Understanding the Heavy Fermion Phenomenology from a Microscopic Model

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We solve the 3D periodic Anderson model using a two impurity cluster dynamical mean field theory. We obtain the temperature versus hybridization phase diagram. Approaching the quantum critical point (QCP) both the Néel and lattice Kondo temperatures decrease and they do not cross at the lowest temperature we reached. While strong ferromagnetic spin fluctuation on the Kondo side is observed, our result suggests the critical static spin susceptibility is local in space at the QCP. We observe in the crossover region a logarithmic temperature dependence in the specific heat coefficient and spin susceptibility.

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The heavy fermion phenomenon is among the most intensively studied subjects in condensed matter physics [1]. Experimental information accumulated over the past 30 years has revealed many unconventional aspects of heavy fermion physics [1,2].

Heavy fermion physics is derived from a local moment band hybridizing with an extended conduction band. The hybridization induces the competing Kondo ($J_K$) and RKKY interactions [3]. At $J_K \ll T$, the physics is dominated by those of the two bands separately with the hybridization as a perturbation. As the temperature is lowered, depending on the strength of $J_K$, different physics may develop. In the region where $J_K \ll W$ (the conduction bandwidth), the RKKY interaction prevails. The RKKY is a long range exchange mediated by the conduction electrons near the Fermi surface and oscillates with $k = 2k_F$ asymptotically. At $T \leq J_K^2 / W$, the RKKY interaction can induce a transition to a magnetically ordered phase. In the crossover regime on this side, the Kondo behavior, though subdominant, would still show up in various measurables.

Heavy fermion physics begins at high temperatures $T > \Delta / V$, where $\Delta$ is related to the Fermi energy.

$T_N$ is related to the lattice Kondo energy, $T_0$, which is proportional to the Fermi energy. $T_0$ can be defined in terms of the saturated homogeneous spin susceptibility in the Fermi liquid phase, $\chi_{\bar{L}=0} = C/T_0$, where $C$ is the Curie constant. Should $T_0$ approach zero, strong FM spin fluctuation would be observed.

It turns out that the thermodynamics related to the continuous condensation of the local moments into the heavy fermion fluid is quite universal [5], as contrasted with that more material specific in the low temperature region. At low temperatures, various phases may develop, including a superconducting phase [1]. One interesting possibility is that the competing Kondo and RKKY interactions result in a quantum phase transition (QPT) without the interference of any other phases. Such a situation is observed experimentally in CeCu$_{1-x}$Au$_x$ [6,7] and YbRh$_2$Si$_2$ [8].

Two different scenarios have been proposed for the heavy fermion QPT. One is the Hertz-Millis-Moriya theory [9], which is applicable when the energy scale of the Kondo screening is much higher than that of the magnetic ordering near the quantum critical point (QCP). As a result, the local moments are fully screened in the quantum critical regime. However, by comparing the predicted critical exponents with the experiments, the theory is found to be inapplicable in many cases [1,2,7,8]. In the second scenario, one expects the magnetic ordering ($T_N$) and lattice Kondo screening ($T_0$) energies vanish simultaneously at the QCP [2,7] (see also [10]). Since the critical spin fluctuations at $k = 2k_F$ and $k = 0$ are associated with the $T_N$ and vanishing $T_0$, respectively, it is interesting to know what kind of critical magnetic mode may develop in the neighborhood of the QCP. We will show later that near the QCP the two critical modes strongly interact with each other and the critical spin fluctuation becomes local in space (see also [2]). In this scenario, the local moments survive at the QCP.

It is important that the heavy fermion physics as we understand be obtained from a microscopic model. To this end the main theoretical difficulty lies in treating on equal footing the Kondo and RKKY interactions. Many bosonic
mean field theories, like the Hertz-Millis-Moriya [9] and slave-boson [11] theories, fail because they rely on the order parameter of either the magnetic or Kondo phase and miss the properties of the other. On the other hand, a fermionic mean field theory, like the dynamical mean field theory (DMFT) [12], would allow the possible orders to develop and compete and is more desirable.

The DMFT extends the Weiss mean field theory to describe fermions. The single impurity DMFT was applied to both the Kondo and magnetically ordered phases in heavy fermions [13] (see also Ref. [14]). Besides the over estimation of the Neel temperature due to the lack of the magnetic fluctuation, this approach cannot capture properly the renormalization of the RKKY interaction. A partial solution to this problem is to add an RKKY interaction to the model and allow the renormalization of the RKKY by extending the Weiss approximation to the interaction [15]. The resulted formalism, the so-called extended DMFT (EDMFT), is able to describe qualitatively the heavy fermions [16] in both the Kondo and antiferromagnetic (AFM) phases. However, it predicts a first order phase transition due to the local mean field treatment of the RKKY interaction [16,17].

A parallel path in studying the heavy fermions is through the two impurity problem. This model contains the dynamics of both the Kondo and RKKY interactions and is solvable [18,19]. It was shown, at the particle-hole symmetry, there was a non-Fermi liquid fixed point separating the Kondo and magnetic phases. However, it is difficult to extend the properties of a multi-impurity model to a lattice of impurities.

In this Letter, we combine the two impurity model with DMFT so that a lattice of impurities can be described. This overcomes the difficulties of both the single impurity DMFT and multi-impurity approaches. It handles the Kondo and RKKY interactions in a more balanced way.

We consider the periodic Anderson model in 3D

\[
H = \sum_{\mathbf{k},\sigma} (\epsilon_{\mathbf{k}} - \mu) c_{\mathbf{k},\sigma}^\dagger c_{\mathbf{k},\sigma} + (E_f - \mu) \sum_{j,\sigma} n_{j,\sigma}^f \\
+ U \sum_{f} (n_{j,1}^f - 1/2)(n_{j,1}^f - 1/2) + V \sum_{j,\sigma} (f_{j,\sigma}^1 c_{j,\sigma}^\dagger + c_{j,\sigma}^\dagger f_{j,\sigma})
\]

(1)

with \(\epsilon_{\mathbf{k}} = -\frac{1}{2} \sum_{i=1}^{3} \cos k_i\). We divide the lattice into two interpenetrating sublattices, A and B. The unit cell is then doubled. Applying the cavity method [12], we obtain an effective local action:

\[
S_0 = -\int_0^\beta d\tau \int_0^\beta d\tau' \sum_{\mathbf{r},\sigma,y} f_{\mathbf{r},0,\sigma}(\tau) [G_{y\alpha,\sigma}(\mathbf{r}) - \mathbf{i}(\tau - \tau')] f_{\mathbf{r},1,\sigma}(\tau') + U \int_0^\beta d\tau \sum_{\mathbf{r},\sigma,y} \left( n_{\mathbf{r},0,\sigma}^f - \frac{1}{2} \right) \left( n_{\mathbf{r},1,\sigma}^f - \frac{1}{2} \right)
\]

(2)

The Weiss function \(G_{y\alpha,\sigma}(\mathbf{r})\) is determined self-consistently as follows. First, we use the quantum Monte Carlo method (QMC) [20] and obtain the impurity Green’s function. Then, from the Dyson equation for the impurities, we get the impurity self-energy \(\Sigma_{\sigma}^{\text{imp}}(\mathbf{r})\). The lattice self-energy is constructed to be:

\[
\Sigma^{ff}_{\sigma}(\mathbf{k}, \mathbf{r}_n) = -2d\delta f^f(\mathbf{k}) \sum_{\sigma,\mathbf{r},\mathbf{r}'} G^{\text{imp}}_{\mathbf{r} \mathbf{r}' \sigma}(\mathbf{k})(\mathbf{r}_n, \mathbf{r}') \sum_{\sigma,\mathbf{r},\mathbf{r}'} G^{\text{imp}}_{\mathbf{r} \mathbf{r}' \sigma}(\mathbf{k})(\mathbf{r}_n, \mathbf{r}')
\]

(3)

from which we obtain the local Green’s function:

\[
G^{ff}_{\mathbf{r}_n,\sigma}(\mathbf{r}_n) = \sum_{\mathbf{k}} \left[ \begin{array}{cc}
ip_n + \mu - E_f & \epsilon_{\mathbf{k}} f(\mathbf{k}) \\
0 & ip_n + \mu - E_f
\end{array} \right]^{-1} - \frac{V^2}{(ip_n + \mu)^2 - \epsilon_{\mathbf{k}}^2 f(\mathbf{k}) (ip_n + \mu)} \left[ \begin{array}{cc}
\epsilon_{\mathbf{k}}^2 f(\mathbf{k}) & \epsilon_{\mathbf{k}}^2 f(\mathbf{k}) \\
\epsilon_{\mathbf{k}}^2 f(\mathbf{k}) & \epsilon_{\mathbf{k}}^2 f(\mathbf{k})
\end{array} \right]^{-1},
\]

(4)

with \(f(\mathbf{k}) = \exp(ik_z)\). By identifying the local (within a unit cell) Green’s function on the lattice with that of the impurity model, we form a self-consistent loop [12,21]. While solving the impurity model we can measure the z-direction spin susceptibility \(\chi_{\mathbf{r} \mathbf{r}'}(\tau) = \langle T_\tau \hat{S}_z(\mathbf{r}, \tau) \hat{S}_z(\mathbf{r}', 0) \rangle\), with \(\hat{S}_z(\mathbf{r}) = n_{\mathbf{r},1} - n_{\mathbf{r},1}^\dagger\). The lattice susceptibility is obtained in the same way as the lattice self-energy given in Eq. (3).

We study the phase diagram of temperature versus \(V\) at fixed \(U = 1.2\) and \(E_f = -0.15\). To avoid crossing the band gap, we change the chemical potential \(\mu\) along with \(V\) so that the free (\(U = 0\)) particle density per site at \(T = 0\) is fixed at \(n_{\text{tot}}^{\text{free}} = 2.5423\). The resulting physical density changes slightly as \(V\) increases and is always greater than and close to the half-filling [22]. We study the AFM and paramagnetic (PM) phases and the transition between them. In solving the two impurity problem we use the QMC. We always use \(U\Delta \tau \approx 1\) where \(\Delta \tau = \beta/L\) and \(L\) is the number of time slices in QMC. In each DMFT iteration, we perform QMC sweeps \(\sim 10^5\). Away from the phase transitions around 10 DMFT iterations are usually enough to converge the results. Near the phase transition a lot more are needed due to the critical slowing down.

Figure 1 is the phase diagram we obtained. Two technical remarks are in place. First, to exam if the AFM to PM transition is continuous, we checked the inverse static spin susceptibilities at \(\mathbf{k} = (\pi, \pi, \pi)\) which becomes very critical at the corresponding transition values of \(V\) [22].
the static homogeneous spin susceptibility at $V = 0.28$.

Second, the crossover temperature $T_0$ is obtained using the saturated static homogeneous spin susceptibility $\chi(k = 0, i0) \to C/T_0$ at low temperatures. [We used Curie constant $C = 1/2$ which is obtained in the high temperature limit.] An example of the saturation behavior is shown in the inset of Fig. 1. Note that $T_0$ being small means the FM spin fluctuation becomes very strong.

To study the critical magnetic fluctuation around the QPT, we plot in Fig. 2 $\chi_{AA}(i0)$ and $\chi_{AB}(i0)$. From the result we see that $\chi_{AB}(i0)$ changes sign when $V$ increases. This sign change is actually a special manifestation of a more general evolution of the RKKY correlation from being AFM to FM as hybridization is increased, due to the conduction electrons mediating the RKKY change from those near the Fermi surface to around the band bottom. Meanwhile, $\chi_{AA}(i0)$ is always positive and becomes very strong in the crossover regime. This scenario should extend to $T = 0$ and there would be a point at which $\chi_{AB}(i0) = 0$ and $\chi_{AA}(i0)$ becomes critical. This point would be the heavy fermion QCP. Note that the static spin response being local does not mean a local electron self-energy. Actually, the nonlocal self-energy becomes stronger at lower temperatures [22].

An important question is how the Kondo screening, which is local in space, becomes coherent in heavy fermions. To this end, we study the evolution of the spatial correlation in the spin responses in Fig. 3. It shows that spin fluctuations are quite local in space down to $T \sim 0.1$ for $V = 0.26$ and $T \sim 0.2$ for $V = 0.50$. At lower temperatures, the FM spin susceptibility becomes dominant. This is similar to that observed experimentally in YbRh$_2$Si$_2$ [23]. Two remarks are in place. (1) the nonlocality in $\chi$ develops at lower temperature for $V = 0.26$ than that for $V = 0.50$. This reflects the local nature of the spin fluctuation in the quantum critical region which extends to lower temperature as the QCP is approached and is consistent with Fig. 2. (2) The logarithmic temperature dependence of $\chi$ as shown in Fig. 3 is similar to those observed in many heavy
fermion compounds [1]. A logarithmic temperature dependence is also found in the total energy shown in Fig. 4.

To conclude, using a two impurity DMFT, we studied the periodic Anderson model on cubic lattice at finite temperatures. We obtained the phase diagram which is consistent with the picture that both the Neel and Kondo temperatures vanish at the QCP. As the QCP was approached from the Kondo side, we found strong ferromagnetic spin fluctuations. From the sign change of the static nearest-neighbor spin susceptibility, we conjectured that the critical static spin fluctuation become local at the heavy fermion QCP. We explored the crossover region and observed logarithmic temperature dependences in the specific heat coefficient and spin susceptibility.

Our results presented in this Letter imply that a two impurity Anderson model combined with DMFT might serve as a minimal model in describing most of the thermodynamics of the heavy fermions. However, we should also keep in mind that there could exist other phases, like the superconducting phase [24], whose existence may also keep in mind that there could exist other phases, like.

Since the $4f$ or $5f$ orbital contributes to the local moment physics in most heavy fermion compounds, there are inevitably many physical properties oversimplified by the periodic Anderson model, like the orbital degeneracy and the subsequent crystalline field splitting. These are further complicated by the spin-orbital coupling, lattice frustration, disorder, hybridization with other bands, etc. Actually all these contribute to a much richer physics observed in experiments [1] than we have obtained. To this end, our current work can be considered as a useful guide to distinguish the “universal” heavy fermion features from those specific to individual materials.

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