Heavy Fermionic Systems

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I. HEAVY FERMIONS: AN OVERVIEW

Heavy fermion systems are a collection of materials whose properties are governed by a lattice that carries f-electron magnetic ions at crystalographic sites [Hewson]. The electrons within these f-orbitals interact magnetically with conduction electrons within the system. This state of affairs is reminiscent of the single ion Kondo problem that arises in other systems. In single ion Kondo systems, magnetic impurities are injected into the system which then interact with passing conduction electrons. This results in a direct exchange coupling J forming between the localized spin impurities and those of the itenerant conduction electrons. A direct consequence of this interaction is that resistivity within these systems breaks from standard Fermi liquid theory at very low temperatures . Moreover, with decreasing temperature the coupling can screen out the spin impurities by binding conduction electrons to them to form a singlet state. The temperature at which this screening occurs is refered to as the Kondo Temperature, T_K. Jun Kondo was the first to solve this problem and provide a logarithmic correction term to the resistivity [Kondo]. This term accurately accounted for the peculiar upturn in resistivity at low temperatues but also asymptotically diverged as the temperature was suppressed to absolute zero. Further work by others succeded in fixing the divergence and today the solution to the Kondo problem stands as an impressive achievement in solid state physics.

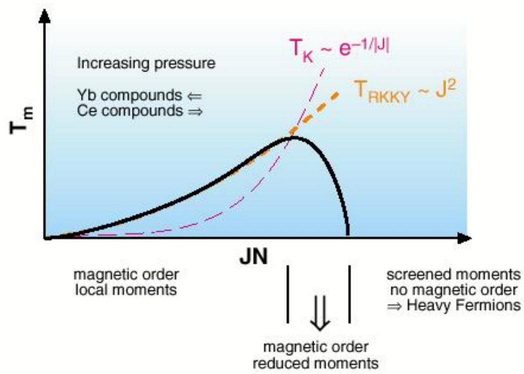

FIG. 1: Possible ground states resulting from the competition between the Kondo and RKKY interactions. T_m is the ordering temperature and JN is the f-d exchange coupling times the f density of states at the Fermi energy. [Sanchez]

A fundamental and very important difference between single ion Kondo systems and heavy fermion systems is that in the former, the ions exist as impurities scattered within the system and, as a result, interactions are isolated short range events. However, in heavy electron systems, the f-orbitals are part of the crystalographic structure and form a Kondo lattice of magnetic ions. Although this may at fist appear to be a direct extension of the original Kondo problem, in practice a solution is much less tractable. Moreover, due to the periodicity of the magnetic lattice, an indirect exchange coupling mediated by the conduction electrons is established between the sites; the so-called RKKY interaction [Kittel]. Indeed, many models of heavy fermion systems are treated as a competition between a RKKY interaction that acts to set up long range magnetic order at a temperature T_{RKKY} and a Kondo effect that begins screening the sites as the temperature drops below T_K. In the Kondo effect the onset of magnetic screening is given by

$$T_K = \rho^{-1} e^{-\frac{1}{\rho J}} \tag{1}$$

where ρ is the density of states at the Fermi surface and J is the exchange coupling between the conduction electrons and the localized magnetic f-orbitals. However, the onset temperature for magnetic ordering due to the RKKY interaction goes as

$$T_{RKKY} \propto J^2 \rho \tag{2}$$

Consequently, as the temperature is supressed, the moments associated with the long range magnetic order will begin to be screened away (Fig 1) as the Kondo interaction begins to dominate the system.

Aside from T_{RKKY} and T_K, there exists a third temperature that plays a vital role in determining the onset of property changes within heavy fermion systems. This temperature, which is referred to as T^*, corresponds to a point where the bound f-electrons become (at least partially) itenerate. Interestingly, it is for temperatures below T^* that the f-electrons begin to unbind. T^* is typically very low, around 1-10K depending on the system [Fisk]. A calculation of the change in entropy over this temperature range reveals a sharp climb which is attributed to this unbinding process. The magnitude of this change is fairly consistent from system to system, around $Rln(2)$ where R is Rydbergs constant, and is accompanied by significant changes in properties such as reduced resistivity, modified spin sucseptibility, an observed Knight shift, etc [Yang, PRL]. Because of this, it is convinient to define T^* as

$$\Delta S = \int_0^{T_*} \gamma dT = Rln(2) \tag{3}$$

Recent work [Yang, Nature] has demonstrated that T^* can be modeled very well as

$$T^* = cJ^2\rho \tag{4}$$

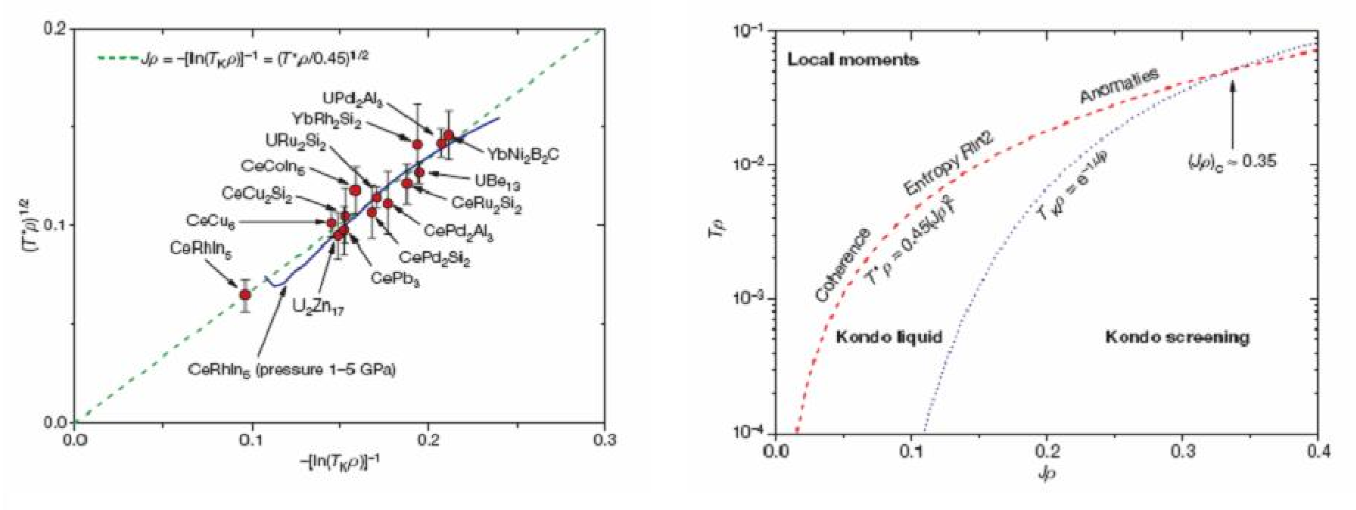

FIG. 2: a) Confirmation of T^* given by the intersite RKKY interaction for a variety of Kondo lattice materials; c = 0.45. b) Updated Doniach diagram for Kondo lattice materials. [Yang]

where c is a parameter to be determined. Combining this with (1) gives the relation:

$$[ln(T_K\rho)]^{-1} = \sqrt{c^{-1}T^*\rho} \tag{5}$$

A value of c = 0.45 was determined by fitting (5) to experimenal values of T^*, T_K, and γ for a variety of Kondo lattices (Fig 2a). From this, a modified version (Fig 2b) of the Doniach diagram [Doniach] was generated that relates the general phase diagram behavior of the system to the fundamental quantities that drive this behavior.

II. PROPERTIES AND GROUNDSTATES

Surprisingly, there exist a vast number of ground states that support the formation of heavy fermions from system to system. For instance, systems such as UBe_{13} form a Non-Fermi Liquid which crosses over into an exotic unconventional superconductivity at very low temperatures. However, other systems like UPt_3 begin as an antiferromagnet that transition to a heavy Fermi Liquid phase below T_n before ultimately crossing over to superconductivity at even lower temperatures. Others like $CeAl_2$ and U_2Zn_{17} remain as antiferromagnets at very low temperatures while some like CeNiSn become narrow-gap semiconductors [Misra]. Work on a doped variant of the parent compound $URu_{2-x}Re_xSi_2$ has uncovered new and equally interesting phenomena. For example, $URu_{2-x}M_xSi_2$ (M = Re, Tc) revealed the first instance of a Fermi surface (FS) instability in a heavy fermion system [Bauer]. Re doping, while quickly suppressing the SC phase, also results in a region of non-Fermi-liquid (NFL) behavior. Of particular interest, the NFL behavior is exhibited across a FM quantum critical point (QCP), which is in direct contrast to the majority of studies that exhibit NFL behavior across SG or AFM QCPs [Aronson, Wilson]. Moreover, this NFL behavior extends deeply into the ferromagnetic phase, acting as a rare instance of deep NFL penetration into an ordered magnetic state. Specifically, the electrical conductivity ρ(T), magnetic susceptibility χ(T), and heat capacity C(T) all exhibit logarithmic or power law behavior (hallmarks of NFL behavior) well within both the paramagnetic and ferromagnetic phases, as can be seen in Fig. 3 for ρ(T) [Bauer].

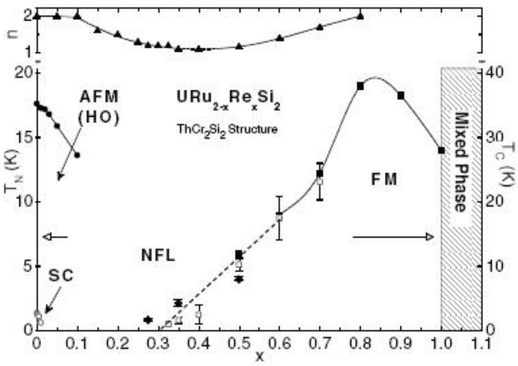

FIG. 3: The power law exponent of ρ, n, is shown in the top part of the phase diagram and clearly deviates from n = 2 Fermi Liquid behavior both at the QCP and well within the ordered ferromagnetic state. [Bauer]

The lack of uniformity in properties and ground states of heavy fermion systems provide a rich playground for experiment while simultaneously making the search for a comprehensive theory of the systems a daunting task. Even systems with striking similarities in their phase diagrams exhibit a wide range of unique characteristics. For example, consider $Y_{1-x}U_xPd_3$ [Seaman] and $Sc_{1-x}UxPd_3$ [Wilson], both of which exhibit Non-Fermi liquid behavior after U-doping to include an f-electron impurity. Yttrium and Scandium are both very similar in that they carry a single d-electron in their valence shell and a similar phase diagram of temperature versus U-concentration have been established as shown in Fig. 4 [Gajewski]. In general, the onset of non-Fermi liquid behavior in heavy fermion systems occurs immediately after crossing a quantum critical point and for these systems in particular the onset of the NFL behavior is established across a spin glass QCP. Theory focused on ascribing heavy fermion behavior to the rise in magnetic fluctions in the systems as the systems approaches a QCP has found much support in the scientific community. Systems whose NFL behavior is governed by a second order phase transition across a quantum critical point should obey ω/T scaling in their dynamic susceptibility of the form:

$$\chi^{"}(q,\omega,T) = \frac{1}{AT^{\alpha}F(\omega/T)} \quad , \qquad F(\omega/T) = e^{\alpha\ \Psi(\frac{1}{2}-\frac{i\omega}{2\pi T})} \tag{6}$$

This ω/T scaling has been seen in several systems, including $Sc_{1-x}U_xPd_3$ (α=1/5) [Wilson], $UCu_{5-x}Pd_x$ (α=1/3) [Aronson], and $CeCu_{5.9}Au_{0.1}$ (α=3/4) [Schroder]. A striking result is that the first of these systems undergoes a spin glass QCP while the other two undergo an anti-ferromagnetic QCP. As mentioned earlier, a ferromagnetic QCP has

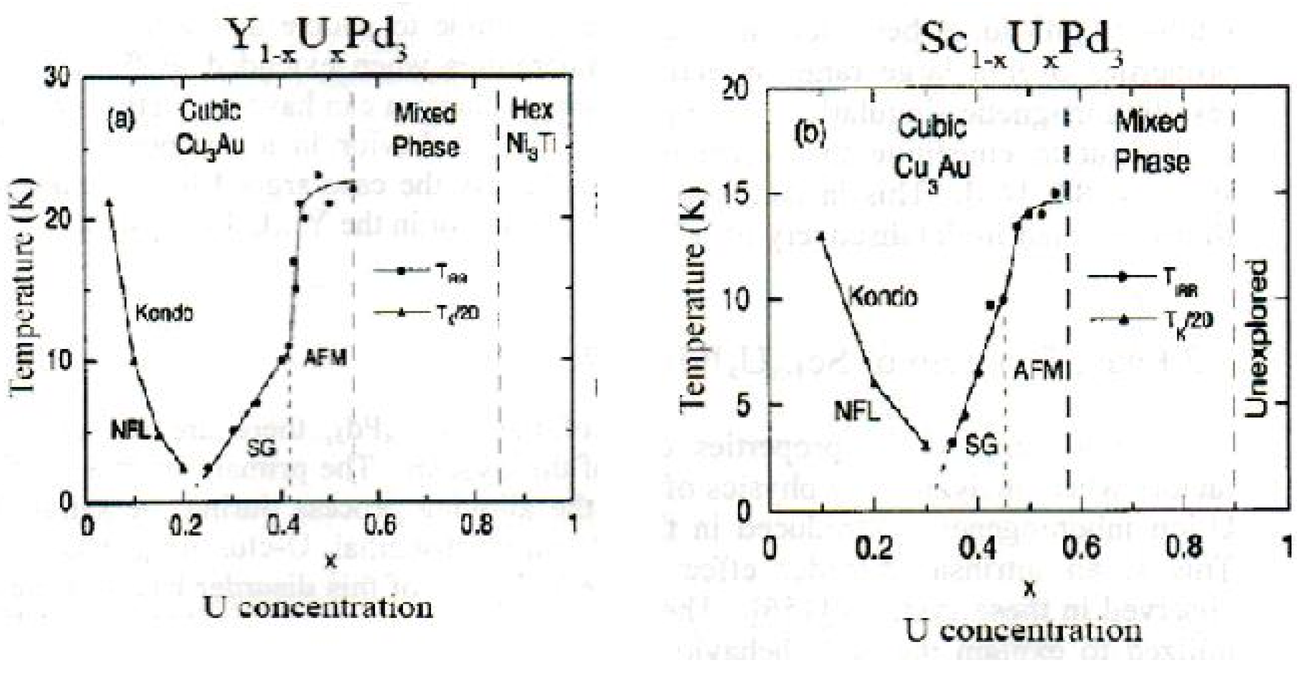

FIG. 4: Two heavy fermion systems with related phase diagrams [Gajewski].

also been observed at the onset of NFL behavior in $URu_{2-x}Re_xSi_2$. Consequently, the details of the phase transition appear to be unimportant, with the only pervading characteristic being the existence of QCP in the phase diagram.

The motivation to ascribe NFL behavior to the existence of a QCP is quite natural due to their ubiquity in these systems. Under such a scheme, critical fluctuations in the system extend over increasing length scales as the temperature is suppressed to zero. These long range interactions lead to behavior that departs from standard Fermi-liquid or mean field theory. For many systems, quantum criticality stems from the competition between a Kondo effect and RKKY interactions. Indeed, a viable microscopic description for the origin of NFL behavior in many systems is a model predicated on competition between these two interactions within disordered, anisotropic regions. However, such a model has yet to be verified definitively or applied universally. Even in the case of $Y_{1-x}U_xPd_3$ and $Sc1-xUxPd_3$, models originally applied to one have failed to apply to the other. Indeed, for $Sc_{1-x}U_xPd_3$ alternative explanations for the NFL behavior that have been ascribed successfully to $Y_{1-x}U_xPd_3$ do not work. Namely, the cubic crystalline electric field (CEF) of the parent compound YPd_3 splits the J=4 degeneracy of the U4+ ions into Γ4 and Γ5 triplets, a Γ1 singlet, and a Γ3 doublet [Lea]. A quadrapolar Kondo effect (QKE) can then be applied to explain the NFL behavior of the doped system [Seaman]; although some controversy does exist over such an interpretation due to conflicts between a (QKE) predicted Γ3 nonmagnetic ground state and neutron experiments that reveal a magnetic ground state [Dai]. However, for the $ScPd_3$ parent compound, (CEF) models are entirely incompatible with inelastic neutron scattering data [Wilson]. The qualitative similarity in the phase behavior of $Y_{1-x}U_xPd_3$ and $Sc_{1-x}U_xPd_3$ reinforces the idea that a similar, if not identical, mechanism should be responsible for establishing an NFL phase in both of these systems. Consequently, the incompatibility of a CEF model in $Sc_{1-x}U_xPd_3$ makes a working (but controversial) CEF model in $Y_{1-x}U_xPd_3$ even less convincing. A second model also exists that fails to explain NFL behavior in both systems. $Y_{1-x}U_xPd_3$ suffers from an intrinsic disorder due to the inhomogeneous doping of Uranium into the system. This led to speculation that the disorder was a potential source of the NFL behavior in this system. However, $Sc_{1-x}U_xPd_3$ shows no such inhomogeneity. Considering the similarity of their phase diagrams, disorder became a much less favorable candidate for explaining the NFL behavior in either of these systems. The conficiting hyposthesis between two closely related systems underlines the inherent difficulty of properly describing a broad theory that can be applied to a host of heavy fermion systems.

The study of $Sc_{1-x}U_xPd_3$ has also led to some surprising inconsistencies in the magnetic scattering. Specifically, inelastic neutron scattering experiments designed to map the doping evolution of spin excitations revealed a featureless magnetic scattering background over broad spectrums of Q. Indeed, even at x = 0.48 doping most of the scattering resided in an incoherent magnetic signal mixed with a modulated signal intrinsic to SG order. This is in contradiction to the bulk susceptibility phase diagram that predicted AF Bragg peaks at that concentration. This supports a model consistent with non-interacting U-ions or at most U-ions in short range correlated clusters showing only slight long range AF order at high doping concentrations. This scheme is strongly supported by the scattering data. However, the interpretation becomes worrisome when compared to a plot of integrated intensity versus U-concentration as shown in figure 5 [Wilson]. The current picture describing the NFL behavior as a QCP magnetic signal surrounded by incoherent magnetic scattering due to non-interacting U-ions implies that one would expect the integrated intensity to increase linearly when doping in these non-interacting ions. Consequently, there exists an inherent inconsistency in interpreting the magnetic scattering as arising solely from isolated U sites.

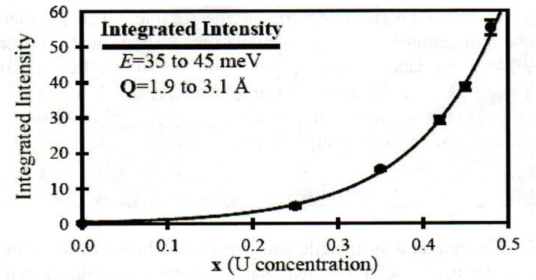

FIG. 5: Exponential rise in scattering as a function of doping in direct contradiction to non-interacting f-election sites. [Wilson]

The above examples of heavy electron systems provide a glimpse of some of the considerations that must be accounted for when studying these systems, while simultaneously reinforcing the fact that many of the results conflict with current ideas about why the systems form in the first place. As it stands, intense study of heavy fermions has dropped off as new areas, such as the pnictide superconductors, have taken center stage. However, these heavy electron systems are still rife with unsolved puzzles and offer a unique and interesting physics. It will be interesting to see how their story unfolds as research continues to slowly unravel the details of these compounds.

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