PHENOMENOLOGICAL MODEL OF SUPERCONDUCTIVITY IN $U_{1-x}Th_xBe_{13}$

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A phenomenological model taking into account the interaction between superconductivity and the coherence of Kondo screening is introduced. This model describes the main experimental data on UBe_{13}, including the behaviour of T_c in $U_{1-x}Th_xBe_{13}$ under ambient and elevated pressures.

1. The available experimental data on $CeCu_2Si_2$ and UBe_{13} heavy fermion superconductors (HSF) indicate that both compounds can be considered as superconducting Kondo lattices (KL) with an associated Kondo temperature $T_K < \Delta_{CF}$ and a doublet as a ground crystal field split state [1–6]. In this case the narrow (~ 1 K) Abrikosov–Suhl resonance (ASR) of giant amplitude is formed at the Fermi energy E_F. Heavy quasiparticles with $E \approx E_F$ ("heavy fermions") correspond to the ASR in KL. There are two characteristic temperature regimes in nonmagnetic KLs [3]: at temperatures $T \gtrsim T_K$, KLs can be considered as a collection of many independent Kondo scatterers, whereas at low temperatures $T < T_{c2} < T_K$ a coherence of Kondo screening at different lattice sites sets in.

The transition to the coherent state in KL can be described as an increase (abrupt or continuous) of the size of the spacial region r_{coh} where Kondo screening is correlated. This collective behaviour can be thought of as a static or a dynamic effect. In the first case, coherence corresponds to a static magnetic ordering of partially screened Kondo scattering centers. A dynamic process is also conceivable, where due to fluctuations, the array of collectively ordered spins is strongly time dependent. We should mention here, that besides the exchange interaction, coherence can be induced by the Coulomb interaction.

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The substitution of a Kondo scattering center by a nonmagnetic impurity results in the formation of an induced magnetic moment at this particular site for $T < T_{c2}$ [7,8]. Heuristically speaking, at this location there is a compensating electron cloud but nothing to compensate. For low concentrations of nonmagnetic impurities, the induced magnetic moment therefore appears to be higher than a Kondo screened localized magnetic moment. This effect looks similar to that of nonmagnetic impurities in antiferromagnetic superconductors [9]. In the case of dynamical coherence of Kondo scatterers in KL, these induced magnetic moments are strongly fluctuating.

The observation of a giant specific heat anomaly at $T = T_c$ [1,2] shows that heavy quasiparticles are responsible for the superconducting transition. The associated Fermi energy $E_F \sim T_K/2 \sim 3$–5 K is extremely low. Therefore the Ginzburg number [10,11]

$$\mathrm{Gi} \approx 10^2 (T_c/E_F)^4 \qquad (1)$$

in UBe_{13} is about 0.1–0.7, being compared to the usual value 10^{-14} in ordinary superconductors. Taking into account the smallness of the mean free path in UBe_{13} [12,13], estimates in the dirty limit yield even higher values for Gi. This implies the importance of fluctuations for a description of superconductivity in HFS.

In known HFS there are different cases as regards the relative value of the two characteristic temperatures T_c and T_{coh} [14]. For example, in $CeCu_2Si_2$,

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$T_c \sim 0.5$ K $< T_{c2} \sim 1$–2 K, whereas in UBe_{13}, $T_c = 0.86$ K $\sim T_{c2}$. The proximity of T_c and T_{c2} in the HFS UBe_{13} as well as the large Gi value, enforces the following approach: coherence and superconductivity in UBe_{13} substantially affect each other. As has been put forward by one of the authors [15], both phenomena ought to be described on equal grounds, using the fluctuation theory of phase transitions: due to strong fluctuations, one cannot use the standard Landau formalism of phase transition lines crossing. Here we present the results obtained for UBe_{13} within the framework of a model taking into account strong fluctuations.

2. To describe the interaction between superconductivity and coherence we have used the following Landau–Ginzburg model hamiltonian:

$$H = \int \mathrm{d}\boldsymbol{x}\, [a_1(T - T_{c,0})\varphi_1^2 + b_1\varphi_1^4 + a_2(T - T_{c2,0})\varphi_2^2 + b_2\varphi_2^4 + \lambda_{12}\varphi_1^2\varphi_2^2] \,. \qquad (2)$$

(This ansatz for the hamiltonian cannot be derived strictly. In particular, gradient terms were omitted and a specific interaction term was assumed.) Here $\varphi_1 = \Delta/\Delta_0$ and $\varphi_2 = r_{coh}/r_{coh,0}$ stand for the superconducting and coherence order parameter respectively, $T_{c,0}$ and $T_{c2,0}$ are associated with the bare superconducting and coherence temperatures without the interaction term proportional to the parameter λ_{12}; a_1, b_1, a_2, b_2 are temperature independent coefficients. In this model we assumed that the interaction parameter λ_{12} depends on the Th concentration x:

$$\lambda_{12} = -\lambda_0 + Cx \,. \qquad (3)$$

This choice yields the enhancement of superconductivity due to the $-\lambda_0 < 0$ factor (pure superconducting KL UBe_{13}), and its suppression due to induced magnetic moments on Th sites (see positive Cx term in eq. (3)).

Minimizing eq. (2) with respect to φ_1^2 and φ_2^2, one obtains the critical temperatures for superconductivity and coherence:

$$T_c = T_{c,0} - (\lambda_{12}/a_1)\langle\varphi_2^2\rangle = T_{c,0} - (\lambda_{12}/a_1)\,|T_c/(T_c - T_{c2})|^{2\nu} \,, \qquad (4a)$$

$$T_{c2} = T_{c2,0} - (\lambda_{12}/a_2)\langle\varphi_1^2\rangle = T_{c2,0} - (\lambda_{12}/a_2)\,|T_{c2}/(T_c - T_{c2})|^{2\nu} \,. \qquad (4b)$$

In eq. (4) we have insert the values $\langle\varphi_1^2\rangle \sim |T_{c2}/(T_c - T_{c2})|^{2\nu}$ and $\langle\varphi_2^2\rangle \sim |T_c/(T_c - T_{c2})|^{2\nu}$ from the fluctuation theory of phase transitions [11].

When $T_{c,0}$ and T_{c2} in (4) are small with respect to the interaction term, the solution gives two close transition temperatures $T_c(x)$ and $T_{c2}(x)$ decreasing with the thorium concentration x. This means that superconducting and coherent transitions occur close to each other when $\lambda_{12} < 0$ and the proximity of T_c and T_{c2} decreases the free energy.

On the other hand, when $T_{c,0}$ and $T_{c2,0}$ are much higher than the interaction term, the two close temperatures $T_c(x)$ and $T_{c2}(x)$ split into $T_c \sim T_{c,0}$ and $T_{c2} \sim T_{c2,0}$. When λ_{12} changes its sign at x_i, both transitions "repel" each other. This results in two quite different transition temperatures $T_c(x_i)$ and $T_{c2}(x_i)$.

Eqs. (4) can be solved analytically if we replace $2\nu \sim 4/3$ by $2\nu \sim 1$:

$$T_c = [-\lambda_{12}/a_2 - (a_1/a_2)T_{c,0} + (T_{c,0} + T_{c2,0})/2 + \mathscr{D}^{1/2}](1 - a_1/a_2)^{-1} \,, \qquad (5)$$

$$\mathscr{D} = [(T_{c,0} - T_{c2,0})/2]^2 + (\lambda_{12}/a_2)T_{c,0} - (\lambda_{12}/a_1)T_{c2,0} + \lambda_{12}^2/a_1a_2 \,, \qquad (6)$$

$$T_{c2} = T_c(x)\left(1 + \frac{\lambda_{12}/a_1}{x - T_{c,0} - \lambda_{12}/a_1}\right) . \qquad (7)$$

When $\mathscr{D}$ is decreased from $\mathscr{D} > 0$ to $\mathscr{D} < 0$, the $T_c(x)$ and $T_{c2}(x)$ dependences change qualitatively, as illustrated in fig. 1. For $\mathscr{D} \geqslant 0$, $T_c(x)$ is the sum of a linear plus a $\mathscr{D}^{1/2}$ dependence on x. If $\mathscr{D}$ is negative, there are no solutions (see shaded areas in figs. 1c, 1d). Different concentration dependences of T_c and T_{c2} in fig. 1 are obtained by changing only one parameter a_1, with λ_{12} given by (3). Qualitatively this corresponds to the $T_c(x)$ behaviour in $U_{1-x}Th_xBe_{13}$ under pressure [16], see fig. 2. Better quantitative agreement between theoretical (fig. 1) and experimental (fig. 2) curves can be obtained if one assumes that at least one more parameter depends on pres-

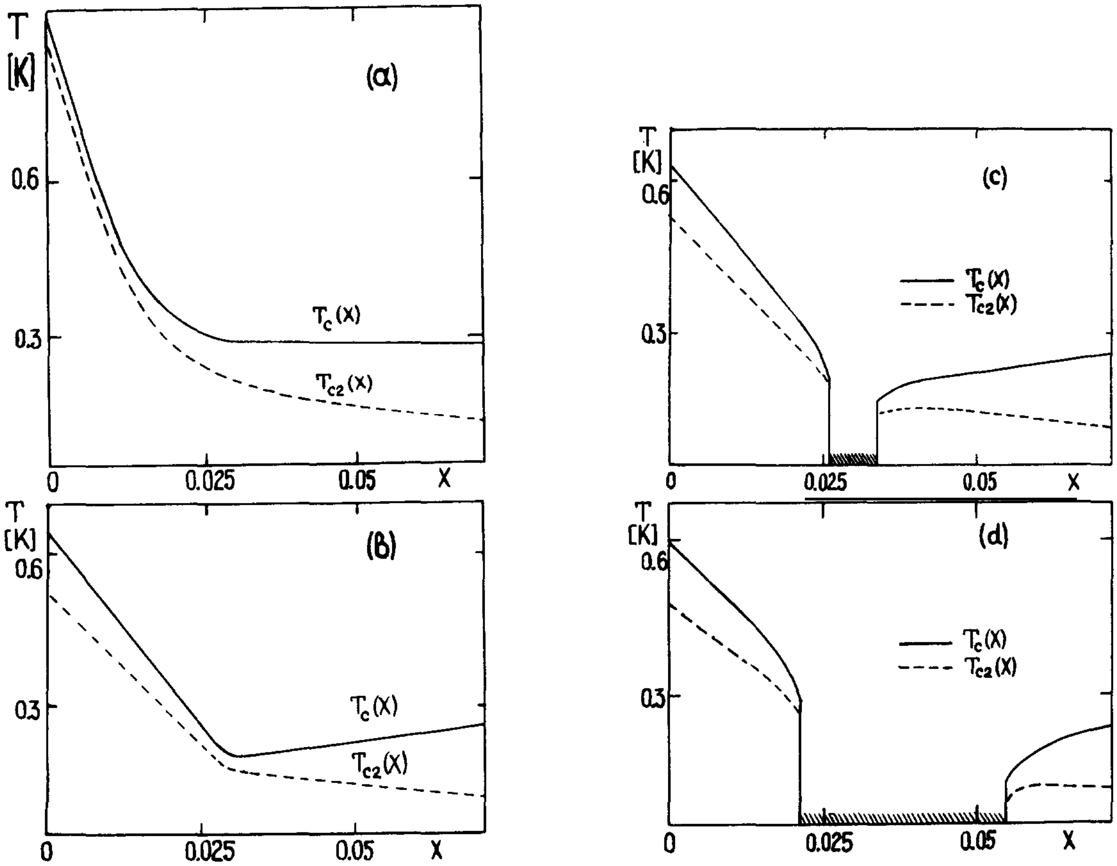

Fig. 1. Dependence of the critical temperatures of superconductivity T_c and coherent Kondo screening T_{c2} on the Th concentration x in $U_{1-x}Th_xBe_{13}$ for different parameters in eqs. (2) and (3): For all diagrams $a_2 = 1.1$ K^{-1}, $T_{c,0} = 0.3$ K, $T_{c2,0} = 0.246$ K, $\lambda_0 = 0.0463$ and $C = 2.4$. a_1 is varied with pressure according to (a) $a_1 = 1$ K^{-1}, (b) $a_1 = 0.74$ K^{-1}, (c) $a_1 = 0.73$ K^{-1}, (d) $a_1 = 0.6$ K^{-1}. Reentrant behaviour appears in (c) and (d).

sure and concentration. The main features of $T_c(x)$ and $T_{c2}(x)$ dependences are persistent for the value of $2\nu = 4/3$ in eq. (4).

3. We shall discuss some consequences of the above phenomenological model:

(i) In the case of pure UBe_{13} it seems reasonable to interpret the sound attenuation peak at $T \sim 0.82$ K [17] to be associated with strong fluctuations and related to the transition into a coherent state at $T = T_{c2}$. The coincidence of T_c (~ 0.86 K) and T_{c2} (~ 0.82 K) might be used to describe the heat capacity behaviour $C(T)$ [2,18] in UBe_{13} in the following way: Below T_c, C versus T should be considered as a sum of contributions to the heat capacity both from superconducting and coherent transitions, leading to a higher specific heat anomaly [2,8] than predicted by the BCS theory. The power laws $C(T) \sim T^{2.88}$ [19], spin–lattice relaxation rate $1/T_1 \sim T^3$ [20], specific behaviour of the penetration depth $\lambda_L(T)$ [21] caused various suggestions as regards the character of superconducting gap zero-points [19] or lines [20,21]. One of the probable solutions of the discrepancy is that there are, in fact, lines of zeroes, but $C(T)$ data show T^3 dependence due to the summation of two contributions from coherent and superconducting transitions.

(ii) Since the transition at T_{c2} for $\lambda_{12} < 0$ tends to stay very close to T_c (see fig. 1a), it is quite natural to expect almost identical $T_c(H)$ and $T_{c2}(H)$ shifts in magnetic fields H, as has been observed in $U_{1-x}Th_xBe_{13}$ [22].

(iii) The shape of the $C(T)$ peak in $U_{1-x}Th_xBe_{13}$ alloys strongly resembles the form of a typical fluctuation rather than a mean field phase transition (see ref. [23]).

(iv) The huge derivative $dH_{c2}/dT|_{T\to T_c} \sim 200$–$500$ kOe/K in UBe_{13} [12] is to be expected for strong fluctuating superconductivity, which is almost insensitive to small magnetic fields.

(v) The enhancement of fluctuations at $T=T_c$ in UBe_{13} might be also responsible for the negative proximity effect reported in ref. [25].

(vi) In UBe_{13}, substitution of Th for U yields induced magnetic moments on Th sites in the coherent regime at $T \lesssim T_{c2}$ and in turn increases the interaction parameter λ_{12} in eqs. (2) and (3) at the same time. Moreover, for low concentrations x we expect the specific heat peak at $T=T_{c2}$ to be roughly proportional to x, in agreement with experimental data [23]. Hence Th in $U_{1-x}Th_xBe_{13}$ at $T=T_{c2}$ helps to observe the second transition in specific heat measurements. Eventually for $x \gtrsim 0.02$, λ_{12} changes sign and $T_c(x)$ and $T_{c2}(x)$ dependences are split, making possible the distinct observation of the second phase transition at T_{c2} for $x=0.02$–0.04 [23]. In $U_{1-x}Th_xBe_{13}$, the sound attenuation peak is seen exactly at $T=T_{c2}$ [24].

(vii) We suggest that the branching of $T_{c2}(x)$ from $T_c(x)$ is a privilege of Th impurities in UBe_{12} due to a much stronger T suppression by other impurities [26]. For these systems the complete suppression of superconductivity seems to occur before a branching of $T_{c2}(x)$ from $T_c(x)$ takes place.

(viii) The interaction of two phases in the vicinity of the crossing point is characterized by the same critical behaviour, independent of microscopic mechanisms, when fluctuations are strong [11]. In particular, the length scales for superconductivity and collective Kondo screening at temperatures close to $T_c=T_{c2}$ should be of the same order.

The main features of $T_c(x)$ behaviour in $U_{1-x}Th_xBe_{13}$ (fig. 2), especially the appearence of reentrant superconductivity under pressure, are qualitatively reproduced by our model (compare figs. 1 and 2). As for the nature of the second $C(T)$ peak at T_{c2} we do not need any suggestions about a possible second exotic superconducting phase [2]. The peak below T_c simply corresponds to the transition into a coherent state, which at the same time renders induced magnetic moments at Th sites for temperatures below T_{c2}.

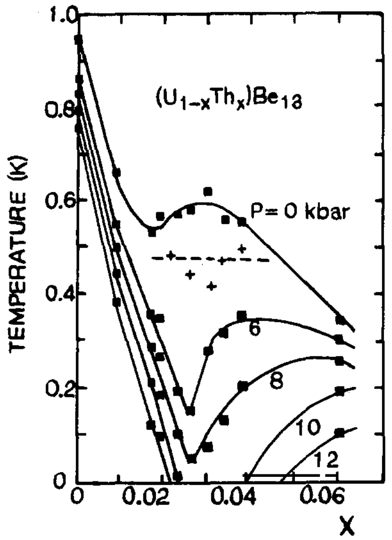

Fig. 2. Pressure dependence of the superconducting transition temperature T_c on Th concentration x, obtained in ref. [16]. These curves compare with the model results, fig. 1. For pressures of 10 and 12 kbar a reentrant behaviour appears. The dashed line indicates the critical temperature for the second phase transition T_{c2} at ambient pressure, as obtained in ref. [23].

Even in the case when the superconducting transition is treated in terms of the standard (nonfluctuating) approach, the coupling of the low-lying strongly fluctuating transitions greatly enlarges the critical region for the superconducting order parameter, in which fluctuations should be observable [27].

In conclusion, we would like to emphasize the fluctuating character of superconductivity in HFS, associated with the extremely high density of electrons participating in the superconducting transition. In this context, the interaction of superconductivity with another collective phenomenon, coherent Kondo screening, will be very important, especially for those HFS where T_{c2} is close to the superconducting transition temperature.

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