortex cores, particular behaviors are expected for unconventional superconductors with nodes in the superconducting gap. However, the usual models for dirty superconductors, i.e., the normal core model [1] and the time-dependent Ginzburg-Landau (TDGL) model [2,3] cannot be applied to unconventional superconductors except under restricted conditions (e.g., $H \sim H_{c2}$) since unconventional superconductors are usually in the clean limit [4], and ρ_f is particularly sensitive to the nodes in gap for the clean case.

Recently, Kopnin *et al.* have developed a kinetic equation (KE) theory for clean superconductors including the unconventional case [5,6]. They have found that ρ_f obeys a particular scaling law for $H \ll H_{c2}$ when a localized quasiparticle excitation contribution is dominant. The scaling behavior of ρ_f is classified into two cases, i.e., the moderately clean case for $l > \xi$ (where l is the electronic mean-free path and ξ is the superconducting coherence length) and the superclean case for $l \gg \xi E_F/\Delta(T)$ [$\Delta(T)$ is the superconducting gap amplitude] [5]. The scaling relation for the moderately clean case (situation of UPt_3) is

$$r \approx \alpha[T_c/\Delta(T)]h\,, \tag{1}$$

where $r \equiv \rho_f/\rho_n$ (ρ_n is the resistivity of the normal state) and $h \equiv H/H_{c2}$. The factor α is inversely proportional to the Fermi surface average ($\langle\ \rangle_{FS}$) of the quantity $\omega(k) \approx \Delta(k)^2/E_F$ known as the vortex minigap and related to the energy spacing of the localized excitations. For the case of a nonisotropic gap, $\omega(k)$ depends on k and $\omega(k) \sim 0$ at the nodes so that $\alpha \sim \omega_0/\langle\omega(k)\rangle_{FS}$ can be larger than 1 [ω_0 is the maximum value of $\omega(k)$ over the Fermi surface], whereas $\alpha \sim 1$ for an isotropic s-wave superconductor.

Few measurements of ρ_f have been performed in conventional moderately clean cases [7,8], and even then ρ_f at low field far below H_{c2} has not been determined due to defects which induce the pinning force on the flux lines. The same difficulty is quite serious in high T_c oxides (unconventional superclean case) [9]. No measurements have been done in unconventional moderately clean cases. Therefore the predicted scaling law [Eq. (1)] has not been observed experimentally to date. The heavy fermion superconductor UPt_3 is an excellent candidate to test the above predictions. First, a high quality single crystal with quite a few defects can be obtained. In fact, the pinning force F_p of our crystal is on the order of $\sim 10^4\ N/m^3$, which is at least 20 times smaller than that of Nb [7] and $NbSe_2$ [8] single crystals. Then the pinning effect which disturbs measurements of ρ_f is not important here. In UPt_3, three distinct superconducting phases as a function of temperature (T) and field (H) are observed, i.e., the A phase at (high T, low H), the B phase at (low T, low H), and the C phase at (low T, high H), suggesting strongly the unconventional superconducting state [4]. Measurements such as thermal conductivity have led to the consensus that the superconducting gap has symmetry enforced nodes along certain crystallographic directions [10,11]. From typical values of our crystal [$l \approx 5000$ Å, $\xi \approx 100$ Å, $\Delta(0) \approx 1$ K, and an effective Fermi energy $E_F \approx 20$ K [12]], UPt_3 is in the moderately clean case. In this study for the unconventional superconductor UPt_3, the predicted behavior by the KE theory is clearly observed.

Our measurements were made on a single crystal ($T_c = 0.53$ K, the normal state residual resistivity $\rho_0 = 0.20\ \mu\Omega$ cm). The current was applied parallel to the c direction along the length of the sample. The magnetic field was applied perpendicular to a crystallographic c direction for which $H_{c2}(0) = 3$ T. One end of the sample was clamped mechanically onto the mixing chamber of a dilution refrigerator. An otherwise thermally isolated thermometer was connected to the sample by a gold wire to observe and permit the regulation of the real temperature of the sample. The standard four contact method was used to measure the resistance of the sample which had a cross section of 0.04 mm^2 with a distance of 1 mm between the voltage contacts. The contacts to the sample were made by spot welding gold wires with a micromanipulator. The applied current $I = I_{dc} + I_{ac}$ consisted of a dc part I_{dc} sufficient to drive the system into the flux-flow state and a small ac current I_{ac} of amplitude 0.01 mA at a frequency of 17 Hz. By detecting

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the voltage at 17 Hz with a phase sensitive detector, dV/dI was measured directly as a function of I_{dc}. The I vs V relation could be obtained by integrating dV/dI with respect to I_{dc}. The flux-flow resistivity is defined as $\rho_f = s * dV/dI$ (s is the geometric factor of our sample $s = 0.04$ mm) over the wide linear region of the I vs V curve above the superconducting critical current I_c. The power dissipated due to the direct joule heating of the sample and the contacts was around 50 nW for $I_{dc} = 10$ mA. The sample temperature was therefore higher than that of the mixing chamber. Based on our previous determination of the thermal conductivity of UPt_3 in the mixed state [10], we calculate that any temperature difference across the sample was less than 2 mK. A temperature difference occurred between the sample and the mixing chamber. Its magnitude was found to be consistent with that deduced from the estimated thermal conductance of the contact determined from the Wiedemann-Franz law and the electrical contact resistance. Since we measured and regulated the temperature of the sample directly, we avoided the significant error that would occur if only the temperature of the mixing chamber was available.

Typical data for $T = 0.35$ K in the form I_{dc} vs V are shown in Fig. 1. At other temperatures, similar behaviors have been observed. A linear response above I_c characteristic of the pinning-free flux-flow state is found over a wide range of currents, which suggests that ρ_f is indeed independent of the pinning character. Although there is a significant temperature dependence of the normal state resistivity [10], the measured I-V curve is linear above H_{c2} ($H = 1.1$ T), which confirms that the sample temperature is indeed correctly regulated and independent of I_{dc}. For $T < 0.4$ K, our sample shows a mild peak in I_c for fields near H_{c2}. This "peak effect" which we have reported previously [13] is a characteristic of the pinning and is similar to effects observed in several other type-II superconductors. The important point is that ρ_f shows no corresponding anomaly and increases monotonically with increasing field over this region. This confirms our assertion that ρ_f is not sensitive to details of the flux-lattice pinning. This conclusion is further supported by the fact that ρ_f was not found to depend at all on the field history nor samples which have different critical currents. In addition, ρ_f was identical for the samples with different surfaces prepared by polishing, indicating the surface effect is not important here.

Figure 2 shows the reduced field, $h = H/H_{c2}$, dependence of flux-flow resistivity $r = \rho_f/\rho_n$. It is remarkable that r is always larger than h. To our knowledge, such a behavior has never been observed in a conventional superconductor. As shown in Fig. 2, it cannot be explained by the normal core model [1] or by the TDGL model for dirty superconductors [3]. As we discuss below, this observed highly resistive behavior corresponds to a large α, which is naturally explained by considering the presence of nodes in the superconducting gap.

A careful examination of the data suggests that r undergoes a small discontinuous change around the $A-B$ transition at $T \sim 0.45$ K in intermediate fields as illustrated in the inset of Fig. 2. This may correspond to the change of symmetry of the gap at the transition [4]. At lower and higher field the experimental resolution is perhaps not sufficient to observe the change. No anomaly in ρ_f at the transition between the B and C phases (e.g., at $h \sim 0.5$ at 0.35 K) could be detected within the resolution of the present measurements. As reported previously [14], this

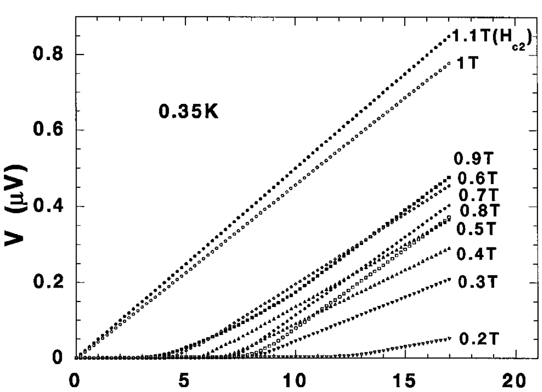

FIG. 1. Voltage vs current (I_{dc}) curves, in different magnetic fields (H), measured at 0.35 K in the mixed superconducting state ($H < 1.1$ T) and normal state ($H = 1.1$ T) of a single crystal of UPt_3 ($j \parallel c$ axis, $H \perp c$ axis).

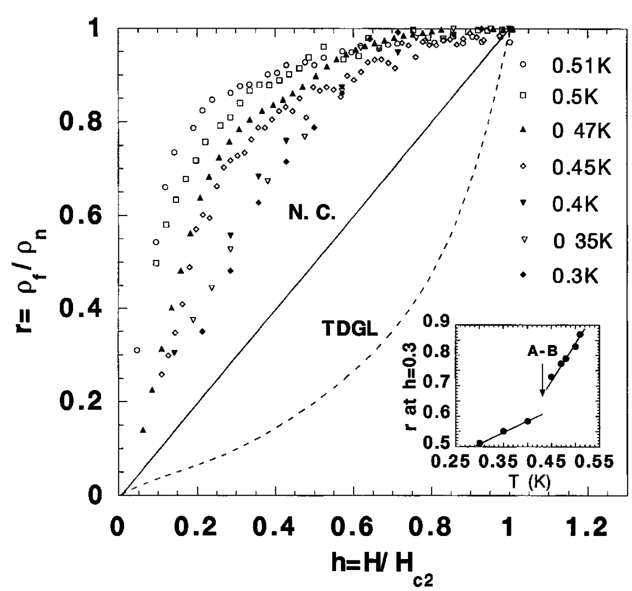

FIG. 2. The flux-flow resistivity ρ_f divided by the normal state resistivity ρ_n is plotted against the applied field normalized to $H_{c2}(T)$. Only a representative selection of the experimental points are presented for clarity. The lines represent calculations for (i) the normal core model (solid line) [1], and (ii) the TDGL (dirty case) model (dashed curve) [3]. The inset shows the temperature dependence of r at $h = 0.3$. The lines here are to guide the eye.

confirms that the change of physical properties at the $B-C$ transition is very small.

Following Eq. (1) for moderately clean superconductors, we plot the field dependence of r normalized by the superconducting gap function $\Delta(T)/\Delta(0)$ for the E_{2u} order parameter symmetry case in Fig. 3. The inset (a) of Fig. 3 shows the numerically calculated T dependence of the normalized superconducting gap amplitude $\Delta(T)/\Delta(0)$ for the frequently quoted model E_{2u} for the gap symmetry of UPt_3, as well as for the isotropic BCS case. As the difference in $\rho_f\Delta(T)/\rho_n\Delta(0)$ for the two models is insignificant compared to our experimental errors, we present only the result for the E_{2u} case in the main figure. Returning to Fig. 3 for $H \ll H_{c2}$, all the curves are seen to fall on a universal scaling line $[\rho_f\Delta(T)/\rho_n\Delta(0) \approx 1.7H/H_{c2}]$. The value of α is estimated to be 3.2 from this scaling relation if we use the plausible value of $\Delta(0) \approx 1.9T_c$ obtained from Andreev reflection measurements [15].

The large α (>1) suggests nodes in the gap or at least a strongly anisotropic gap. The previous indications for the nodes in UPt_3 are mainly the power law T dependence of physical properties (instead of the exponential law) [4]; however, such a behavior may be sensitive to impurity and particular T dependence of the normal state property. Moreover, the interpretation of these studies, other than in zero field or very close to the upper critical field H_{c2}, is complicated by sample dependent pinning of the flux lattice. In contrast, the large α here corresponds directly to the enhancement of quasiparticle excitations due to the nodes without any ambiguity, and is independent of the pinning since α is estimated in the free flux-flow state.

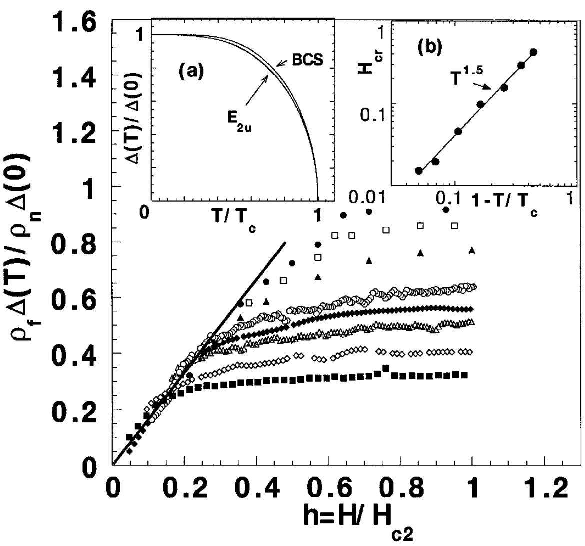

FIG. 3. The field dependence of the quantity $\rho_f\Delta(T)/\rho_n\Delta(0)$ for the E_{2u} case is shown. The different symbols correspond to different temperatures (descending from the top of the figure at $h \sim 1$, $T = 0.30$, 0.35, 0.40, 0.45, 0.47, 0.48, 0.50, and 0.51 K). The solid line represents the universal relation at low field $\rho_f\Delta(T)/\rho_n\Delta(0) = 1.7H/H_{c2}$. Inset (a) shows the explicit temperature dependence of $\Delta(T)/\Delta(0)$ for the E_{2u} and the BCS gap symmetries. The weak difference in the temperature dependence of the gap between the two models makes a negligible difference to the normalization of the data. Inset (b) shows the T dependence of the crossover field H_{cr} in tesla. The solid line represents the relation $H_{cr}/H_{c2}(0) \sim 0.5(1 - T/T_c)^{1.5}$.

Combining with the large α, the observed scaling confirms the KE prediction that the localized quasiparticle excitations around the nodes of the gap dominate ρ_f at low field in the clean case. It was unclear whether true localized excitations exist in the vortex cores for a superconductor with gap nodes. Indeed for the case of the $d_{x^2-y^2}$ symmetry state in the high T_c oxides, it is proposed theoretically that no true localized state exists in the vortex [16]; however, a small number of localized states have been observed by scanning tunneling microscopy measurements [17]. Obviously an experimental distinction between "true" and "strongly" localized states is difficult. At least our observation supports the strongly localized quasiparticles around the nodes for the A and B phases of UPt_3 ($H \perp c$ axis). This fact will favor the identification of the order parameter symmetry.

For larger fields above a crossover field H_{cr} which corresponds to the onset of a deviation from Eq. (1), the contribution of the delocalized excitations becomes comparable to that from the localized excitations considered above. This crossover is a predicted particular property of clean superconductors in the KE theory. The observed T dependence of H_{cr} obeys the relation $H_{cr}/H_{c2}(0) \sim 0.5(1 - T/T_c)^{1.5}$ at high temperatures [inset (b) of Fig. 3]. This relation is also naturally explained using the KE theory. From Eq. (1), we find $H/H_{c2}(0) \sim \frac{r\Delta(0)}{\alpha T_c}(1 - T/T_c)^{1.5}$ around T_c; thus the observed relation means that the crossover happens when r reaches a certain value r_{cr}. In the present case $r_{cr} \sim 0.84$ from $\frac{r_{cr}\Delta(0)}{\alpha T_c} \sim 0.5$, indicating that the localized excitation is dominant up to $\rho_f \sim 0.84\rho_n$ as expected in clean superconductors. The weakness of H_{cr} around T_c reflects the large value of α, i.e., the nodes in gap. From the KE theory, the observation of the relation also suggests that the conductivity due to the delocalized quasiparticle around vortex is proportional to $\omega_c\tau$ at low field as in the normal state (ω_c: cyclotron resonance frequency; τ: the relaxation time), whereas the conductivity due to the localized quasiparticle is proportional to $\langle\omega(k)\rangle\tau$. Consistently with the existence of these two characteristic energies, the cyclotron radius $l_H = v_F/\omega_c \sim 600$ Å for H_{cr} ($\sim$1.5 T at $T \ll T_c$) is smaller than the mean-free path $l \approx 5000$ Å (v_F is the Fermi velocity $\sim 3 \times 10^3$ m/s [12]).

At high magnetic field very near to H_{c2}, the observed ρ_f is interpreted by the TDGL model since in this field regime the delocalized excitations are dominant and the effect of the gap nodes can be ignored. In the TDGL model and its extension to the clean case for $H \simeq H_{c2}$ [18,19], the slope of ρ_f; dr/dh at $h = 1$ can be expressed

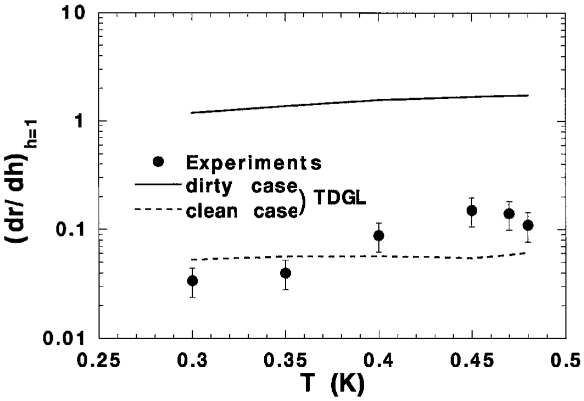

FIG. 4. The experimental temperature dependence of the gradient of $r = \rho_f/\rho_n$ with respect to $h = H/H_{c2}$ at H_{c2} is plotted as a function of temperature (circles). The calculated values for the TDGL model for the dirty and clean cases are shown as solid and dashed curves, respectively.

in dimensionless units for the clean case as

$$\left(\frac{dr}{dh}\right)_{h=1} = \frac{2m^* H_{c2}}{ne\hbar\pi[1.16(2\kappa_2^2 - 1) + 1]}, \tag{2}$$

while for dirty case

$$\left(\frac{dr}{dh}\right)_{h=1} = \frac{4\kappa_1^2}{[1.16(2\kappa_2^2 - 1) + 1]}, \tag{3}$$

where m^* is the effective mass and n is the superconducting electron density. In order to calculate $(dr/dh)_{h=1}$, the T dependence of m^*/n is estimated from the penetration depth λ [20] using the relation $\lambda^2 = m^* c^2/(4\pi n e^2)$. The estimated T dependence of the Maki-Ginzburg-Landau parameters κ_1 (evaluated at $T = 0$ K) and $\kappa_2(T)$ is deduced from specific heat measurements [21]. Experimentally, $(dr/dh)_{h=1}$ has been determined with a slow field sweep around H_{c2}. The experimental $(dr/dh)_{h=1}$ agree with the TDGL model for the clean case as expected, especially at low temperatures (Fig. 4). At high temperatures, the experimental points deviate slightly towards the dirty case [Eq. (3)], indicating perhaps that the clean case no longer describes the data well at $T \sim T_c$. A similar behavior has been observed in conventional clean case compounds such as pure V [22] and Nb [7].

In conclusion, we have achieved the free flux-flow state at low field in UPt_3 owing to the weak pinning effect of our high quality single crystal. The observed ρ_f is quite large compared to the value expected for conventional superconductors, but consistent with the recent kinetic equation theory for unconventional clean superconductors. The predicted scaling law for ρ_f with a large field slope is observed, confirming that ρ_f reflects the nodes of the gap and the localized quasiparticle excitations around the vortex in the low field (A and B) phases. The crossover from the localized ($H < H_{c2}$) to the delocalized ($H \sim H_{c2}$) excitation regime is observed. This crossover is a characteristic of clean superconductors and is interpreted consistently with the other properties in the superconducting and normal states using the same theory. Against the usual statement, the flux-flow resistivity can give important information of quasiparticle excitations in the mixed state of type-II superconductors.

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