# $p$-Wave Superconductivity in $UBe_{13}$

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The specific heat in the superconducting state of $UBe_{13}$ shows marked deviations from BCS theory and obeys a $T^3$ rather than an exponential law at low temperatures. A good description is obtained by the assumption of an Anderson-Brinkman-Morel $p$-wave superconducting state at all temperatures. The value of the spin-fluctuation parameter deduced is large and consistent with the stability of such a state.

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Since the discovery that liquid $^3He$ is a $p$-wave superfluid at very low temperatures[1–4] there has been an added impetus to the search for anisotropic superconducting metals, but there have been no clear candidates among conventional superconductors. However, it has been recognized for some time that certain rare-earth and actinide intermetallic compounds starting with $CeAl_3$[5] are electronic Landau Fermi liquids at low temperatures with very strongly enhanced specific heat and susceptibility analogous to the normal state of liquid $^3He$. The recent discovery of superconductivity in such materials has led naturally to the speculation that the analogy to $^3He$ may be closer for these materials than for conventional superconductors. The first system, $CeCu_2Si_2$,[6] has had materials difficulties[7–9] but the new heavy-electron superconductor[10] $UBe_{13}$ has proved free of such problems and also has a higher superconducting transition temperature. In this Letter we wish to demonstrate that the analogy to $^3He$ is very close for $UBe_{13}$ by analyzing the specific heat in the normal and superconducting states. The specific heat in the superconducting state at low temperatures shows a very clear deviation from the BCS form but a good description can be obtained by the assumption that the superconducting state is an Anderson-Brinkman-Morel (ABM) $p$-wave superconductor at all temperatures. Further, the value of the spin-fluctuation parameter deduced is large, larger than that attained in liquid $^3He$ at the solidification curve and consistent with an ABM state at all temperatures in the Brinkman-Serene-Anderson (BSA) theory.[4,11] We propose from this comparision between theory and experiment that $UBe_{13}$ is the first $p$-wave electronic superconductor and also the first case of an electronic superconducting transition driven by an interaction other than the electron-phonon interaction.

To demonstrate the analogy, mentioned above, between the normal-state properties of the electronic subsystem of $UBe_{13}$ and liquid $^3He$, we plot the experimental specific heats $C_p$ of both $UBe_{13}$ and $^3He$ at low temperatures in the form of $C_p/T$ vs $T$ in Fig. 1. For $UBe_{13}$, our measurements were made on a small ($\sim 200$ mg) polycrystal which showed a rather sharp superconducting transition. Although almost negligible in this temperature range, the lattice contribution to $C_p$ was subtracted for $UBe_{13}$ and

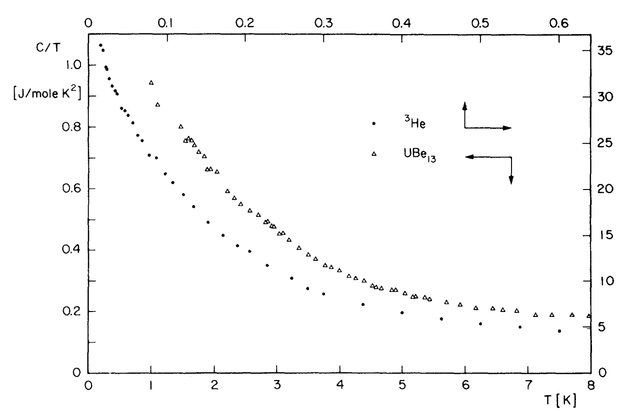


FIG. 1. Low-temperature specific heats $C_p/T$ of liquid $^3He$ and of $UBe_{13}$. Note the different scales for the two materials. The data for $^3He$ are from Refs. 12 and 13.

 

is not contained in the points shown in Fig. 1. The data for $^3$He were taken from Brewer, Daunt, and Sreedhar[12] and Anderson, Reese, and Wheatley.[13] It should be noted that the scales for both $C_p/T$ and $T$ are different for the two materials.

The effective Fermi temperatures $T_F^*$ for $UBe_{13}$ and $^3$He are $\approx 8$ K and a few tenths of a kelvin, respectively, as estimated from the $C_p(T)/T$ data shown in Fig. 1. While $^3$He is superfluid only below 2.7 mK ($\approx 10^{-2} T_F^*$), $UBe_{13}$ is superconducting below 0.9 K ($\approx 10^{-1} T_F^*$). The specific heat above and below $T_c$ is shown in Fig. 2. The occurrence of superconductivity with a large anomaly in $C_p(T)$ is the final proof that the large values of $C_p(T)/T$ are of electronic origin.

The ratio of $C_p$ in the superconducting $[C_s(T)]$ and normal $[C_n(T)]$ states is shown in Fig. 3 in the form $\log[C_s(T)/C_n(T_c)]$ vs $T_c/T$. A comparison with the universal BCS curve shows considerable deviations both at $T \approx T_c$ and as $T \rightarrow 0$. The extrapolated discontinuity $\Delta C/C_n(T_c)$ is $\approx 2.5$ (BCS value, 1.43). As $T \rightarrow 0$, $C_s/C_n(T_c)$ is not exponential but has a power law form: $C_s/C_n(T_c) \approx 2.8(T/T_c)^3$.

The fact that $C_s(T) \propto T^3$ as $T \rightarrow 0$ rather than exponential suggests that the gap function goes to zero someplace on the Fermi surface and that we have a *p*-wave (or possibly higher quantum

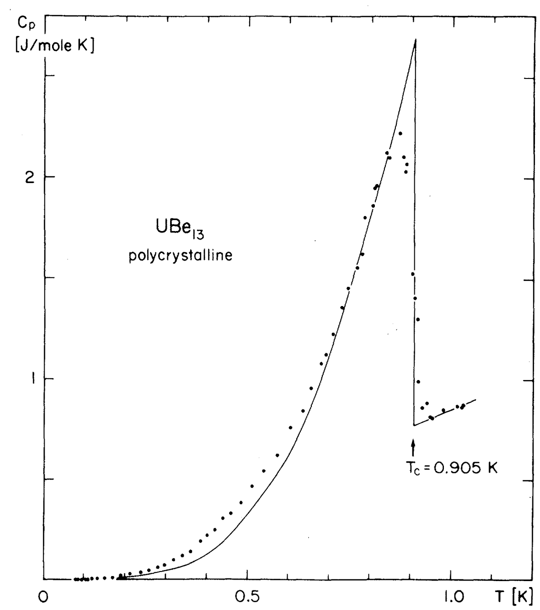


FIG. 2. Low-temperature specific heat of $UBe_{13}$ between 0.07 and 1.1 K. The solid line is calculated for an ABM *p*-wave superconductor with strong-coupling corrections (see also Fig. 3).

number) superconducting state. Note that alternative origins of nonexponential behavior in $C_s(T)$, such as gapless *s*-wave pairing due to magnetic impurities or multiphase samples with normal regions, lower the value of $\Delta C/C_n(T_c)$ and lead to $C_s(T) \sim T$ as $T \rightarrow 0$ in contradiction to experiment. *p*-wave superconductivity has been extensively studied in $^3$He. Two *p*-wave superconducting states have been observed there.[2-4] One is the Balian-Werthamer (BW) state in which the magnitude of the gap is constant over the Fermi surface. It has thermodynamic properties identical to an *s*-wave superconductor. The second is the ABM state observed near solidification with a gap function[2,4]

$$\Delta_{\sigma\sigma'}(\vec{k}) = \frac{3}{2}^{1/2} \Delta(T) \begin{pmatrix} \hat{k}_x + i\hat{k}_y & 0 \\ 0 & \hat{k}_x + i\hat{k}_y \end{pmatrix}. \quad (1)$$

The magnitude of the gap, $(\Delta^+\Delta)^{1/2}$, now vanishes at the points $k_x = k_y = 0$, $k_z = \pm k_F$. In the weak-coupling limit it is straightforward to evaluate the specific heat of the ABM state and the result is shown also in Fig. 3 (identified as w.c.). The jump in specific heat at $T_c$ is smaller than BCS by a factor of $\frac{5}{6}$. At low temperatures the zeros of the gap give rise to a power law leading to the following results in weak coupling[2,4]:

$$[\Delta C/C_n(T_c)]_{T=T_c} = 1.18;$$
$$\frac{C_s(T)}{C_n(T_c)} = 9.21 \frac{T}{T_c} \left( \frac{T}{\Delta(0)} \right)^2, \quad (2)$$

with $\Delta(0) = 1.65 k_B T_c$. A comparison with the experimental values of $C_p(T)$ for $UBe_{13}$ shows that while the overall form of the curve is in better agreement than the BCS form, shown also in Fig. 3, there are substantial quantitative disagreements. The experimental value of $(\Delta C/C_n)_{T_c}$ is much larger, not smaller, than BCS while the measured values of $C_p(T)$ as $T \rightarrow 0$ are smaller than Eq. (2). Both discrepancies suggest that substantial strong-coupling corrections to the theory are necessary.

We have examined the leading strong-coupling corrections to the ABM state. This problem was investigated and applied to $^3$He by Brinkman and coworkers.[4,11] In the BSA theory the fermions are coupled to a low-lying set of spin fluctuations or paramagnons and the ABM state is stabilized relative to the BW state for sufficiently strong coupling. The jump at $T_c$ in the specific heat is increased while the magnitude of the zero-temperature gap is also increased. Their results, with strong-coupling corrections parametrized by a single parameter $\delta$, are

$$\left.\frac{\Delta C}{C_n}\right|_{T=T_c} = \frac{2.38}{2-1.05\delta};$$
$$\frac{\Delta(0)}{k_B T_c} = 1.65(1+0.13\delta) \tag{3}$$

and $\delta$ is given by

$$\delta = \frac{150\pi^2}{7\zeta(3)}\left(\frac{T_c}{E_F}\right)\int_0^1 dx\left(\frac{\bar{I}}{1-\bar{I}+\alpha\bar{I}x^2}\right)^2, \tag{4}$$

where $1-\bar{I}$ is the Stoner enhancement factor and $\alpha$ is a parameter determining the width of the enhancement region in $\vec{k}$ space. Thus with a single parameter we can determine the specific heat both near $T_c$ and at low temperatures through Eqs. (2) and (3). If we wish to calculate the complete curve of $C_s(T)$ we will need to specify the form of the spin susceptibility $\chi(\vec{q},\omega)$ more precisely and so we have chosen to adopt a simpler scheme with only a single free parameter, $\delta$. We use Eqs. (2) and (3) to determine the $C_s(T)$ for $T\approx T_c$ and $T\to 0$ and interpolate between the two regions with a form that gives an entropy $C_n(T_c)$ at $T_c$:

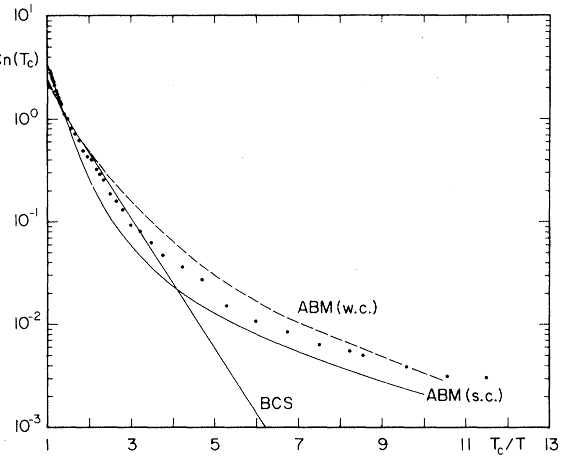


FIG. 3. $C_s/C_n(T_c)$ for superconducting $UBe_{13}$. Dashed line: weak-coupling ABM state; solid lines: BCS and strong-coupling ABM state from Eq. (5).

$$\frac{C_s(T)}{C_n(T_c)} = \frac{3.35}{1+0.27\delta}\left(\frac{T}{T_c}\right)^3 \exp\left[\frac{-\epsilon T}{T_c-T}\right] + \left(1+\frac{2.38}{2-1.05\delta}\right)\left(\frac{T}{T_c}\right)\exp\left[\frac{-\epsilon(T_c-T)}{T}\right], \tag{5}$$

where the interpolation parameter $\epsilon$ is determined by the entropy condition to be $\epsilon = 2.05$. For the parameter $\delta$, we obtain 1.0 from the value of the specific-heat jump at $T_c$. The values obtained from Eq. (5) are shown as curve ABM (s.c.) in Fig. 3.

A consistent application of BSA theory requires that $\delta > 0.87$ in order to have the ABM state stable down to $T=0$. Our value of $\delta$ exceeds this critical value. A comparison to the experiment shows a considerable deviation at low temperatures but the overall shapes of the experimental and theoretical curves are similar. The remaining discrepancy shows the inadequacy of the BSA strong-coupling theory. Signs of such strong-coupling effects are (a) the 10% enhancement of the observed entropy over the value linearly extrapolated from $C_n(T_c)/T_c$, which is most likely due to a further increase in $C_p(T)/T$ in a hypothetical normal state, e.g., as observed in $U_{1-x}Th_xBe_{13}$[14] and $CeCu_2Si_2$[15]; (b) the relatively large value ($\sim 0.1$) of $T_c/T_F^*$; and (c) the large intrinsic resistivity at $T \gtrsim T_c$, which is compatible with $p$-wave pairing only because strong-coupling effects reduce the coherence length to a value of a few lattice spacings, comparable to the mean free path.

The close analogy in the physical properties between the heavy Fermi liquids $UBe_{13}$ and $^3He$ suggests that there should be an underlying similarity in the microscopic descriptions. Anderson and Brinkman[4] and also Vollhardt[16] have successfully interpreted the low-temperature properties of $^3He$ in terms of an almost localized Fermi liquid using the Brinkman-Rice theory[17] for the Hubbard model near the Mott transition. If in $UBe_{13}$ we assume that the heavy-electron liquid arises because of coherence between the $5f$ states and that the only role of the Be bands is to determine the size of the U-U hopping-matrix element and the Fermi level is the $5f$ band, then it is clear that the Brinkman-Rice theory of a strongly correlated Fermi liquid can be directly applied to the U $5f$ states. The important requirement is that the number of $5f$ electrons be very close to an integral number per U and that the on-site Coulomb interactions be strong. These conditions are quite reasonable for $UBe_{13}$. In such a model the very large specific heat arises from the low-lying excitations which arise from the rearrangement of the moments on the U sites with integral occupation, and it is the coupling of the quasiparticles through these excitations which leads to the $p$-wave superconducting state.

In summary, an examination of the low-temperature specific heat of $UBe_{13}$ leads us to identify it as the first metal with an ABM $p$-wave superconducting state, driven by interactions other than the electron-phonon interaction. There are important strong-coupling corrections and the strong-coupling parameter is larger than that attained in $^3He$ and is consistent with the ABM state being stable at all temperatures. We suggest that there is a close analogy between the heavy-electron liquid in $UBe_{13}$ and the nearly localized Fermi liquid in $^3He$.

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