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On the competition between magnetic order and local Kondo effect in Kondo lattice

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Abstract. We investigate the competition between magnetic order and local Kondo effect in a Kondo lattice model (*i.e.* the Coqblin-Schrieffer Hamiltonian extended to a lattice) in a mean-field approximation, taking account of the spin-orbit degeneracy $N_{\text{s.o.}}$ of each localized f level. This leads to the definition of a $N_{\text{s.o.}}$ dependent Kondo temperature. We study the Kondo phase and compare its energy with the energies of magnetic phases, when the number of the conduction band electron per site is near one. We present a phase diagram which shows the occurrence of three phases: Kondo, antiferromagnetic and paramagnetic phases. Our model in the mean-field approximation also shows a somewhat flat Kondo temperature, for large values of $N_{\text{s.o.}}$, as a function of the exchange coupling J between conduction and localized f electrons. Finally we show some scaling effects between $N_{\text{s.o.}}$ and J and we define a corresponding Kondo temperature.

PACS. 75.20.-g Diamagnetism and paramagnetism – 75.30.-m Intrinsic properties of magnetically ordered materials – 71.28.+d Narrow-band systems; intermediate-valence solids

1 Introduction

Strong electron-electron correlations are one of the central issues of the current condensed matter physics. The so-called heavy fermion materials represent a typical example of systems in which strong correlation effects are essential in determining their physical properties. The Kondo lattice model (KLM) [1] or periodic Anderson model (see for example [2] and Refs. therein) has attracted much interest as one of the typical standard models of heavy fermion systems. But the theoretical description of KLM remains an unsolved problem of solid state physics.

The KLM describes the interaction, at each lattice site, between conduction electrons and a localized f electron. In this model, conduction electrons are spin polarized due to their exchange interaction with the f electrons. This spin polarization propagates from one lattice site to another one. Since f electrons are present at each lattice site, two of these sites are related by the indirect exchange interaction between f electrons. This interaction is the so-called Ruderman-Kittel- Kasuya-Yosida (RKKY) interaction. So, there is a competition between local Kondo effect and magnetic order.

On the other hand, in weakly hybridized Kondo compounds, there are three types of compounds at low temperatures: nonmagnetic, antiferromagnetic and ferromagnetic ones. The nonmagnetic compounds, such as $CeAl_3$, $CeCu_2Si_2$ [3–5] are characterized by a heavy fermion behavior. The effects of the magnetic Ce ions are compensated by the conduction electrons, which leads to a Fermi liquid behavior at low temperature. Enormous values of the static magnetic susceptibility and the electronic specific heat coefficient are found as well as a T^2 behavior for the magnetic resistivity at very low temperatures. The other compounds such as $CeAl_2$ [6,7] and CeAg, $CeGe_2$ [8,9] have antiferromagnetic order or ferromagnetic order respectively.

In order to explain these anomalous behaviors, several authors studied the present problem of the competition between local Kondo effect and magnetic order. Doniach [10] introduced, many years ago, the KLM and applied the mean-field theory to this model. He described the competition between Kondo effect and RKKY interactions by a presently well known "Doniach diagram". He found that the ground state is antiferromagnetic for low $|J|/D$ and it is a nonmagnetic Kondo singlet for large $|J|/D$ (where J is the exchange parameter between 4f and conduction electrons and D is the half width of the conduction band). Lacroix and Cyrot [11] used the Functional Integral method with a static approximation for the non-degenerate f orbital case. Coleman [12] used the Scaling theory for the degenerate f case of the total angular momentum. Recently, Iglesias *et al.* [13] reexamined the Doniach diagram by including explicit short-range antiferromagnetic interactions in the Kondo lattice for the nondegenerate case and found a rather flat Kondo temperature, as observed in several cerium compounds.

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In this work, we treat the problem for the degenerate f case of the total angular momentum and use a mean field approximation as well as order parameters both for the local Kondo effect and the magnetic orders. For example, at a cerium site, the large spin-orbit coupling leads to a ground state of total angular momentum $I = l - \frac{1}{2}$ with $l = 3$. In this paper we disregard further splitting due to crystal field. In Section 2, we describe the KLM by adding the spin-orbit degeneracy of the localized f level and by using a mean-field approximation, but without explicitly including nearest-neighbor magnetic exchange interactions as was done in reference [13]. In Section 3, we study some ground state properties of the Kondo phase, in particular the Kondo binding energy and the Kondo lattice temperature T_1 at which the Kondo order parameter vanishes. In Section 4, we determine the ferromagnetic (antiferromagnetic) binding energy and the temperature $T_{\rm F}$ ($T_{\rm AF}$) at which the magnetic order parameter of the ferromagnetic (antiferromagnetic) phase vanishes. Finally, in Section 5, some consequences are calculated on the phase boundaries between the energetically favoured phases, *i.e.* the Kondo, antiferromagnetic and possible paramagnetic phases.

2 The Hamiltonian

We start with the Kondo lattice Hamiltonian [11,12] (*i.e.* the Coqblin-Schrieffer [14] Hamiltonian extended to a lattice), taking the spin-orbit degeneracy $N_{\rm s.o.}$ of each local level into account:

$$H = H_0 + H_{\rm I}, \tag{2.1}$$

$$H_0 = \sum_{\mathbf{k},\sigma} \varepsilon_{\mathbf{k}} c^\dagger_{\mathbf{k}\sigma} c_{\mathbf{k}\sigma} + E_0 \sum_{n,M} f^\dagger_{nM} f_{nM}, \tag{2.2}$$

$$H_{\rm I} = -\frac{J}{N} \sum_{n,\mathbf{k},\mathbf{k}'} \mathrm{e}^{\mathrm{i}(\mathbf{k}-\mathbf{k}')\cdot\mathbf{R}_n} \left\{ \sum_{M\sigma} \gamma_{\sigma M}(\mathbf{k}) \gamma^*_{\sigma M}(\mathbf{k}') \nu_{nM} c^\dagger_{\mathbf{k}\sigma} c_{\mathbf{k}'\sigma} + \sum_{(M,\sigma)\neq(M'\sigma')} \gamma_{\sigma M}(\mathbf{k}) \gamma^*_{\sigma' M'}(\mathbf{k}') f^\dagger_{nM'} f_{nM} c^\dagger_{\mathbf{k}\sigma} c_{\mathbf{k}'\sigma'} \right\}, \tag{2.3}$$

with

$$\nu_{nM} = f^\dagger_{nM} f_{nM} - \frac{1}{N_{\rm s.o.}} \sum_{M'} f^\dagger_{nM'} f_{nM'}, \tag{2.4}$$

$$\gamma_{\sigma M}(\mathbf{k}) = \sum_{m=-l}^{l} \langle lm \frac{1}{2} \sigma | IM \rangle Y^m_l(\mathbf{k}), \tag{2.5}$$

where $c^\dagger_{\mathbf{k}\sigma}$ creates a band electron with momentum $\mathbf{k}$ and spin component σ, $f^\dagger_{nM}$ creates an f electron in a total angular momentum state $|IM\rangle$ at site n, $\langle \dots | \dots \rangle$ is a Clebsch-Gordon coefficient, Y^m_l labels a spherical harmonic of rank l, J is the s-f exchange constant which is negative here, N designates the number of lattice sites. Equation (2.3) is an exchange-type Hamiltonian taking into account combined spin and orbital exchange scattering.

The quantity E_0 and the chemical potential μ must satisfy [11]

$$\frac{1}{N} \sum_{\mathbf{k}\sigma} \langle c^\dagger_{\mathbf{k}\sigma} c_{\mathbf{k}\sigma} \rangle = n_{\rm c}, \tag{2.6}$$

$$\frac{1}{N} \sum_{nM} \langle f^\dagger_{nM} f_{nM} \rangle = 1, \tag{2.7}$$

where $\langle A \rangle$ is a thermal average of A; $n_{\rm c}$ is the band electron number par site, which is assumed to be close to one. For example, in Ce intermetallic compounds (like $CeRu_2Ge_2$) about 10% of the $5d$ states are occupied, as it is more generally the case for many rare earth compounds. Thus the trivalent Ce ground state configuration is $4f^1 5d^1$.

We approximate equation (2.3) in the following way:

$$H_{\rm I} = \frac{J}{N} \sum_{n\mathbf{k}\mathbf{k}'} \mathrm{e}^{\mathrm{i}(\mathbf{k}-\mathbf{k}')\cdot\mathbf{R}_n} \left\{ \sum_{(M\sigma)\neq(M'\sigma')} \gamma_{\sigma M}(\mathbf{k}) \gamma^*_{\sigma' M'}(\mathbf{k}') \times \left(\langle f^\dagger_{nM'} c_{k'\sigma'} \rangle c^\dagger_{\mathbf{k}\sigma} f_{nM} + \langle c^\dagger_{\mathbf{k}\sigma} f_{nM} \rangle f^\dagger_{nM} c_{\mathbf{k}'\sigma'} - \langle c^\dagger_{\mathbf{k}\sigma} f_{nM} \rangle \langle f^\dagger_{nM'} c_{\mathbf{k}'\sigma'} \rangle \right) - \sum_{M\sigma} \gamma_{\sigma M}(\mathbf{k}) \gamma^*_{\sigma M}(\mathbf{k}') \times \left(\langle \nu_{nM} \rangle c^\dagger_{\mathbf{k}\sigma} c_{\mathbf{k}'\sigma} + \langle c^\dagger_{\mathbf{k}\sigma} c_{\mathbf{k}'\sigma} \rangle \nu_{nM} - \langle c^\dagger_{\mathbf{k}\sigma} c_{\mathbf{k}'\sigma} \rangle \langle \nu_{nM} \rangle \right) \right\}. \tag{2.8}$$

Then, at site n, the quantity

$$w_n = \frac{J}{\sqrt{N}} \sum_{\mathbf{k}M\sigma} \mathrm{e}^{\mathrm{i}\mathbf{k}\cdot\mathbf{R}_n} \gamma_{\sigma M}(\mathbf{k}) \langle c^\dagger_{\mathbf{k}\sigma} f_{nM} \rangle \tag{2.9}$$

is assumed to play the role of a Kondo effect order parameter, the quantities

$$v_n = \frac{1}{N} \sum_{\mathbf{k}\mathbf{k}'\sigma} \mathrm{e}^{\mathrm{i}(\mathbf{k}-\mathbf{k}')\cdot\mathbf{R}_n} \sigma \langle c^\dagger_{\mathbf{k}\sigma} c_{\mathbf{k}'\sigma} \rangle \tag{2.10}$$

and $\langle \nu_{nM} \rangle$ to play the role of magnetic order parameters.

3 The Kondo phase

For that phase, both magnetic order parameters vanish, since we recover a non-magnetic (singlet) regime.

$$v_n = 0 \tag{3.1}$$

$$\langle \nu_{nM} \rangle = 0. \tag{3.2}$$

Therefore, equation (2.8) gives

$$H_{\rm I} = \frac{J}{N} \sum_{n\mathbf{k}\mathbf{k}'} \mathrm{e}^{\mathrm{i}(\mathbf{k}-\mathbf{k}')\cdot\mathbf{R}_n} \sum_{(M\sigma)\neq(M'\sigma')} \gamma_{\sigma M}(\mathbf{k}) \gamma^*_{\sigma' M'}(\mathbf{k}') \times \left\{ \langle f^\dagger_{nM'} c_{\mathbf{k}'\sigma'} \rangle c^\dagger_{\mathbf{k}\sigma} f_{nM} + \langle c^\dagger_{\mathbf{k}\sigma} f_{nM} \rangle f^\dagger_{nM'} c_{\mathbf{k}'\sigma'} - \langle c^\dagger_{\mathbf{k}\sigma} f_{nM} \rangle \langle f^\dagger_{nM'} c_{\mathbf{k}'\sigma'} \rangle \right\}. \tag{3.3}$$

Equation (3.3) is an effective hybridization Hamiltonian term giving high density of states near the Fermi energy [11,15–17] which is observed in Kondo lattice systems [18]. In the following, we assumed $N_{\mathrm{s.o.}}$ to be large, therefore the restricted summation, in equation (3.3), on $M\sigma M'\sigma'$ can be approximately replaced by unrestricted one, and one obtains:

$$H_{\mathrm{I}} = \frac{1}{\sqrt{N}} w \sum_{n\mathbf{k}M\sigma} \left(\mathrm{e}^{\mathrm{i}\mathbf{k}\cdot\mathbf{R}_n} \gamma_{\sigma M}(\mathbf{k}) c^{\dagger}_{\mathbf{k}\sigma} f_{nM} + \mathrm{e}^{-\mathrm{i}\mathbf{k}\cdot\mathbf{R}_n} \gamma^{*}_{\sigma M}(\mathbf{k}) f^{\dagger}_{nM} c_{\mathbf{k}\sigma} \right) - N\frac{w^2}{J} \tag{3.4}$$

where w is given in equation (2.9) and is a constant for each site n.

Let us examine first the situation at absolute zero temperature. Minimization of the internal energy functional with respect to w or Green's function technique gives the Kondo effect order parameter:

$$w = D\mathrm{e}^{D/N_{\mathrm{s.o.}}J} \left(1 - \mathrm{e}^{2D/N_{\mathrm{s.o.}}J}\right)^{-1} \tag{3.5}$$

where a rectangular band density of states with D as half bandwidth has been used. The Kondo binding energy is then:

$$\frac{\Delta\xi_{\mathrm{K}}}{D} = N\mathrm{e}^{2D/N_{\mathrm{s.o.}}J} \left(1 - \mathrm{e}^{2D/N_{\mathrm{s.o.}}J}\right)^{-1} \tag{3.6}$$

This result is similar to that of Lacroix and Cyrot [11] who took the non-degenerate orbital case and used the Functional Integral method with a static approximation.

Now, let us look for the temperature T_1 at which the Kondo order parameter w vanishes. Green's function technique gives a B.C.S. type of equation for T_1 (similar to the equation giving the superconducting critical temperature) and one obtains:

$$1 = \frac{N_{\mathrm{s.o.}}J}{2D} \left\{ \log\left[\frac{N^2_{\mathrm{s.o.}}}{2(N_{\mathrm{s.o.}}-1)}\frac{T_1}{D}\right] + \frac{1}{N_{\mathrm{s.o.}}} \log\left[\frac{N_{\mathrm{s.o.}}}{2(N_{\mathrm{s.o.}}-1)}\frac{T_1}{D}\right] \right\}. \tag{3.7}$$

which can be exactly resolved and gives :

$$\frac{T_1}{D} = \frac{2(N_{\mathrm{s.o.}}-1)}{N_{\mathrm{s.o.}}^{\frac{2N_{\mathrm{s.o.}}+1}{N_{\mathrm{s.o.}}+1}}} \exp\left\{\frac{2D}{(N_{\mathrm{s.o.}}+1)J}\right\}. \tag{3.8}$$

But when $N_{\mathrm{s.o.}}$ is supposed to be large, one has:

$$\frac{T_1}{D} = \frac{2}{N_{\mathrm{s.o.}}} \mathrm{e}^{2D/N_{\mathrm{s.o.}}J} \tag{3.9}$$

and we recover the usual Kondo temperature, except the introduction of $N_{\mathrm{s.o.}}$, instead of a coefficient of 2 in the historical "$\boldsymbol{s}\bullet\boldsymbol{S}$" model (non-degenerate case). Actually, there are $N_{\mathrm{s.o.}}$ channels for changing the quantum number M at site n and these channels add independently to each other.

4 The magnetically ordered phase

For this phase, one has

$$w_n = 0 \tag{4.1}$$

therefore, equation (2.8) gives

$$H_{\mathrm{I}} = -\frac{J}{N} \sum_{n\mathbf{k}\mathbf{k}'M\sigma} \mathrm{e}^{\mathrm{i}(\mathbf{k}-\mathbf{k}')\cdot\mathbf{R}_n} \gamma_{\sigma M}(\mathbf{k}) \gamma^{*}_{\sigma M}(\mathbf{k}') \times \left(\langle \nu_{nM} \rangle c^{\dagger}_{\mathbf{k}\sigma} c_{\mathbf{k}'\sigma} + \langle c^{\dagger}_{\mathbf{k}\sigma} c_{\mathbf{k}'\sigma} \rangle \nu_{nM} - \langle c^{\dagger}_{\mathbf{k}\sigma} c_{\mathbf{k}'\sigma} \rangle \langle \nu_{nM} \rangle \right). \tag{4.2}$$

4.1 The ferromagnetic case

For this case, all the lattice sites are equivalent, they have the same value of the f-magnetic moment. At the absolute zero, the f-magnetic moment has its maximum value, and the use of equations (2.1, 2.2, 2.8, 4.2) gives for the ferromagnetic binding energy:

$$\frac{\Delta\xi_{\mathrm{F}}}{D} = \frac{N}{8} \left(\frac{J}{D}\right)^2. \tag{4.3}$$

Now, let us look for the temperature T_{F} at which the ferromagnetic order parameter vanishes. Green's function technique gives two differing results corresponding to both spin-orbit components:

$$\frac{T_{\mathrm{F}}}{D} = \begin{cases} \frac{1}{12}\frac{(N_{\mathrm{s.o.}}-1)^2}{N_{\mathrm{s.o.}}(N_{\mathrm{s.o.}}-1)} \left(\frac{J}{D}\right)^2 & \text{for } I = l - \frac{1}{2} \\ \frac{1}{12}\frac{N_{\mathrm{s.o.}}+1}{N_{\mathrm{s.o.}}} \left(\frac{J}{D}\right)^2 & \text{for } I = l + \frac{1}{2}, \end{cases} \tag{4.4}$$

For large $N_{\mathrm{s.o.}}$ one has, whatever the spin-orbit component:

$$\frac{T_{\mathrm{F}}}{D} = \frac{1}{12} \left(\frac{J}{D}\right)^2. \tag{4.5}$$

4.2 The antiferromagnetic case

For this case the lattice is divide in two sublattices A and B, and the f-magnetic moment of a site belonging to sublattice A is directly opposite to that belonging to sublattice B; then one has an extra-periodicity. At the absolute zero, the absolute value of the f-magnetic moment of each lattice site has its maximum value, and the use of equations (2.1, 2.2, 2.8, 4.2) gives the antiferromagnetic binding energy for $N_{\mathrm{s.o.}}$ large:

$$\frac{\Delta\xi_{\mathrm{AF}}}{D} = N \left\{ \frac{1}{8} \left(\frac{J}{D}\right)^2 \log\left(\frac{2D + \sqrt{4D^2 + J^2}}{-J} \right) + \frac{1}{2} \sqrt{1 + \left(\frac{J}{2D}\right)^2} - \frac{1}{2} \right\}. \tag{4.6}$$

Table 1. Numerical values of exchange coupling J_0/D (in the large $N_{\rm s.o.}$ limit), $J_{\rm T}/D$ and corresponding temperatures as well as binding energies.

$N_{\rm s.o.}$	J_0/D	$\Delta\xi_{\rm K}/D = \Delta\xi_{\rm AF}/D$	$J_{\rm T}/D$	$T_{\rm T}/D = T_{\rm AF}/D$
2(no degenaracy)			-1.6×10^{-1}	1.0×10^{-2}
$6(I = 5/2)$	-5.2×10^{-2}	1.6×10^{-3}	-5.5×10^{-2}	1.6×10^{-3}
$8(I = 7/2)$	-3.5×10^{-2}	0.8×10^{-3}	-4.0×10^{-2}	1.0×10^{-3}
14(limit)	-1.7×10^{-2}	0.2×10^{-3}	-2.0×10^{-2}	2.8×10^{-4}

Now, let us look for the temperature $T_{\rm AF}$ at which the antiferromagnetic order parameters vanish. Minimization of the Free energy functional leads, for large $N_{\rm s.o.}$, to the following equation in $T_{\rm AF}$

$$\frac{T_{\rm AF}}{D} = \frac{1}{12}\left(\frac{J}{D}\right)^2 \log\left(\frac{D}{T_{\rm AF}}\right). \tag{4.7}$$

5 On the phase boundaries

From equations (4.3-4.7), one sees that as far as a magnetically ordered phase is concerned, the antiferromagnetic case (because of the logarithmic factor) is more favoured than the ferromagnetic one, since here $n_{\rm c}$ is near one. This is somewhat related with the fact that the heavy-fermion behavior is much more often observed in anti- or nearly anti-ferromagnetic systems than in ferromagnetic ones. Also, the competition is now between the antiferromagnetic order and the local Kondo effect.

First, at the absolute zero, let us look for the value J_0 of the exchange constant J such that the system is in the Kondo phase for $|J|$ greater than $|J_0|$ and in the antiferromagnetic phase for $|J|$ less than $|J_0|$. J_0 is the solution of the equation obtained by setting the Kondo binding energy $\Delta\xi_{\rm K}$ (Eq. (3.6)) equal to the antiferromagnetic binding energy $\Delta\xi_{\rm AF}$ (Eq. (4.6)):

$$\mathrm{e}^{2D/N_{\rm s.o.}J_0}\left(1-\mathrm{e}^{2D/N_{\rm s.o.}J_0}\right)^{-1} = \frac{1}{8}\left(\frac{J_0}{D}\right)^2 \times \log\left(\frac{2D+\sqrt{4D^2+J_0^2}}{-J_0}\right) + \frac{1}{2}\sqrt{1+\left(\frac{J_0}{2D}\right)^2} - \frac{1}{2}. \tag{5.1}$$

Of course, by solving equation (5.1), we obtain the exact value of J_0, but only in the large $N_{\rm s.o.}$ limit because equation (4.6) is also deduced in that limit. Now, let us look for the value $J_{\rm T}$ of the exchange constant J such that the Kondo order parameter and the magnetic order parameters vanish together (paramagnetic phase limit). $J_{\rm T}$ is then the solution of the equation obtained by setting $T_{\rm T}$ from equation (3.9) equal to $T_{\rm AF}$ from equation (4.7)

$$\frac{2(N_{\rm s.o.}-1)}{N_{\rm s.o.}^{\frac{2N_{\rm s.o.}+1}{N_{\rm s.o.}+1}}}\exp\left\{\frac{2D}{(N_{\rm s.o.}+1)J}\right\} = \frac{T_{\rm AF}(J_{\rm T})}{D} \tag{5.2}$$

where $T_{\rm AF}$ can be computed from equation (4.7). In fact, the point of coordinates $J_{\rm T}$ and $T_{\rm T} = T_{\rm AF}(J_{\rm T})$ is a kind of triple point in the (J,T) diagram, because it belongs to the three phases (Kondo phase, antiferromagnetic phase, paramagnetic phase).

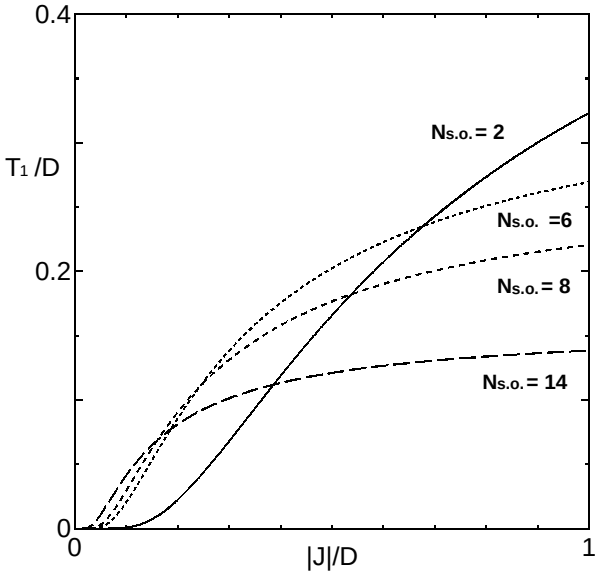

Fig. 1. Kondo temperature $|T_1|/D$ as a function of $|J|/D$ for the spin-orbit degeneracy $N_{\rm s.o.} = 2$ (solid), 6 (dot), 8 (short dash) and 14 (dash) cases. Each Kondo temperature is taken from equation (3.8).

From equations (5.1, 5.2), one can guess that $|J_{\rm T}|$ is greater than $|J_0|$, because of the $2(N_{\rm s.o.}-1)/N_{\rm s.o.}^{\frac{2N_{\rm s.o.}+1}{N_{\rm s.o.}+1}}$ factor in the left hand side of equation (5.2). Therefore, for the value of the exchange constant J between J_0 and $J_{\rm T}$, the system crosses three phases (Kondo, Antiferromagnetic, Paramagnetic) if one increases the temperature (see Tab. 1).

Figure 1 shows plots of Kondo temperature T_1/D (Eq. (3.8)) as a function of $|J|/D$ for values $N_{\rm s.o.} = 2$ (no degenaracy), 6, 8 and 14 (large limit of $N_{\rm s.o.}$). T_1 decreases with decreasing $N_{\rm s.o.}$ for small $|J|/D$ (< 0.2), but for large $|J|/D$ values such as $|J|/D \simeq 1$, T_1 decreases with increasing $N_{\rm s.o.}$. Also for large $|J|/D$ values, a rather flat behaviour is obtained for T_1 but only for large $N_{\rm s.o.}$ (degeneracy effect). It is interesting to give (Fig. 2) a magnified view of the various values of T_1 (Fig. 1) when $|J|/D$ is restricted to $|J|/D \leq 0.2$. We also plot $T_{\rm AF}$ to show the limit between Kondo, Antiferromagnetic (AF) and Paramagnetic phases: Larger degeneracies $N_{\rm s.o.}$ decreases the AF area along the $|J|/D$ axis to the benefit of the Kondo area. Actually, Figure 3 shows a scaling effect which was

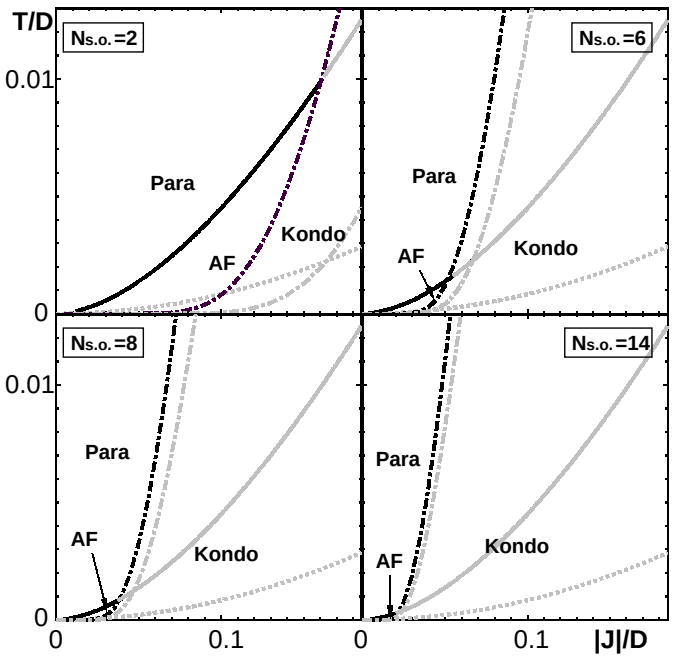

Fig. 2. Ground state phase diagram for $N_{\text{s.o.}} = 2$, 6, 8 and 14 for small values of $|J|/D$. The full line taken from equation (4.7) and the dotted line from equation (4.5) do not depend upon $N_{\text{s.o.}}$. The dashed-dotted line from equation (3.8) is compared to the pale dashed-dotted line from asymptotic equation (3.9); equation (3.9) is a good approximation of equation (3.8) for $2 < N_{\text{s.o.}} \leq 14$.

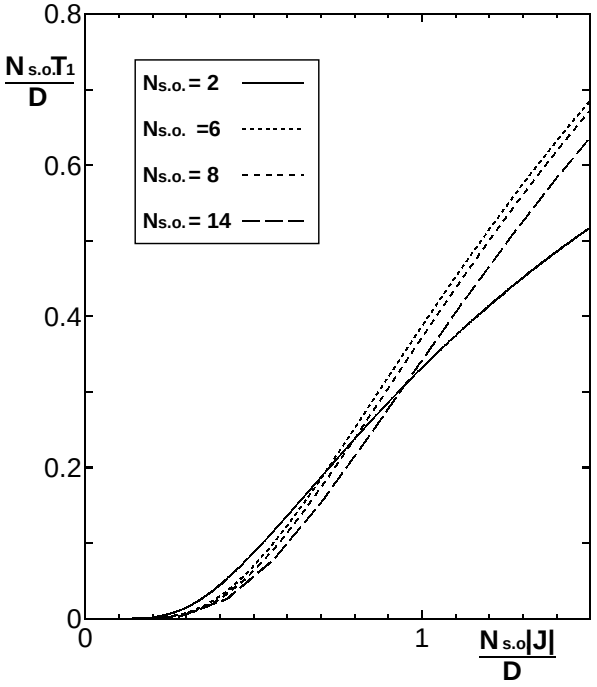

Fig. 3. The orbital degeneracy $N_{\text{s.o.}}$ dependence of the scaled Kondo temperature $T_{\text{K}} = N_{\text{s.o.}} T_1/D$ where T_1/D is given by equation (3.8). The scaled Kondo temperature T_{K} depends on the orbital degeneracy $N_{\text{s.o.}}$, but its essential feature is almost unchanged if $N_{\text{s.o.}}\ |J|/D$ is kept constant, especially in the large $N_{\text{s.o.}}$ limit.

already mentioned by Kotani *et al.* [19] within the impurity Anderson model characterized by an hybridization V. More precisely core-level photoemission spectra [19] were shown to depend on the orbital degeneracy $N_{\text{s.o.}}$ of the $4f$ state but their essential features appeared unchanged if $N_{\text{s.o.}} V^2$ was kept constant. In our case also (Fig. 3), the scaled Kondo temperature is almost the same if $N_{\text{s.o.}}|J|/D$ is kept constant: in this way we can define: $T_{\text{K}} = N_{\text{s.o.}} T_1/D$ for the scaled Kondo temperature (see also Eq. (3.9) for the asymptotic limit).

6 Concluding remarks

Within a mean field approximation the Kondo lattice system has been studied using the Coqblin-Schrieffer Hamiltonian extended to a lattice and including the degeneracy $N_{\text{s.o.}}$ of the total angular momentum I of the f states. This leads to a definition of a Kondo temperature depending on $N_{\text{s.o.}}$. The competition between magnetic order and local Kondo effect has been described when the band electron number per site $n_{\text{c}} \sim 1.0$ and for large values of $N_{\text{s.o.}}$. Our study gives some trends of the phase diagram and shows possibility of three phases (antiferromagnetic, Kondo or paramagnetic ones) depending on the values of the effective constant J and the temperature T. Also we point out important scaling effects between $N_{\text{s.o.}}$ and J and we define a new Kondo temperature almost independent of $N_{\text{s.o.}}$ if $N_{\text{s.o.}}\ |J|/D$ is kept constant.

Finally, to test our results we suggest experimental studies on pressure effect on an antiferromagnetic Kondo lattice at low temperature ($CeAl_2$ for instance), since the parameters $|J|/D$ increases with pressure. Then the question is if we can see, beyond a critical value of pressure, the system crossing and entering the Kondo phase. Of course a similar pressure effect might be obtained through alloying, for example, replacing partially the Al of $CeAl_2$ by another appropriate element. Similarly for the $CeRu_2Si_2$ system, Ce can be replaced either by La or Y element [20].

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