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1990 J. Phys.: Condens. Matter 2 10475

(http://iopscience.iop.org/0953-8984/2/51/019)

J. Phys.: Condens. Matter 2 (1990) 10475–10486. Printed in the UK

Re-entrant and spin-glass-like behaviour of the Anderson lattice

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Received 14 March 1990

Abstract. The magnetic properties of intermediate valence systems described by the periodic Anderson model are studied within a mean-field approximation for a wide range of localized 4f-level occupation n_f. It is found that the linear susceptibility χ_0 shows divergence at two critical temperatures T_{c1} and T_{c2}, indicating instability in magnetic ordering. Ferromagnetic order exists for $T_{c1} \leqslant T \leqslant T_{c2}$, indicating a re-entrant behaviour. For another range of n_f, the system exhibits spin-glass-like magnetic response with almost divergent non-linear susceptibility as the external magnetic field goes to zero, at a temperature T_0 where the linear susceptibility has a cusp-like behaviour. As n_f is scanned, it is found that T_{c1} and T_{c2} merge for a particular value of 4f-level position. This temperature T_g ($= T_{c1} = T_{c2}$) is found to increase with the increase in hybridization between the localized and band states. It is observed that the linear response for $T \geqslant T_0$ can be represented by a scaling relation.

1. Introduction

The intermediate valence state of the 4f ion in rare-earth compounds or alloys appears due to the quasi-degeneracy of two neighbouring valence states. A simplified description of the intermediate valence system is based on the observation that two types of electrons, one moving in a wide band and the other in a narrow band, are present in the system (Lawrence *et al* 1981, Czycholl 1986, Newns and Read 1987). Interaction between these two types of electrons is predominantly reflected in the magnetic behaviour of these compounds. The interplay between the valence fluctuation and the ferromagnetic order (Nolting and Matlak 1984, Eyert and Nolting 1986, Nolting and Ramakanth 1986, 1987, Gangadhar Reddy and Ramakanth 1986, 1987) and the antiferromagnetic order (Leder and Muhlschlegel 1978, Stratkotter and Nolting 1987, Bulk and Nolting 1988) has been examined on the basis of the periodic Anderson lattice in a mean-field approximation. The existence and nature of the magnetic order depend on the occupation n_f of the localized 4f level and the effective Coulomb interaction among the 4f electrons (Leder and Muhlschlegel 1978, Gangadhar Reddy and Ramakanth 1986, 1987). When the valence of the rare-earth ion is nearly integral ($n_f \simeq 1$), magnetic order exists below the transition temperature, which decreases as the 4f-level position moves towards the centre of the wide band. For a certain range of parameter values, as the temperature is increased, a re-entrant magnetic transition has been noted (Leder and Muhlschlegel 1978, Entel *et al* 1978). When the system is driven into a state of lower occupation of the

4f level, the susceptibility shows a sharp peak at a finite temperature. A closer study of the magnetic behaviour of this situation reveals spin-glass-like behaviour (Gangadhar Reddy *et al* 1989). Recent experimental results on $CeCu_{6.5}Al_{0.5}$ (Rauschschwalbe *et al* 1985) and $CePd_3B_{0.3}$ (Dhar *et al* 1989) also suggest the existence of a spin-glass-like state in the mixed valence phase. Hybridization between the localized and band states introduces an element of randomness in the magnitude of the local moment and the effective interaction between the moments at different sites. In contrast the spin-glass in the localized moment system has randomness only in the interaction between the moments at neighbouring sites. In this paper, we present a detailed analysis of the magnetic phase transition boundary as the occupation of the 4f level changes, based on the periodic Anderson model in a mean-field approximation. The evolution of the re-entrant phase and spin-glass-like phase is studied following a detailed calculation of the magnetization and the linear and non-linear magnetic responses.

2. Theory

The Hamiltonian of the periodic Anderson model can be written as the sum of three terms,

$$H = H_s + H_f + H_v. \tag{2.1}$$

The first term

$$H_s = \sum_{i,j,\sigma} (T_{ij} - Z_\sigma h) d_{i\sigma}^+ d_{j\sigma} \tag{2.2}$$

describes the uncorrelated conduction electron states. Here $d_{i\sigma}^+$ and $d_{i\sigma}$ are, respectively, the creation and annihilation operators of an electron in the conduction band with spin σ at site $\boldsymbol{R}_i$; T_{ij} are the usual hopping integrals, related to the Bloch energies $\varepsilon(\boldsymbol{k})$ by

$$T_{ij} = \frac{1}{N} \sum_{\boldsymbol{k}} \varepsilon(\boldsymbol{k}) \, e^{i\boldsymbol{k}\cdot(\boldsymbol{R}_i - \boldsymbol{R}_j)} \tag{2.3}$$

$h = Sg\mu_B B$ is the magnetic energy, with S the electron spin, g the Landé g-factor, μ_B the Bohr magneton and B the external magnetic field; and Z_σ is a sign factor,

$$Z_\uparrow = +1 \qquad Z_\downarrow = -1. \tag{2.4}$$

The Hamiltonian H_f describes a periodic array of strongly correlated 4f electrons,

$$H_f = \sum_{i,\sigma} [(E_f - Z_\sigma h) n_{fi\sigma} + \tfrac{1}{2} U n_{fi\sigma} n_{fi-\sigma}]. \tag{2.5}$$

Here $n_{fi\sigma} = f_{i\sigma}^+ f_{i\sigma}$ is the number operator for the 4f electrons, with $f_{i\sigma}^+$ and $f_{i\sigma}$ the corresponding creation and annihilation operators, respectively; the 4f-level position E_f is measured relative to the centre of gravity of the conduction band; and U is the intrasite Coulomb repulsion.

The final term H_v describes the hybridization,

$$H_v = \sum_{i,\boldsymbol{k},\sigma} \frac{V_{\boldsymbol{k}i}}{N^{1/2}} (f_{i\sigma}^+ d_{\boldsymbol{k}\sigma} \, e^{-i\boldsymbol{k}\cdot\boldsymbol{R}_i} + d_{\boldsymbol{k}\sigma}^+ f_{i\sigma} \, e^{+i\boldsymbol{k}\cdot\boldsymbol{R}_i}). \tag{2.6}$$

The hybridization allows the hopping of an electron from the localized 4f level to the conduction band and vice versa; the hybridization strength V_{ki} is taken to be independent of the lattice site.

The Hamiltonian (2.1) cannot be solved exactly. In order to solve the problem, in the literature, several approximate methods have been proposed (for reviews see Czycholl 1986, Newns and Read 1987). In the present investigation we use the mean-field approximation (Leder and Muhlschlegel 1978). In this approximation the conduction band occupation is given by

$$n_{\mathrm{d}\sigma} = \frac{1}{N}\sum_k \sum_{t=1}^{2}\left(1 - \frac{\partial E_\sigma(\boldsymbol{k}, t)}{\partial E_{\mathrm{f}\sigma}}\right) f(E_\sigma(\boldsymbol{k}, t)) \tag{2.7}$$

and the 4f-level occupation by

$$n_{\mathrm{f}\sigma} = \frac{1}{N}\sum_k \sum_{t=1}^{2} \frac{\partial E_\sigma(\boldsymbol{k}, t)}{\partial E_{\mathrm{f}\sigma}} f(E_\sigma(\boldsymbol{k}, t)). \tag{2.8}$$

Here $E_\sigma(\boldsymbol{k}, t)$ are the single-particle energies, which are the poles of the Green functions and are given by

$$E_\sigma(\boldsymbol{k}, t) = \tfrac{1}{2}\{E_{\mathrm{f}\sigma} + \varepsilon_\sigma(\boldsymbol{k}) \pm [(E_{\mathrm{f}\sigma} - \varepsilon_\sigma(\boldsymbol{k}))^2 + 4V^2]^{1/2}\} \tag{2.9}$$

with $E_{f\sigma} = E_f + Un_{f-\sigma}$ and $f(E) = \{\exp[(E-\mu)\beta] + 1\}^{-1}$ is the Fermi function with $\beta = 1/K_{\mathrm{B}}T$. The chemical potential μ is determined from conservation of the total number of electrons per lattice site,

$$n = \sum Z_\sigma(n_{\mathrm{f}\sigma} + n_{\mathrm{d}\sigma}). \tag{2.10}$$

Equations (2.7), (2.9) and (2.10) constitute the equations for the determination of μ and $n_{\mathrm{f}\sigma}$ self-consistently as functions of the model parameters. For numerical simplification we have used a constant density of states for the conduction band,

$$\rho(x) = \begin{cases} 1/W & \text{if } -W/2 \leqslant x \leqslant W/2 \\ 0 & \text{otherwise} \end{cases} \tag{2.11}$$

where W is the bandwidth of the d state.

3. Magnetization

The self-consistent equations (2.7), (2.9) and (2.10) have been solved numerically for the total number of electrons per lattice site $n = 1$ with and without the magnetic field as a function of temperature for various sets of model parameters, namely 4f-level position E_{f} and hybridization strength V. The parameters of the model are normalized in terms of bandwidth W. Throughout the calculation we have taken $U = 0.25$. The magnetization is defined as

$$M(h, T) = \mu_{\mathrm{B}} m(h, T) \tag{3.1}$$

where

$$m(h, T) = m_{\mathrm{d}}(h, T) + m_{\mathrm{f}}(h, T). \tag{3.2}$$

Here m_{f} and m_{d} are dimensionless magnetizations for the localized and conduction electrons, which we define as

$$m_{\mathrm{f}} = \sum_\sigma Z_\sigma n_{\mathrm{f}\sigma} \qquad m_{\mathrm{d}} = \sum_\sigma Z_\sigma n_{\mathrm{d}\sigma}. \tag{3.3}$$

In the absence of an external magnetic field, the magnetization due to the conduction

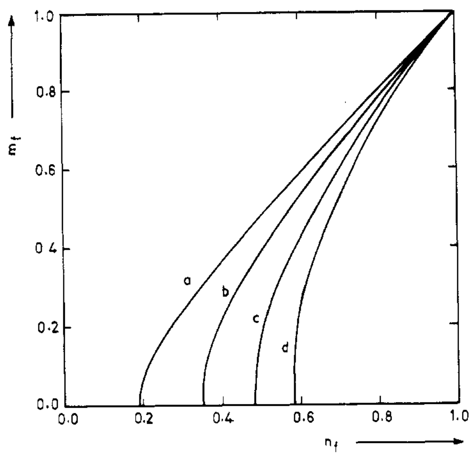

Figure 1. The dependence of the localized spontaneous magnetization m_f as a function of the 4f-level occupation n_f for various values of V: (*a*) $V = 0.05$, (*b*) 0.1, (*c*) 0.15 and (*d*) 0.25.

electrons is always found to be much smaller than that due to the localized electrons for any choice of model parameters and so $m(0, T) \simeq m_f(0, T)$. In the following, the results are given in terms of $m_f(T)$, which depend sensitively on the model parameters and the temperature. At $T = 0$ the magnetic moment varies linearly with n_f when n_f is large (figure 1). But near a critical value of n_f, which depends on the values of the hybridization and the Coulomb interaction, the moment decreases rapidly; and below the critical value of n_f, no spontaneous magnetization exists. A system with n_f less than and close to the critical value of n_f exhibits re-entrant spin-glass-like behaviour as discussed in the next section. The thermal behaviour of the spontaneous magnetization m_f per unit 4f-level occupancy is shown in figure 2 for different positions of E_f and for the same hybridization constant $V = 0.1$. For n_f close to 1, the ferromagnetic state at $T = 0$ is close to saturation ($m_f = n_f$) and goes to zero monotonically as T approaches T_c, the transition temperature. As E_f increases towards the d-band centre, n_f decreases and the state at $T = 0$ is an unsaturated one with $m_f < n_f$ but still exhibiting normal T dependence. As n_f decreases further, a different thermal behaviour of m_f is found. For $E_f = -0.3$ the magnetization goes up as T is increased from $T = 0$ and passes through a broad maximum before vanishing at $T = T_c$. Such behaviour of m_f persists with further increase of E_f. The spontaneous magnetization at $T = 0$ falls off faster and ultimately vanishes at a particular value of E_f. For E_f greater than this value, two transitions are observed, marking a re-entrant type of behaviour. The magnetization that appears at the higher transition temperature disappears at the lower transition temperature. The re-entrant phenomenon occurs within a small region of E_f (or n_f). As the 4f-level occupation is depleted further, no ferromagnetic state appears at any temperature. The thermal variation of n_f for the above values of E_f is given in figure 3. We note the different nature of the thermal dependence of n_f. For large values of n_f, temperature increase causes slow depopulation. On the other hand, for smaller value of n_f, average occupancy goes up with increasing temperature. For a range of values of n_f where the re-entrant magnetic transition is found, two kinks appear in n_f versus T behaviour. At higher temperatures or still smaller values of n_f, almost a linear increase of 4f-level occupancy is observed.

4. Linear and non-linear magnetic responses

The static magnetic susceptibility in the state without magnetic order is given by

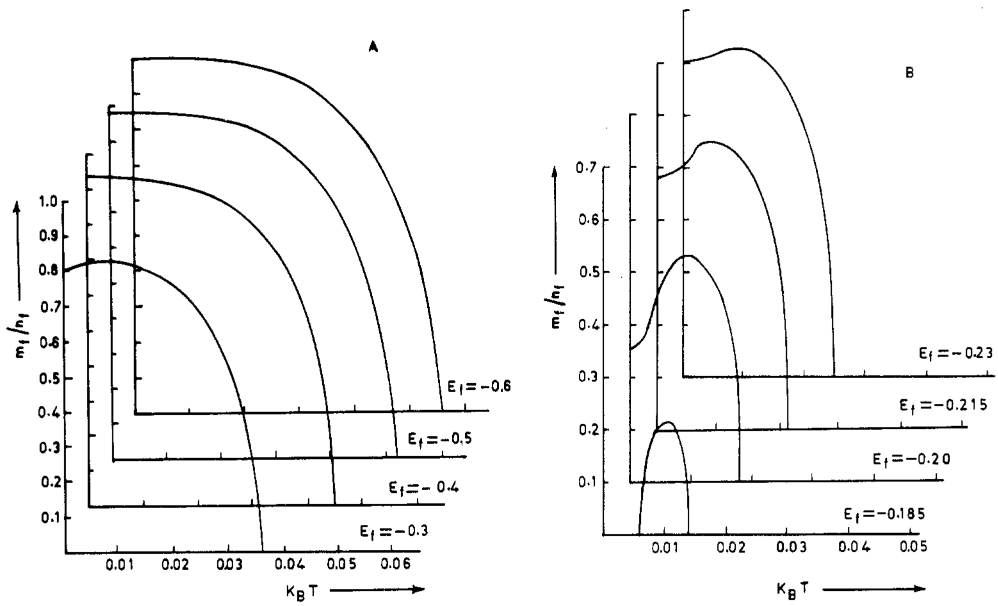

Figure 2. The normalized 4f magnetization as a function of temperature for various values of E_f: (A) in the region where m_f has normal T dependence and (B) in the region where m_f has peak-like behaviour. $V = 0.1$.

$$\chi_0/\mu_B^2 = 2[I_1 + UI_2^2/(1 - UI_3)] \tag{4.1}$$

where

$$I_1 = \frac{\beta}{4N}\sum_k \sum_{t=1}^{2} \cosh^{-2}\{\tfrac{1}{2}\beta[E(\boldsymbol{k}, t) - \mu]\} \tag{4.2}$$

$$I_2 = \frac{\beta}{4N}\sum_k \sum_{t=1}^{2} \frac{\partial E(\boldsymbol{k}, t)}{\partial \bar{E}_f} \cosh^{-2}\{\tfrac{1}{2}\beta[E(\boldsymbol{k}, t) - \mu]\} \tag{4.3}$$

$$I_3 = \frac{1}{N}\sum_k \sum_{t=1}^{2} \left[\frac{\beta}{4}\left(\frac{\partial E(\boldsymbol{k}, t)}{\partial \bar{E}_f}\right)^2 \cosh^{-2}\{\tfrac{1}{2}\beta[E(\boldsymbol{k}, t) - \mu]\} - \frac{\partial^2 E(\boldsymbol{k}, t)}{\partial \bar{E}_f^2} f(E(\boldsymbol{k}, t))\right] \tag{4.4}$$

with

$$E(\boldsymbol{k}, t) = \tfrac{1}{2}\{\bar{E}_f + \varepsilon(\boldsymbol{k}) \pm [(\bar{E}_f - \varepsilon(\boldsymbol{k}))^2 + 4V^2]^{1/2}\} \tag{4.5}$$

$$\bar{E}_f = E_f + Un_f/2. \tag{4.6}$$

The ferromagnetic phase sets in when the Stoner criterion $UI_3 = 1$ is satisfied. The susceptibility diverges at T_c, which is determined by the equation

$$UI_3(T_c) = 1. \tag{4.7}$$

For large occupation of the 4f level, there is only one value for T_c, but within a range of

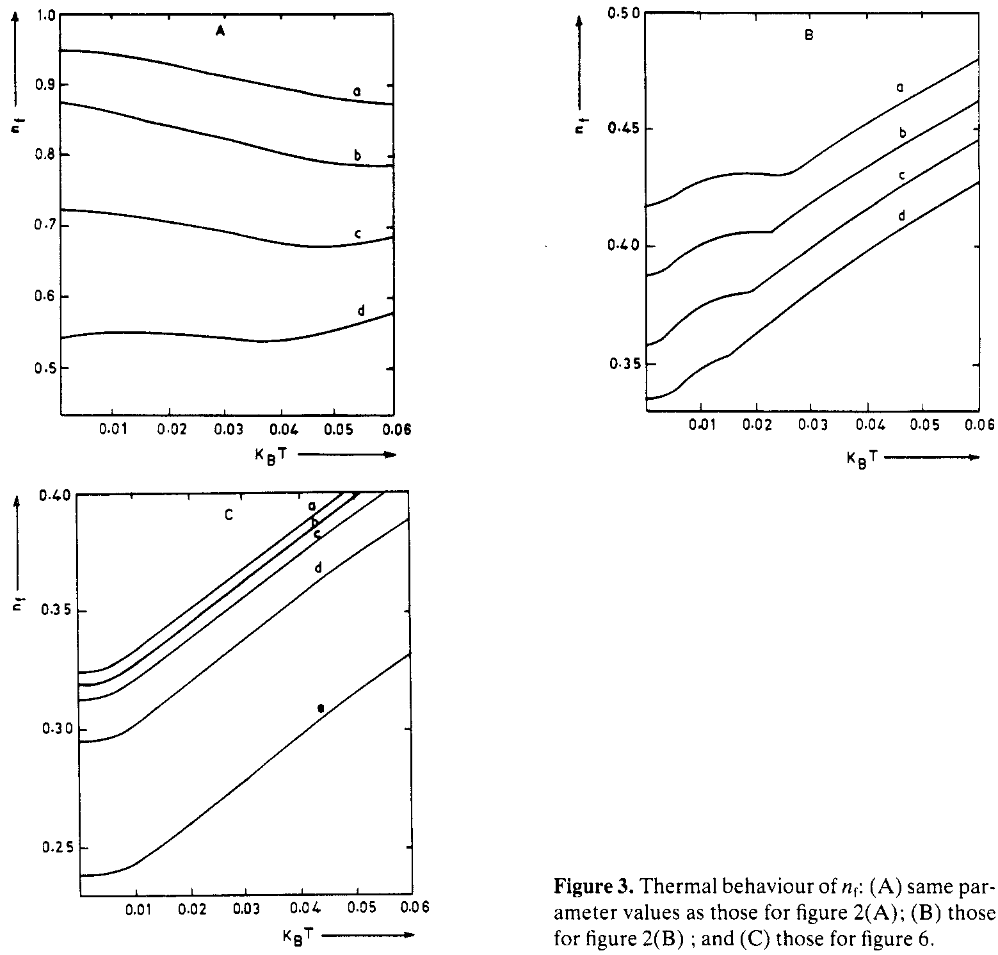

Figure 3. Thermal behaviour of n_f: (A) same parameter values as those for figure 2(A); (B) those for figure 2(B) ; and (C) those for figure 6.

smaller occupation, two solutions are found (figure 4). This is consistent with the magnetization result (figure 2(B)). The locus of T_c for the region of E_f where re-entrant behaviour is found is shown in figure 5. As E_f increases, at a particular value of E_f, the two T_c values merge and this temperature $T_{c1} = T_{c2}$ is defined as T_g. The variation of T_g with hybridization is shown in the inset of figure 5. For a given V, when E_f increases beyond the critical value at which the two T_c values are the same, the susceptibility exhibits a cusp-like behaviour at a temperature T_0 (figure 6). As n_f decreases, the height of the peak decreases whereas T_0 goes up. For very small values of n_f the susceptibility is nearly constant and Pauli-like at low temperature. In all cases, at high temperature, $T > T_0$, the susceptibility has a Curie–Weiss-like decrease. An interesting correlation between T_0 and the inverse of the susceptibility at $T = 0$ is displayed in the inset of figure 6. This bears a striking resemblance to the experimental results (Klaasse *et al* 1981) discussed by Newns and Read (1987).

In order to determine the magnetic state of a system that shows a maximum in χ_0 at a temperature T_0, the magnetization and the non-linear response are evaluated. The magnetization at a finite field with $n_f = 0.3303$ at $T = 0$ passes through a maximum as T

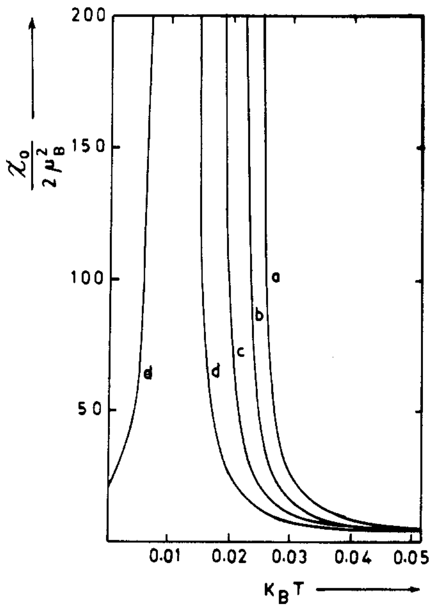

Figure 4. The linear susceptibility χ_0 as a function of temperature for different values of E_f: (a) $E_f = -0.23$, (b) -0.215, (c) -0.20 and (d) -0.185.

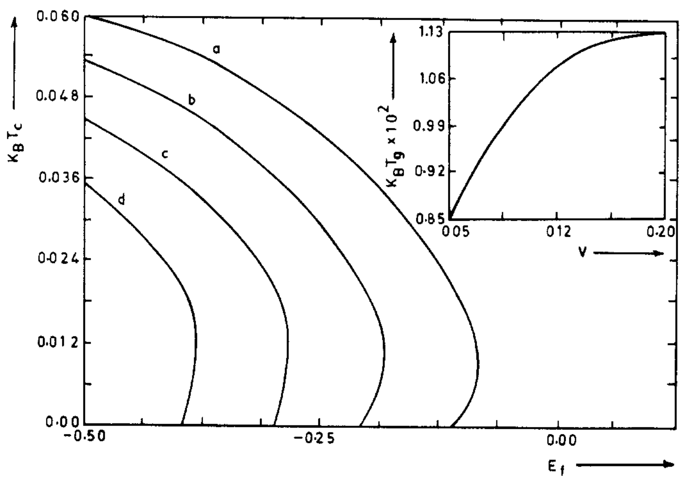

Figure 5. The Curie temperature T_c as a function of 4f level E_f for different values of the hybridization strength V: (a) $V = 0.05$, (b) 0.1, (c) 0.15 and (d) 0.2. Inset: T_g as a function of V.

is increased (figure 7). The maximum shifts to higher values of temperature as the magnetic field is increased. In the limit of small field,

$$M(B) = \chi_0 B + \chi_2 B^3 \tag{4.8}$$

where χ_2 is the non-linear susceptibility. Defining

$$\chi_2/\mu_B^4 = [m(h)/h - \chi_0/\mu_B^2]/h^2 \tag{4.9}$$

the numerical results for χ_2 for different fields are displayed in figure 8 along with χ_0.

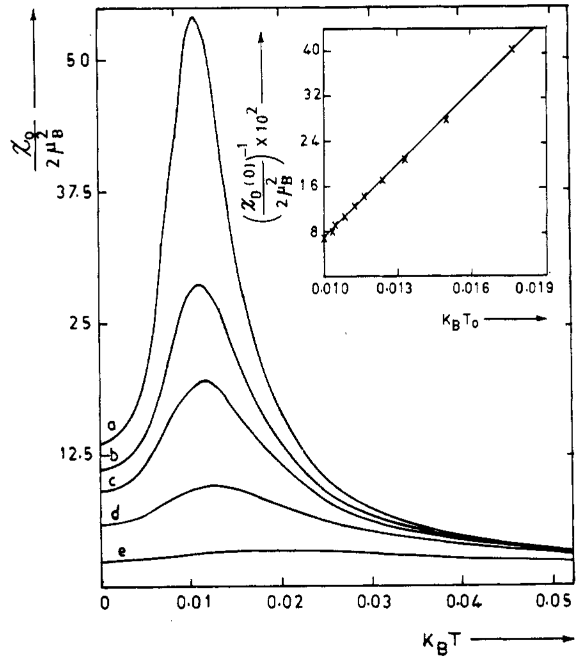

Figure 6. The temperature dependence of the linear susceptibility χ_0 for different values of E_f: (a) $E_f = -0.175$, (b) -0.17, (c) -0.165, (d) -0.15 and (e) -0.1. $V = 0.1$. Inset: the inverse susceptibility at $T = 0$ versus the peak temperature.

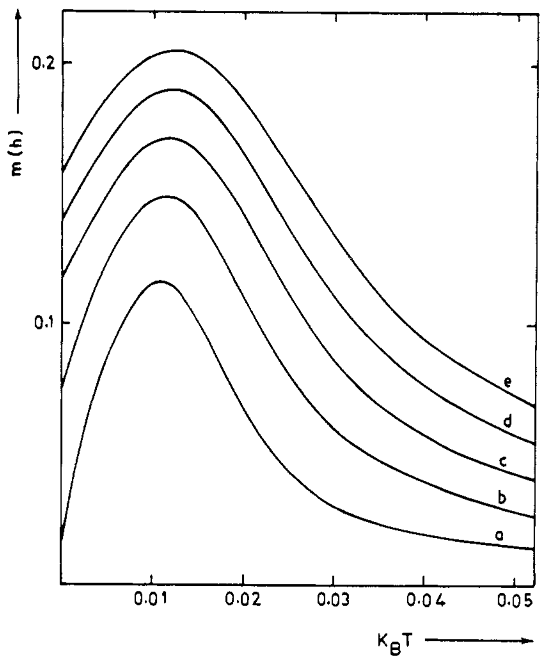

Figure 7. The temperature variation of the field-dependent magnetization for different values of the magnetic energy h: (a) $h = 0.001$, (b) 0.002, (c) 0.003, (d) 0.004 and (e) 0.005. $E_f = -0.18$ and $V = 0.1$.

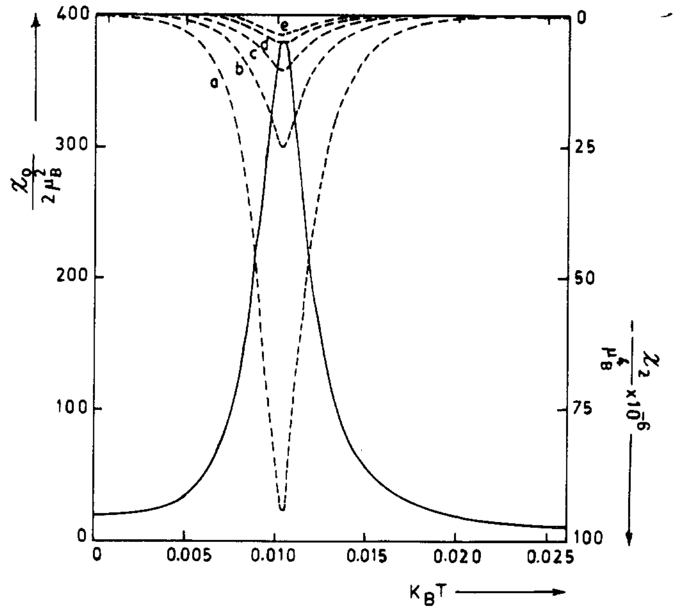

Figure 8. The non-linear susceptibility χ_2 (broken curves) and the linear susceptibility χ_0 (full curves) as a function of temperature for different values of the magnetic field strength h: (a) $h = 0.001$, (b) 0.002, (c) 0.003, (d) 0.004 and (e) 0.005. $E_f = -0.18$ and $V = 0.1$

This shows that $\chi_2(h)$ has a sharp dip at T_0 where χ_0 has a maximum, and its magnitude increases very steeply as the field goes down. In the limit of $h \to 0$, χ_2 tends to diverge at $T = T_0$. This is more evident in the plot of the reduced non-linear response

$$R = 1 - \mu_B^2 m(h)/h\chi_0 \tag{4.10}$$

as a function of the square of the magnetic field at temperature $T \simeq T_0$. For $T \gtrless T_0$, the slope as $h \to 0$ is finite, but for $T = T_0$ it grows to an infinitely large value (figure 9). The above magnetic behaviour closely mimics that of the spin-glass phase. As E_f approaches the centre of the band, both linear and non-linear responses of the system are diminished.

In conventional spin-glass systems the magnetic responses at $T \geqslant T_0$ is found to be represented by the formula (Muzumder and Bhagat 1987)

$$\chi_0(T) = \chi_0(T_0)t^{-\gamma}/[\chi_0(T_0)A + t^{-\gamma}] \tag{4.11}$$

where $t = (T - T_0)/T_0$, γ is a critical exponent and A is a constant. In order to examine whether such a scaling relation exists here also, in figure 10(A) we have plotted $\ln(x)$ as a function of $\ln(t)$, where $x = 1/\chi_0(T) - 1/\chi_0(T_0)$. The above relation is found to be valid for a small region of temperature, similar to the behaviour of the susceptibility in traditional spin-glass systems (Muzumder and Bhagat 1987). When the above results are scaled in terms of the non-linear scaling variable $t' = (T - T_0)/T$ an improved formula

$$\chi_0(T) = \chi_0(T_0)(1 - t')t'^{-\gamma}/[\chi_0(T_0)A + (1 - t')t'^{-\gamma}] \tag{4.12}$$

fits the results over a wider range of temperature (figure 10(B)). This is similar to the recent result on an Ising spin-glass where the relation (4.12) is found to be obeyed for the entire range of temperature (Basu and Ghatak 1990).

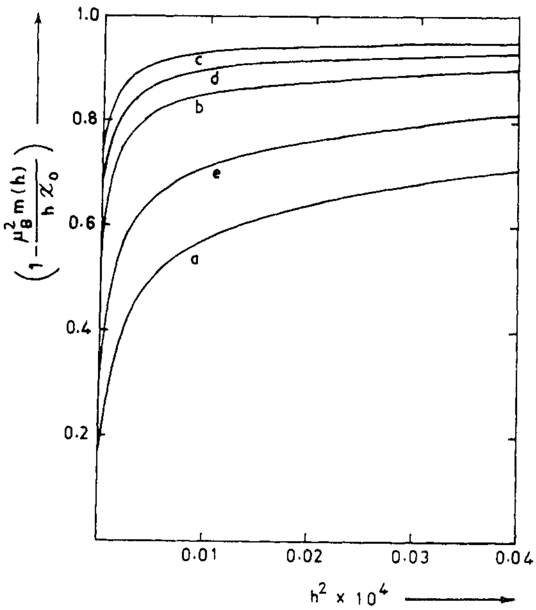

Figure 9. Non-linear response $R = 1 - \mu_B^2 m(h)/h\chi_0$ as a function of the square of the magnetic field strength at different temperatures: (a) $K_B T = 6.45 \times 10^{-3}$, (b) 8.6×10^{-3}, (c) 9.89×10^{-3}, (d) 11.18×10^{-3} and (e) 11.33×10^{-3}. $E_f = -0.18$ and $V = 0.1$. Curve (c) is for $T = T_0$.

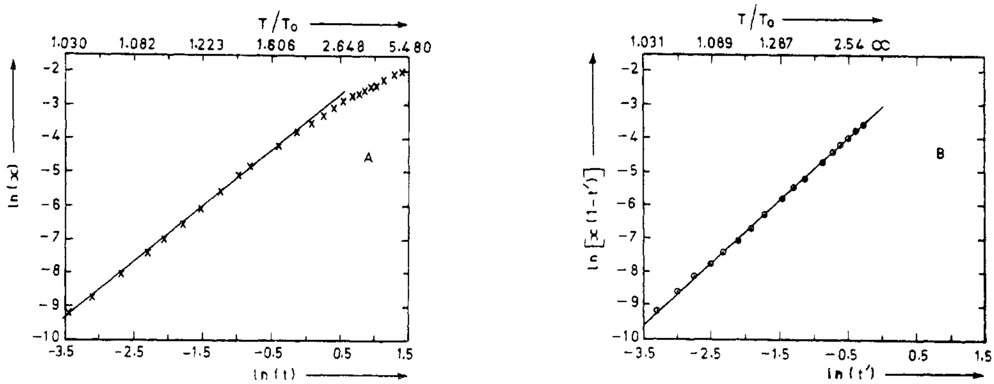

Figure 10. The scaling behaviour of the linear response: (A) for equation (4.11) and (B) for equation (4.12). The value of T/T_0 is marked on the top of each figure.

5. Discussion of the results

The location of the 4f level with respect to that of the conduction band is the crucial factor in determining the properties of the model system (Leder and Muhlschlegel 1978, Gangadhar Reddy and Ramakanth 1986, 1987, Nolting and Ramakanth 1987). When the 4f level is well below the bottom of the conduction band, the effect of the fluctuations is very small, $n_f \simeq 1$, so that the system is integral valent with the magnetic configuration. As the 4f level approaches the conduction band, the fluctuations in valence increase and the system becomes mixed valent, so that n_f deviates substantially from the integral value. When the 4f level is pushed further into the band, 4f-level occupation is depleted to zero, so that the system is again in the integral valence phase but with the non-magnetic configuration. Naturally, the spontaneous magnetization should undergo qualitative changes as the 4f level is swept. Initially, in the integral or near-integral valence phase, m_f decreases monotonically with increasing T, going to zero at a critical temperature T_c (figure 2). The anomalous thermal behaviour of the magnetization for smaller occupation of the 4f state can be correlated with the thermal behaviour of n_f (figure 3). When $m_f(T)$ decreases monotonically (figure 2(A)), $n_f(T)$ also decreases monotonically with increasing T (figure 3(A)). When $m_f(T)$ displays a peak-like behaviour (figure 2(B)), $n_f(T)$ initially increases with increasing T and then decreases, so that it also has a peak-like behaviour (figure 3(B)) though not as prominently as $m_f(T)$. The reduction in n_f leads to a quenching of the local moment so that the spontaneous magnetization will decrease. On the other hand, when the 4f level is near or inside the conduction band at low T, $n_f(T)$ increases with increasing T and there are two competing mechanisms: (i) 4f-level occupation trying to stabilize magnetic order and (ii) thermal agitation trying to destabilize it. As T increases, n_f increases, and so for low T the first mechanism prevails and therefore $m_f(T)$ increases. For higher T, the second one dominates and therefore $m_f(T)$ decreases, resulting in a peak-like structure. The result that as the 4f level is further moved up, the peak in m_f sharpens and shifts to higher temperature, can be understood on similar lines. As E_f moves up, the 4f-level occupation is depleted so it requires higher temperature for the stability of the local moment. Again, if T is too large, the magnetization is destroyed due to the thermal agitation. Therefore, the region in temperature where $m_f(T)$ is finite becomes progressively smaller with increasing E_f, resulting in a sharper peak. Further, the increase in the rate of change of n_f with T (figure 3) is reflected in the shfiting of the peak in m_f to higher temperature.

When the 4f level is sufficiently deep inside the band, its occupation at $T = 0$ is too small to support magnetic order. With increasing T, n_f increases, and at a critical temperature, the magnetization becomes non-zero. On further increase of T, m_f displays peak-like behaviour, vanishing at a higher T, another critical temperature (the curve with $E_f = -0.185$ in figure 2(B)). This re-entrant behaviour is also confirmed by the existence of two singularities in the paramagnetic susceptibility (figure 4). The occurrence of the re-entrant behaviour in this model has been mentioned earlier (Leder and Muhlschlegel 1978) and it also appears in the presence of electron–phonon interaction (Entel *et al* 1978). As the hybridization strength V increases, the lower transition temperature T_{c1} increases and the higher one T_{c2} decreases, so that the region in temperature where the spontaneous magnetization is non-zero decreases. Since the increase in V enhances the valence fluctuations, the situation is unfavourable to magnetic ordering though n_f is large enough (Gangadhar Reddy and Ramakanth 1986). To stabilize the magnetic ordering, n_f has to be slightly increased and this is accomplished by increasing T. Therefore, T_{c1} increases. A large V and higher T both try to destabilize magnetism and therefore T_{c2} decreases with increasing V for a given E_f (figure 5). With increasing V, the value of n_f at which the magnetism exists increases (figure 1). As a result, the temperature T_g increases with increasing V (inset of figure 5). The peak temperature in χ_0 and its magnitude at $T = 0$ depend on the value of E_f, that is, n_f (figure 7). Further, these two have an interesting scaling relation (inset of figure 6), which is also recognized in experiments (Klaasse *et al* 1981).

Since the rare-earth ions fluctuate between magnetic and non-magnetic configurations, when E_f is inside the conduction band, the valence fluctuations are large and consequently the local moment fluctuates. These incoherent fluctuations of the magnetic moment introduce the element of randomness. The magnetic behaviour of the system bears a close resemblance in many respects to that of conventional spin-glass systems. The peak-like behaviour of the field-dependent magnetization as a function of temperature (figure 7), the almost divergent non-linear magnetic susceptibility (figure 8) (Suski 1977, Chalupa 1977, Ghatak 1986), a very large non-linear response R near T_0 and its magnetic field dependence (figure 9) (Omari *et al* 1983, Ghatak 1986) and the scaling behaviour of linear response (figure 10) (Muzumder and Bhagat 1987) are very characteristic of a spin-glass system with randomness in the interaction between moments. The results as derived here suggest that the fluctuation of moment due to valence fluctuations can give rise to spin-glass-like behaviour.

6. Conclusions

Based on the periodic Anderson lattice, with mean-field approximation, the magnetic state and its behaviour are found to depend strongly on the occupancy of the 4f level. Depending upon the occupancy of the 4f level the long-range magnetic order, the re-entrant magnetic phase transition, the spin-glass-like magnetic state and the paramagnetic (Pauli-like at low temperature) state are found to follow. Improved approximation of the intra-atomic Coulomb interaction term in the not-too-large U limit is not expected to change the scenario of the above magnetic phase too drastically.

Acknowledgments

One of the authors (AR) acknowledges CSIR (Government of India) for financial support in the form of a research scheme. GGR thanks ICTP, Trieste, Italy, where part of the work was done, for their hospitality.

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