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# Kondo Coherence in $UBe_{13}$: Magnetoresistance at High Pressure

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We have used resistance measurements at pressures to 95 kbar and magnetic fields to 8 T to probe the establishment of Kondo coherence in the heavy-fermion Kondo lattice $UBe_{13}$. With increasing pressure, the negative magnetoresistance characteristic of noninteracting Kondo impurities evolves into the positive magnetoresistance expected for a coherent Kondo lattice. We show that coherence is established by the suppression of inelastic scattering, over a temperature range which expands as pressure decreases the superconducting transition temperature $T_c$ and increases the Kondo temperature $T_K$.

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Heavy-fermion materials display many of the characteristics of Kondo lattices, whose properties are a problem of continued experimental and theoretical interest. At high temperature, a Kondo lattice behaves like an ensemble of noninteracting Kondo impurities. In this regime, the resistivity increases and the magnitude of the local moments decreases as the temperature is lowered, with a characteristic temperature scale of the Kondo temperature $T_K$. At the lowest temperatures, the Kondo lattice can be described as a periodic Fermi liquid.[1] The crossover from incoherent Kondo impurity scattering at high temperatures ($T \gg T_K$) to the formation of a Fermi-liquid regime at low temperatures ($T \ll T_K$) is commonly referred to as the establishment of Kondo coherence. In addition to the appearance of bandlike character in the transport properties, sign changes in the magnetoresistance[2] and Hall coefficient,[3] as well as peaks in the temperature dependence of the thermoelectric power[4] and specific heat,[5] are characteristic of the Fermi liquid and suggest the presence of low-energy features and gaps at the Fermi surface. Despite extensive theoretical work on the high-temperature state,[6] and more recently on the description of the low-temperature state as a Fermi liquid,[1] a complete theoretical understanding of the Fermi-surface features reflected in these transport measurements remains lacking. We will demonstrate that, in contrast, comparison of elastic and inelastic scattering rates determined from magnetoresistance measurements provides a clear microscopic criterion for the onset of coherence. Our interest has focused on $UBe_{13}$ primarily because of all the known heavy-fermion systems it shows the least evidence for the existence of a Fermi-liquid regime above the superconducting transition temperature at 1 bar. We have studied the crossover in $UBe_{13}$ between incoherent and coherent scattering in a controlled way by variation of pressure.

Polycrystalline samples of $UBe_{13}$ were prepared by arc melting high-purity uranium and beryllium in an argon atmosphere.[7] The sample and a lead manometer were placed in series in the Bridgman anvil cell.[8] Platinum leads are attached to both with silver filled epoxy, allowing simultaneous four-probe measurements of the lead and sample resistances. The average pressure in the completed cell was determined from the depression of the lead superconducting transition temperature.[9] Our modest uncertainty in the pressure determination results from the broadening of the lead transition by pressure inhomogeneities, reflecting the quasihydrostatic nature of the Bridgman clamp. Nonmagnetic tungsten carbide anvils are used to minimize the effects of residual fields, which are estimated to be several hundred oersteds.

The electrical resistivity $\rho$ of $UBe_{13}$ is plotted as a function of temperature in Fig. 1(a) for pressures ranging from 1 bar to 95 kbar. The application of pressure increases the temperature $T_{max}$ at which the resistance is a maximum from 2.2 K at 1 bar to $\sim$50 K at 95 kbar. Most importantly, as depicted in Fig. 1(b), pressure expands the temperature range $T < T^*$ for which $\rho(T) = \rho_0 + AT^2$. Although it is doubtful that the parameters derived from such a fit are accurate at low pressures, a substantial temperature range emerges and expands with pressures above $\sim$40–50 kbar in which $\rho$ is unquestionably quadratic in temperature. Because the presence of a large quadratic term in the temperature dependence of the resistivity accompanies coherence,[10] these data indicate that pressure establishes coherence in $UBe_{13}$.

We have used magnetoresistance experiments to probe the microscopic mechanisms responsible for the establishment of Kondo coherence with pressure. The magnetic field dependence of the normalized resistivity of $UBe_{13}$ at 1.2 K is plotted in Fig. 2(a), for pressures ranging from 35 to 95 kbar. As previously observed,[11] the magnetoresistance at the lowest pressures is negative, with the field dependence expected for isolated spin-$\frac{1}{2}$ Kondo impurities.[12] The magnetoresistance decreases with pressure and changes sign between 53 and 95 kbar.

 

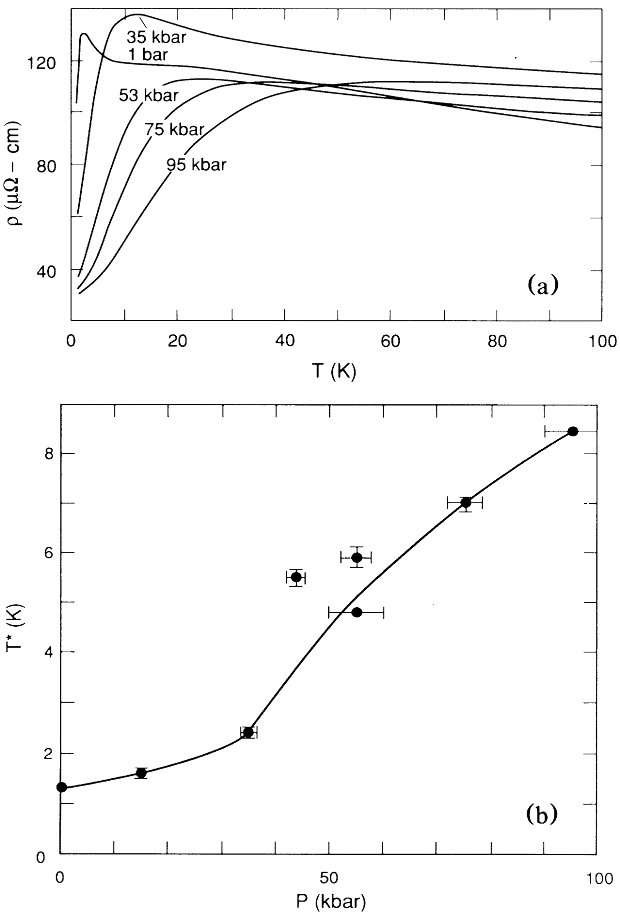


FIG. 1. (a) Electrical resistivity $\rho$ of $UBe_{13}$ as a function of temperature for pressures ranging from 1 bar to 95 kbar. (b) Pressure dependence of temperature range $T < T^*$ for which $\rho(T) = \rho_0 + AT^2$. Solid lines are guides for the eye.

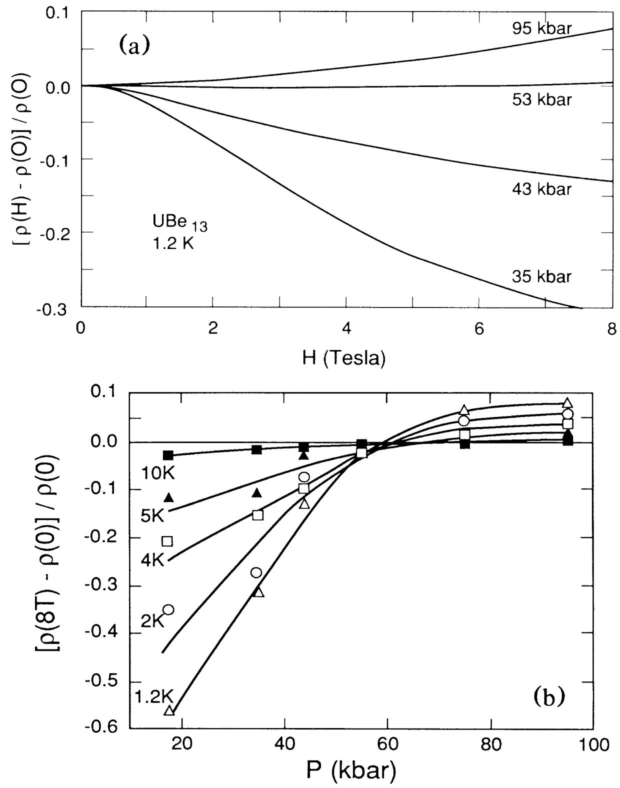


FIG. 2. (a) Normalized magnetoresistance of $UBe_{13}$ at 1.2 K, with pressures ranging from 35 to 95 kbar. (b) Pressure dependence of the normalized magnetoresistance at 8 T, with temperatures ranging from 1.2 to 10 K.

At the highest pressures, the magnetoresistance is positive and quadratic in field. Restoration of this typical metallic behavior at high pressures demonstrates the presence of the bandlike properties associated with Kondo coherence.[2] Overall, these data show that there are two pressure-dependent contributions to the magnetoresistance. As in the spin-$\frac{1}{2}$ Kondo impurity model, the negative contribution which dominates at low pressure results from the field dependence of inelastic scattering at local moment sites. However, this contribution gives way to the positive bandlike response expected for the coherent Kondo lattice as the pressure is increased.

It is instructive to examine the role temperature plays in the establishment of Kondo coherence. To do so, we have plotted the normalized magnetoresistance in a fixed field of 8 T as a function of pressure in Fig. 2(b). Between 1.2 and 10 K, the magnetoresistance is initially negative but decreases in magnitude before changing sign and saturating at a positive value as the pressure is increased. The decrease in the negative contribution can be understood as the result of a Kondo temperature which increases with pressure, as expected for uranium compounds.[13] Namely, since the Kondo temperature $T_K$ and Kondo field $H_K$ scale, a smaller change in the resistance is observed on application of a given field as pressure increases the overall field scale. In contrast, the positive contribution to the magnetoresistance appears to be essentially pressure independent. Since the magnetoresistance of a normal metal is proportional to its conductivity, this result is consistent with the insensitivity of the residual resistivity $\rho_0$ to high pressures.[14] As can be seen in Fig. 2(b), the magnetoresistance experiment identifies a set of temperatures and pressures for which the two contributions to the magnetoresistance are equal, and the magnetoresistance is zero. This pressure-dependent temperature $T_{coh}$ may be interpreted as a coherence temperature, as it separates the low-pressure regime of Kondo impurity behavior from the high-pressure regime having the coherent properties of a normal metal. We have plotted the pressure dependence of $T_{coh}$ in Fig. 3.

The establishment of coherence has a clear signature

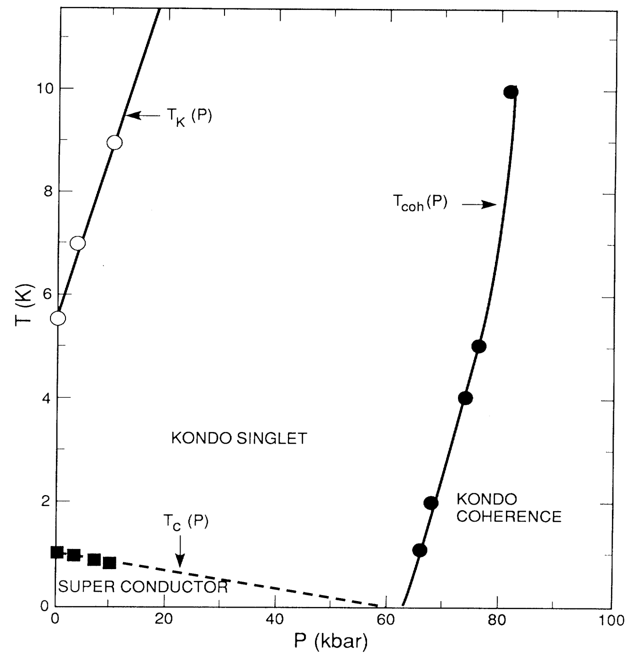


FIG. 3. Pressure dependence of the Kondo temperature $T_K \propto 1/\gamma_0$ (Ref. 15), superconducting transition temperature (Ref. 16), and coherence temperature $T_{\rm coh}$, taken from the magnetoresistance experiment reported in the present work.

in the temperature dependence of the electrical resistivity of a Kondo lattice system. At the highest temperatures, a Kondo lattice may be considered an ensemble of isolated Kondo impurities. The high-temperature resistivity of a Kondo lattice is dominated by the inelastic scattering of conduction electrons from partially compensated local moments. This process has an energy scale of the Kondo temperature $T_K$ and makes a contribution to the resistivity $\rho$ with the temperature dependence $\rho(T) \sim a - b\ln(T)$. A negative magnetoresistance is expected[12] as magnetic field suppresses this spin-flip scattering process. As the temperature drops below $T_K$, coherent electron-electron scattering by Baber processes becomes increasingly important,[10,17] eventually leading to a resistivity which is quadratic in temperature $\rho(T) \propto AT^2$ and in field[2] $\rho(T) \propto H^2$. At intermediate temperatures, neither the incoherent Kondo impurity nor the coherent electron-electron contributions to the scattering dominate and there is no convenient theoretical description of the Kondo lattice. However, observation of a coherent Fermi liquid requires the experimental temperature range to be much lower than the Kondo temperature, so that inelastic Kondo impurity scattering processes are effectively frozen out. In addition, electronic interactions must be sufficiently weak that superconductivity and magnetic order do not set in at higher temperatures than the Kondo coherence. Since the limited temperature range for which $\rho(T) \sim AT^2$ severely compromises the validity of such a fit, and no other signs of Kondo coherence are observed, we conclude that the fully coherent state is not achieved in $UBe_{13}$ at 1 bar before the onset of superconductivity at 0.95 K. This situation is qualitatively changed by the application of high pressures.

The first effect of pressure on $UBe_{13}$ is the variation of the single-impurity Kondo temperature $T_K$. The Kondo temperature for uranium compounds increases with pressure, presumably through increased hybridization of the $f$ level with the conduction band. The pressure dependence of $T_K$ is extracted from that of the electronic specific-heat coefficient $\gamma_0$ using the relation $\gamma_0 = 0.68R/T_K$,[15] where $R$ is the gas constant. We note that this expression is strictly valid only at the low pressures for which $UBe_{13}$ displays Kondo impurity behavior and $T_K$ is well defined. As shown in Fig. 3, below 8 kbar pressure increases the energy scale $T_K$ indicating that the inelastic scattering process is frozen out at progressively higher temperatures, revealing the intrinsic elastic scattering resulting from the underlying periodicity of the lattice.

The second effect of pressure is the suppression of superconductivity. The pressure dependence of the superconducting transition temperature $T_c$[16] is plotted in Fig. 3. Assuming that $T_c$ continues to decrease linearly with pressure, extrapolation indicates that ~60 kbar must be applied to suppress the superconductivity so that $T_{\rm coh} > T_c$. Thus, the net effect of the pressure dependences of $T_K$ and $T_c$ is that an expanding range of temperatures $T$ satisfying $T_c < T < T_{\rm coh} < T_K$ emerges in which elastic scattering dominates the resistivity and coherent behavior may be observed. The onset of a positive magnetoresistance and the peak in the thermoelectric power[18] observed above 60 kbar are strong evidence for the establishment of Kondo coherence under pressure in $UBe_{13}$. Pressure studies of $UBe_{13}$ are thus particularly interesting, as they provide a means of tuning the onset of Kondo coherence into the experimental temperature range without the introduction of disorder which accompanies doping.

The major result of this work is the demonstration that Kondo coherence is established in $UBe_{13}$ under pressure by the suppression of inelastic scattering. The conditions under which the elastic scattering rate characteristic of a coherent Kondo lattice first becomes equal to the inelastic scattering rate characteristic of incoherent Kondo impurities are used to define a coherence temperature $T_{\rm coh}$ for a Kondo lattice. As depicted in Fig. 3, the emergence and subsequent increase of $T_{\rm coh}$ with pressure results from both the depression of the superconducting transition temperature and the increase in the single-impurity Kondo temperature. Details of electronic interactions prevent Fig. 3 from being a universal phase diagram for the Kondo lattice in that the relative magnitudes of $T_K$ and $T_c$ dictate whether superconductivity develops from a coherent or incoherent state. This is an important theoretical consideration for microscopic descriptions of the superconducting state. However, despite the fact that coherence is absent in $UBe_{13}$ at low pressures but not in, for instance, $UPt_3$ at 1 bar, it is interesting to note that the two systems share common features. As an example, it is found that the $T^2$ resistivity coefficient $A$ and the electronic specific-heat coefficient $\gamma_0$ are related by

$$A/\gamma_0^2 = 1.0\times10^{-5}\,\mu\Omega\,\mathrm{cm}(\mathrm{mole\,K/mJ})^2 \qquad (1)$$

in a large number of heavy-fermion systems,[19] including $UBe_{13}$ under pressure.[20] This result suggests that despite the nature of quasiparticle interactions stabilizing the various ground states (superconducting, magnetic, or normal), the many-body enhancements equally affect both $A$ and $\gamma_0^2$. It remains both a theoretical and an experimental challenge to explain this result, given the apparently distinct natures of the low-temperature states in the different materials.

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[1]P. A. Lee, T. M. Rice, J. W. Serene, L. J. Sham, and J. W. Wilkins, Comments Condens. Matter Phys. **12**, 99 (1986).

[2]A. S. Edelstein, R. E. Majewski, and T. H. Blewitt, in *Valence Instabilities and Related Narrow Band Phenomena,* edited by R. D. Parks (Plenum, New York, 1977), p. 115.

[3]T. Penney, J. Stankiewicz, S. von Molnar, Z. Fisk, J. L. Smith, and H. R. Ott, J. Magn. Magn. Mater. **54-57**, 370 (1986).

[4]W. Franz, A. Griessel, F. Steglich, and D. Wohlleben, Z. Phys. B **31**, 7 (1978).

[5]C. D. Bredl, S. Horn, F. Steglich, B. Luthi, and R. M. Martin, Phys. Rev. Lett. **52**, 1982 (1984).

[6]See, for instance, T-K. Lee and S. Chakravarty, Phys. Rev. B **22**, 3609 (1980).

[7]J. L. Smith, Z. Fisk, J. O. Willis, B. Batlogg, and H. R. Ott, J. Magn. Magn. Mater. **55**, 1996 (1984).

[8]J. Wittig, Z. Phys. **195**, 228 (1966).

[9]T. F. Smith, C. W. Chu, and M. B. Maple, Cryogenics **9**, 53 (1969).

[10]Assa Auerbach and K. Levin, Phys. Rev. Lett. **57**, 877 (1986); D. W. Hess, J. Low Temp. Phys. **68**, 311 (1987).

[11]B. Batlogg, D. J. Bishop, E. Bucher, B. Golding, A. P. Ramirez, Z. Fisk, J. L. Smith, and H.R. Ott, J. Magn. Magn. Mater. **63-64**, 441 (1987); J. O. Willis, M. W. McElfresh, J. D. Thompson, J. L. Smith, and Z. Fisk, J. Appl. Phys. **64**, 5613 (1988); G. Remenyi, D. Jaccard, J. Floquet, A. Briggs, Z. Fisk, J. L. Smith, and H. R. Ott, J. Phys. (Paris) **47**, 367 (1986); U. Rauchschwalbe, F. Steglich, and H. Rietschel, Europhys. Lett. **1**, 71 (1986).

[12]V. Andrei, Phys. Lett. **87A**, 299 (1982); P. Schlottmann, Z. Phys. B **51**, 223 (1983).

[13]J. D. Thompson, J. Magn. Magn. Mater. **63-64**, 358 (1987).

[14]M. C. Aronson *et al.* (unpublished); M. W. McElfresh, Ph.D. thesis, University of California, San Diego, 1988 (unpublished).

[15]N. E. Phillips. R. A. Fisher, S. E. Lacy, C. Marcenat, J. A. Olsen, J. Floquet, A. Amato, D. Jaccard, Z. Fisk, A. L. Giorgi, J. L. Smith, and G. R. Stewart, in *Theoretical and Experimental Aspects of Valence Fluctuations and Heavy Fermions,* edited by C. C. Gupta and S. K. Malik (Plenum, New York, 1987), p. 141.

[16]J. W. Chen, S. E. Lambert, M. B. Maple, Z. Fisk, J. L. Smith, and H. R. Ott, in *Proceedings of the Seventeenth International Conference on Low Temperature Physics,* edited by U. Eckern (North-Holland, Amsterdam, 1984), p. 325.

[17]K. Andres, J. E. Graebner, and H. R. Ott, Phys. Rev. Lett. **35**, 1779 (1975).

[18]S. Y. Mao, D. Jaccard, J. Sierro, Z. Fisk, and J. L. Smith, J. Magn. Magn. Mater. **76-77**, 241 (1988).

[19]K. Kadowaki and S. B. Woods, Solid State Commun. **58**, 507 (1986); K. Miyake, T. Matsuura, and C. M. Varma (unpublished).

[20]J. D. Thompson, Z. Fisk, and G. G. Lonzarich, Physica (Amsterdam) B (to be published).