Anisotropy of flux-flow resistivity in UPt$_3$

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Abstract

The ac penetration depth, $\lambda_{ac}(T, H, f)$, has been measured in superconducting UPt$_3$ single crystals for $H \parallel \hat{c}$ and $H \perp \hat{c}$ in the range $f$ = 0.01–1 MHz and $T$ = 0.1–0.4 K. The contributions from bulk and surface pinning have been separated to yield the flux-flow resistivity $\rho_f(H)$. With $H \perp \hat{c}$, $\rho_f$ displays magneto-resistance at low-field which agrees with previous dc measurements and the characteristic scaling law of clean crystals with anisotropic gap. When $H \parallel \hat{c}$, the low-field $\rho_f$ is three times larger. We interpret this property as evidence for flux lines with unconventional core structure.

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The reduced symmetry of the superconducting energy gap in unconventional superconductors is generally inferred from the presence of low-energy excitations, which exist due to nodes in the gap function $\Delta(\mathbf{k})$. In the most important heavy-fermion compound UPt$_3$, evidence for gap nodes has been deduced from measurements of the temperature dependences of the London penetration depth $\lambda(T)$ [1], ultrasonic attenuation [2], thermal conductivity [3], and most recently flux-flow resistivity $\rho_f(T)$ [4]. UPt$_3$ also has a unique phase diagram, with two low-field phases $A$ and $B$ and a high-field $C$-phase [5,6] (see insets in Fig. 2). These properties are compatible with a Ginzburg-Landau expansion in terms of a two-component order parameter [7–9]. Such states allow for new and unconventional vortex core structures as well as other topological defects, such as domain walls or textures [7–10]. A multitude of different vortex structures and other topological defects have been discovered in $^3$He superfluids [11], but so far direct experimental evidence for their existence in superconductors is sparse [12].

From neutron diffraction measurements it is known that the $B$ and $C$ phases in UPt$_3$ support distorted lattices of singly quantized vortex lines, with a clear break in the characteristic length scales at the $B \rightarrow C$ transition [13]. At low magnetic field parallel to the $c$-axis, the two-component order-parameter model predicts that three different core structures of singly quantized vortices are possible in the highly symmetric $B$-phase [14]. These have all singular structure with a radius comparable to the coherence length $\xi$. One of them is axisymmetric while the other two, called triangular and crescent cores, are nonaxisymmetric and in contrast to the former, here the order parameter amplitude does not go to zero anywhere within the core. The relative stabilities of these different structures depend on the Ginzburg-Landau parameters, which are still largely undetermined.

The investigation of the vortex-core structure in the bulk superconductor requires measurement of dynamical properties, like flux-flow dissipation or the Hall angle. We report on measurements of flux-flow resistivity $\rho_f(H,T)$ in the low-temperature limit ($T \lesssim 0.3$ K) for $H \perp \hat{c}$ and $H \parallel \hat{c}$. For $H \perp \hat{c}$ our results are consistent with those of Ref. [4] at $T \gtrsim 0.3$ K. When $H \parallel \hat{c}$, the resistivity proves to be much larger than expected. It does not appear possible to explain this anisotropy unless the vortex structure is different in the two crystal directions.

Vortex-dynamic measurements are complicated by pinning due to sample defects, which strongly alter the dc and low-frequency ac responses [14]. At low temperatures in the $B$-phase, the use of large dc currents, to depin vortex lines and to bias the system into the free-flow state, is prohibited due to excessive Joule heating. In the present ac measurements this difficulty is overcome by recording the linear response over two decades in frequency ($f = 0.01$–1 MHz) while the dissipation is maintained below 0.1 nW/mm$^2$. As shown in Ref. [15], the linear regime at low rf excitation level ($\lesssim 1 \mu$T) and small amplitudes of vortex line oscillation ($\lesssim 1$ Å) lends itself to quantitative analysis, in spite of strong surface and/or bulk pinning. From the measurement of the effective ac penetration depth $\lambda_{ac}(f)$ we extract the damping parameter, the flux-flow resistivity $\rho_f$ or the free-flow skin depth $\delta_f = \sqrt{\rho_f/(\pi \mu_0 f)}$. The frequencies are still sufficiently low so that the anomalous skin-depth ($\delta_f \gtrsim 10 \mu$m $\gg l_m \approx 5000$ Å) and relaxation-time effects ($\tau \lesssim 0.2$ ns [4]) can be neglected.

Our measurements have been performed on two large single crystals (labelled B22 and B3b). They were prepared in Grenoble and were spark cut from the same ingot. Their quality is very similar to that of the sample used in the earlier dc measurements of Ref. [4].
After cutting, the samples were annealed, but no surface polishing was applied. Their final dimensions are: $L(x) \times W(y) \times d(z) = 5.5(\bar{a}^*) \times 2.9(\bar{a}) \times 1.16(\bar{a})$ mm$^3$ (B3b) and $5.5(\bar{c}) \times 3.04(\bar{a}^*) \times 0.63(\bar{a})$ mm$^3$ (B22). They have low residual resistivity, $\rho_n = 0.52 + 1.44T^2 + 0.02(\mu_0H)^2$ $\mu\Omega$cm ($\mu_0H$ in Tesla), for currents $J \perp \hat{c}$, which is the case in our measurements. In these and similar samples [15,17], the resistivity for $J \parallel \hat{c}$ is smaller by a factor $0.33 – 0.37$ (for $T \lesssim T_c$). The measured $\rho_n$ corresponds to a mean-free path $l_m \gtrsim 10\xi$, proving that the samples represent the moderately clean limit [18].

The dc field $\mu_0H \leq 3$ T and the vortex lines are along the $\hat{z}$ direction. The excitation field, $he^{-i2\pi ft}$, is applied along the $\hat{x}$ direction so that the vortices oscillate in the $xz$ plane, and currents and electric fields $E$ are induced in the $\hat{y}$ direction. Due to the Lorentz-force anisotropy, $E_y$ is concentrated on the two main faces $z = \pm d/2$ so that the sample mimics the 1-dim geometry of an infinite slab. The measured signal is the flux $\Phi_{ac}$ through a 15-turn pick-up coil wound directly around the sample in the $\hat{x}$ direction. The apparent complex penetration depth is defined as $\lambda_{ac} = \lambda + i\lambda'' = [\Phi_{ac}(H) - \Phi_{ac}(0)]/(2\mu_0hW)$.

The sample with its pick-up coil and the slightly larger excitation solenoid are placed inside the mixing chamber of a $^3$He-$^4$He dilution refrigerator. The measurements are performed in the temperature range $T = 0.1–0.4$ K, by recording the spectrum $\lambda_{ac}(f)$ at fixed field $H$ and by moving from one field value to the next, while the temperature is kept constant with a feed-back loop. The temperature is monitored with a calibrated Ge resistance thermometer inside the mixing chamber. Its field dependence is adjusted by comparing the measured $H_{c2}$ values with those in Ref. [8]. The amplitude and phase of the signal voltage are calibrated against the difference between the responses in the normal and the Meissner states. This normalization procedure yields a phase accuracy better than $1^\circ$ and a resolution $\delta\lambda_{ac} \lesssim 1\mu m$ [18].

Typical spectra of $\lambda_{ac}(f)$ are shown in Fig. [1]. The complete data set consists of more than 200 different spectra. The real and imaginary parts of $\lambda_{ac}$ can be accurately fit with the following formula [13,17]

\[
\frac{1}{\lambda_{ac}} = \frac{1}{L_S} + \sqrt{\frac{1}{\lambda_C^2} - \frac{2i}{\delta_f}}.
\]  

$L_S(H,T)$ and $\lambda_C(H,T)$ are two frequency-independent lengths, which describe surface and bulk pinning respectively. The high-frequency limit, $\lambda_{ac}(f \to \infty) = (1+i)\delta_f/2$, corresponds to the ideal flux-flow response, while the low-frequency limit, the quasistatic response, $\lambda_{ac}(f \to 0) = \lambda'(0) = (\lambda_C^{-1} + L_S^{-1})^{-1}$, is purely inductive and cannot discriminate between surface and bulk pinning. The relative weight of surface and bulk pinning becomes apparent in the crossover regime.

The excitation field penetrates as the sum of two modes. The first mode, localized near the surface, is associated with strong screening currents, whose amplitudes are governed by surface roughness. If the bulk sample is free of defects, the screened field penetrates further over the free-flow depth $\delta_f$. This situation ($\lambda_C \to \infty$) has been systematically the case in the conventional superconductors, which so far have been measured [18]. In contrast, if there are bulk defects, such as those usually introduced in classical theories of pinning [14], the bulk mode is strongly attenuated, but penetrates at low frequencies over a Campbell length $\lambda_C < \delta_f$ at low frequencies).

Curiously enough, our moderately clean UPt$_3$ crystals display a large bulk pinning
strength $1/\lambda_c$, as evident from Fig. 2. The data in Fig. 2 show that in both crystal orientations at low fields pinning is strong and dominated by the surface process. This is not surprising, since after spark cutting the surfaces are visibly rough. With increasing field, surface pinning falls down and finally vanishes at a reversible field value close below the $B \rightarrow C$ transition. This result is consistent with recent observations in point contact spectroscopy \[18\]  \[20\], which suggest a suppression of the order parameter at the surface at higher fields. Bulk pinning decreases more slowly as a function of $(H_{c2} - H)$ and is the dominant source of pinning in the $C$ phase. No clear anomaly can be distinguished at the $B \rightarrow C$ transition \[17\]. Nor do we observe hysteresis as a function of the field-sweep direction. Low-field vortex-creep measurements \[21\] have suggested that new mechanisms could be present, such as intrinsic pinning by domain walls in the bulk \[7\]. If this is the source for the bulk pinning, then the domain walls have to persist in the $C$ phase. Additional work, including variations in surface treatment, should provide better insight in the pinning processes.

The normalized flux-flow resistivity, $\rho_f/\rho_n$, is shown in Fig. 3. As expected for the moderately clean superconductor, $\rho_f(H)$ exceeds the "normal-core limit", $\rho_f = \rho_n H/H_{c2}$, in both crystal directions. When the contributions from bulk and surface pinning have been removed, the $\rho_f(H)$ data become sample independent, as it should be. In the direction $\mathbf{H} \perp \hat{c}$, where comparison is possible, $\rho_f(H)$ agrees quantitatively with the earlier dc measurements of Ref. \[4\]. It is usual to define the initial slope in Fig. 3 as

$$\frac{\rho_f(H)}{\rho_n(H)} = r_0 \frac{H}{H_{c2}}. \tag{2}$$

At low-field the scaling factor $r_0(T)$ obtains, in the perpendicular direction, the value $r_{0,\perp} = 1.6 \pm 0.15$ in the low-temperature limit. In fact, as seen in Fig. 3, below $0.3 \mathrm{K} \simeq 0.6 T_c$, $r_0$ is already temperature independent. This is in agreement with the results in Ref. \[4\], when they are extrapolated to low temperatures, and testifies for good general consistency between two different measuring methods and samples. It should be noted that the temperature independence below $0.3 \mathrm{K}$ also embraces the bulk and surface pinning strengths in Fig. 2.

In effect, Eq. (2) states that, at low field, flux-flow dissipation is additive and therefore proportional to the vortex density, $n = \mu_0 H/\varphi_0$. Assuming that we can identify $\varphi_0/\mu_0 H_{c2}$ with $2\pi b^2$, where $b$ is an effective core radius, as is the case for conventional Abrikosov vortices, then Eq. (2) can be written as

$$\frac{\rho_f(H)}{\rho_n(H)} = r^* 2\pi nb^2. \tag{3}$$

Here $r^* = r_0$ for Abrikosov vortices. More generally, when $H_{c2}$ is not an appropriate scaling parameter due to Pauli paramagnetic limitation of $H_{c2}$ and/or possible changes in the core structure, $r^*(T)$ can still be considered a dimensionless factor which depends on the mechanism of dissipation.

In the dirty limit ($l_m \ll \xi$) one has $r_0 \lesssim 1$ and Eq. (2) reads $\rho_f/\rho_n \approx H/H_{c2}$ (normal-core model). In the clean limit dissipation is governed by the relaxation of the quasiparticle excitations localized in the vortex core. We shall make use of the theory by Kopnin et al. \[22\]. They take $b = \xi$ as the size of the confining vortex-core potential, allow for anisotropy in $\Delta(k_F)$ over a spherical Fermi surface, and predict that
\[ r^*(T) \simeq \alpha \frac{k_B T_c}{\Delta_{\text{max}}(T)}, \quad \alpha = \frac{2 \Delta_{\text{max}}^2}{\langle (1 - (k \cdot \hat{H})^2) \Delta^2(k) \rangle}. \]

The \( \langle \rangle \) brackets denote a weighted average over the Fermi surface. For an isotropic gap (s-wave superconductor) the dimensionless factor \( \alpha \) is unity. If \( \Delta(k) \) has nodal structure, or at least strong anisotropy, then \( \alpha \) is expected to be larger than unity and its value depends on the orientation of the vortex lines with respect to the crystal axes. If we take \( \Delta_{\text{max}}(0) = 1.9 k_B T_c \) (and \( r^* = r_0 \)), Eq. (4) yields \( \alpha_\perp = 1.6 \times 1.9 \approx 3.0 \), which agrees within the combined experimental precisions with the value 3.2 obtained in Ref. [11]. The large \( \alpha \) is as expected because of gap anisotropy.

The analysis of the measurements in terms of Eq. (4) applies for conventional superconductors with \( r_0 = r^* \), or equivalently, \( b^2_\perp = \xi_a \xi_c = \varphi_0/(2\pi \mu_0 H_{c2\perp}) \). Here \( H_{c2\parallel} \) is not a good scaling parameter because of the Pauli paramagnetic limitation at low temperature in this direction. To compare the results in the two orientations, we use Eq. (3) by taking \( b^2_\parallel = \xi_a \xi_c \) and \( b^2_\perp = \xi^2_a \). Furthermore, taking the high-temperature anisotropy \( \xi_a = 0.6 \xi_c \), we find the low-temperature values of \( r^* \) to be \( r^*_\perp = 1.6 \) and \( r^*_\parallel = 4.7 \times 0.6 = 7.8 \). Eq. (4) gives that these correspond to \( \alpha_\perp = 3.0 \) and \( \alpha_\parallel = 7.8 \times 1.9 = 14.8 \), which represents a large anisotropy of \( \alpha_\parallel/\alpha_\perp \approx 5 \).

An anisotropy in \( \alpha \) as large as this cannot be explained by the UPt$_3$ structure, by taking into account the possible anisotropies from the \( D_{6h} \) point group symmetry in the expression of \( \alpha \) in Eq. (4). For instance, let us consider the gap structure in the 2-dim representations \( E_{1g} (\Delta \sim \{k_x k_y, k_x k_y\}) \) and \( E_{2u} (\Delta \sim \{k_z (k_x^2 - k_y^2), 2k_x k_y k_y\}) \), which are appropriate to the \( A \) and \( B \) phases [8]. For \( E_{1g} (\Delta \sim k_z (k_x \pm i k_y), B\text{-phase}) \) we find \( \alpha_\parallel = 1.25 \xi_a \xi_c = 4.4 \). For \( E_{2u} \), which is strongly supported by experiment [4,8], there is no anisotropy at all: \( \alpha_\parallel = \alpha_\perp = 3.9 \). Thus the small jump of about 10% in \( \rho_f/\rho_n \), which was observed in Ref. [4] at the \( A \rightarrow B \) transition for vortices in the \( aa^* \) plane, might be ascribed directly to anisotropy in crystal structure and symmetry breaking, but not the large value of \( r_0/\rho_0 \).

Quasiparticle scattering may additionally contribute to the anisotropy in \( \rho_f(H) \). An estimate can be worked out by comparing the anisotropy in \( \rho_n(J) \) to that in the slope of the critical field \( H_{c2}(T) \) at \( T_c \): \( dH_{c2}/dT \approx -7.2 \ T/K \) and \( dH_{c2}/dT \approx -4.6 \ T/K \). These values yield for the ratio of the coherence lengths \( \gamma = \xi_a/\xi_c \approx 0.64 \), which corresponds to an effective mass ratio of \( m^*_a/m^*_c = \gamma^2 \approx 0.41 \). The anisotropy in \( \rho_n(J) \) amounts to a ratio \( \rho_{n,\parallel}/\rho_{n,\perp} \approx 0.33 \sim 0.37 \). This value includes the anisotropies in effective masses and in quasiparticle relaxation times. Comparing the two estimates we conclude that not much can be allowed for the anisotropy in quasiparticle scattering.

Finally, one more source for anisotropy is a different vortex core structure in the two orientations. The 2-dim Ginzburg-Landau calculations in Ref. [10] suggest the existence of several competing core structures in the axial direction \( Hparallel \hat{c} \) and thus the perpendicular direction is likely to be again different. In Fig. 3 the large \( \rho_f/\rho_n \) measured with \( Hparallel \hat{c} \) could then be explained by assuming \( r^*_\parallel = r^*_\perp \) in Eq. (3) and \( b^2_\parallel = 3 b^2_\perp = 3 \xi_a \xi_c \) rather than \( b^2_\parallel = \xi^2_a \), as was done above. If we ascribe in this way the entire observed anisotropy to an increase in the effective core radius in the axial orientation and take \( b_\parallel = \xi_c \), we estimate the ratio of the fourth order Ginzburg-Landau coefficients to be \( \beta_2/\beta_1 = \xi_a^2/\xi_c^2 \gtrsim 0.2 \). This is in fair agreement with the independent determination \( \beta_2/\beta_1 \sim 0.2 \sim 0.5 \) extracted from the specific-heat jump at \( T_c \) [4,8].
Such an interpretation assumes that the effective core radius in the axial orientation acquires the value $\tilde{\xi} \approx \sqrt{3}\zeta_0\zeta_c \approx 2\zeta_0$. This is possible if the singular “hard vortex core” has reduced rotational symmetry, as is the case in $^3$He-B at low temperatures. A second possibility is that the singular hard core lies embedded within a larger “soft-core” which has an effective radius comparable to $\xi$. In summary, our measurements provide new indication for the presence of unconventional vortex structures in UPt$_3$.

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FIG. 1. Spectra of the apparent penetration depth $\lambda_{ac}(f) = \lambda' + i\lambda''$, showing as a function of frequency the crossover from pinning-dominated to free-flow response (i.e. the high frequency plateau $\lambda'' \simeq \lambda' = \delta f/2$). The solid curves are fits to Eq. (1) with $L_S$, $\lambda_C$, and $\rho_t$ as adjustable parameters. To illustrate their relative influence on the fitting, the dashed curves represent pure bulk-pinning ($1/L_S = 0$) and the dash-dotted curves pure surface-pinning ($1/\lambda_C = 0$).

FIG. 2. Field dependence of the bulk and surface pinning strengths, $1/\lambda_C$ (full symbols) and $1/L_S$ (open symbols) in two crystal orientations. The different symbols denote measurements at different temperatures below 0.33 K. The data are derived from fits to the measured $\lambda_{ac}(f)$ spectra, as shown in Fig. 1. Bulk pinning $1/\lambda_C$ vanishes at $H_{c2}$, while the surface contribution $1/L_S$ decreases more rapidly and vanishes just below the $B \rightarrow C$ transition. The surface-pinning fraction is shown in gray-scale units in the inset, together with the UPt$_3$ phase diagram as a function of field and temperature.

FIG. 3. Flux-flow resistivity $\rho_f(H)$, as deduced by fitting the measured $\lambda_{ac}(f)$ (Fig. 1) to Eq. (1). It is seen that $\rho_f(H)$ becomes temperature independent below 0.3 K as a function of $H/H_{c2}$. The results differ when measured with $H \parallel \hat{c}$ (sample B3b) or $H \perp \hat{c}$ (sample B22). At low fields this anisotropy is quantified by the initial slopes $r_0(T)$ in Eq. (2) (dashed lines), yielding $r_{0\perp} = 1.6 \pm 0.15$ and $r_{0\parallel} = 4.7 \pm 0.3$. The inset shows unscaled data at 0.15 K such that $H_{c2}$ coincides in the two directions. The overshoot in $\rho_f$, just below $H_{c2}$, is an artifact due to the onset of flux penetration at the edges of the sample. In the main frame as well as in Fig. 2 such data points have been omitted.