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The electrical resistivity of the two-conduction-band heavy-fermion alloy system

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Abstract. The electrical resistivity of the heavy-fermion alloys is calculated through a two-conduction-band slave boson model using the self-consistent coherent potential approximation method. The results indicate that for very low temperatures the resistivity of the alloy system follows the $\rho_0 + AT^2$ law and the coefficient A changes from negative values to positive values as the alloy concentration increases. The occurrence of the resistivity maximum at a finite temperature is also obtained on increasing the concentration of heavy-fermion alloys.

1. Introduction

Much progress has been made in experimental and theoretical research on the heavy-fermion (HF) or the Kondo lattice (KL) system [1]. The HF system, which covers a wide range of inter-metallic rare-earth or actinide compounds, is characterised by a very large effective mass of the electrons near the Fermi level. Since the magnetic ions in this system usually have a 4f or 5f shell with an extremely small radius, very strong intra-site correlation and relatively weak hybridisation with the conduction band, theoretically the Anderson Hamiltonian with strong intra-site correlation is often used. Because the problem of dilute magnetic ions in a metal (Kondo impurity) has been extensively studied in the past [2], the research interest of the theoretical and experimental workers was mostly focused on the KL system [3–8]. In recent years, attention has been devoted to the Kondo alloy system and the disorder effect on the HF system [9–17]. Experimentally, the crossover from single-impurity behaviour to the KL in the HF alloy $Ce_xLa_{1-x}Cu_2Si_2$ has been investigated [9]. They have found that the low-temperature maximum of the electrical resistivity was observed at $T = T_{max}$, where T_{max} increases with the increasing magnetic ion concentration x. More accurate and systematic results have been given in [10, 11]. Theoretically, the alloy analogue approximation of the Anderson lattice has been treated previously within the coherent-potential approximation (CPA) approach [18, 19] so that the mixed valence in the system can be discussed. A CPA calculation of the electrical resistivity of the HF alloys has been given in [14], where the Anderson alloy model with small f-band dispersion was used. However, in the calculation in [14], a drastic variation in the Kondo temperature T_K with the concentration of the alloy was assumed. This is not in agreement with the experimental

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results in [10]. In a previous paper [15], we have investigated the mean-field electron density of states (DOS) of the HF alloy system through a two-conduction-band slave boson alloy Hamiltonian. We have found that the electron DOS changes gradually from that of the concentrated limit (KL) to the dilute limit (Kondo impurity). We have also demonstrated that, in the mean-field approximation, the Kondo temperature of the system does not vary with the concentration. This result is in agreement with the experiment in [10]. In the present work, we intend to discuss the resistivity of the system through the slave boson alloy model. We organise the paper as follows: § 2 is a brief review of the mean-field CPA formalism of the two-conduction-band slave boson alloy model. In § 3, the mean-field resistivity of the system is formulated through the Green functions in the CPA effective medium. The numerical result of the residual resistivity for different concentrations is also given in § 3. We present the numerical results of the mean-field finite-temperature resistivity in § 4. In § 5, by adding a Gaussian fluctuation contribution of the resistivity to the mean-field result according to the Matthiesson rule, the temperature T_{max} of the resistivity maximum for various concentrations is estimated. Finally, § 6 contains some concluding remarks.

2. Model and CPA formalism of the alloy system

The alloy system that we shall investigate is constituted of two types of rare-earth or actinide atoms A and B. In this system, some of the A (Ce- or U-like) atoms are substituted by B (La- or Lu-like) atoms which are neighbours of the A atoms in the periodic system and have no f electrons. Examples of the alloys are $Ce_xLa_{1-x}Cu_2Si_2$, $Ce_xLa_{1-x}Pb_3$ and $Ce_xLa_{1-x}Cu_6$. The system may change from a pure lattice of Ce- or U-like atoms (i.e. KL) to a dilute alloy when the concentration x (percentage of Ce- or U-like atom) varies from unity to zero. The Anderson Hamiltonian and its lattice form with extremely strong f repulsion ($U = \infty$) are usually adopted as a starting point for the discussion of the HF problem. A powerful method for handling the low-temperature properties of the system is the slave (or auxiliary) boson technique [20–22]. In order to study the crossover from the single impurity to the concentrated regime, the many-impurity Anderson model must be introduced. In accordance with [20], the slave boson Hamiltonian of the many-impurity Kondo system can be written as

$$H = \sum_{i,k,\varepsilon} \varepsilon_{ik} c^+_{ik\varepsilon} c_{ik\sigma} + \sum_{i,l,\varepsilon} V\xi_l (c^+_{il\sigma} f_{l\sigma} b_l + \text{HC}) + \sum_{i,\varepsilon} [E_f \xi_l + E_L(1-\xi_l)] f^+_{l\sigma} f_{l\sigma} \tag{1}$$

with the constraint for each A site given by

$$\left\langle \sum_\sigma f^+_{l\sigma} f_{l\sigma} + b^+_l b_l \right\rangle = 1 \qquad l \in A \tag{2}$$

where i (= 1, 2) is the index for the conduction bands, $\varepsilon_{1,2} = \varepsilon_k \pm \varepsilon_0$ and $N(\varepsilon) = 1/2D$ when $|\varepsilon| < D$ (D is the half-width of the conduction band). Here we assume that there are two conduction bands with the same dispersion ε_k, and $2\varepsilon_0$ is the separation between two bands. For $\varepsilon_0 < D$, the metallic behaviour of the system can be ensured even in the pure lattice case [23–25]. ξ_l is the random variable ($\xi_l = 1$ for $l \in A$; $\xi_l = 0$ for $l \in B$). b^+_l and b_l are slave boson operators introduced on each A site. σ is the spin index. Here we take $j = s = \frac{1}{2}$ for brevity. E_f is the energy level of the f electrons on the A sites and E_L is the energy level of the f electrons on the B sites. In order to assure no f-electron occupation on the B sites, we should take $E_L \to \infty$. In the mean-field approximation, the

operator b is replaced by a c number and the Lagrange multiplier λ_l is introduced to satisfy the constraint. The parameters b_l and λ_l are determined by the equations

$$\begin{aligned} \lambda_l b_l &= -V \sum_\sigma \langle c_{l\sigma}^+ f_{l\sigma} \rangle \\ \langle n_f^l \rangle &= 1 - |b_l|^2 \qquad l \in A. \end{aligned} \tag{3}$$

For simplicity, we assume that

$$b_l = b \qquad \lambda_l = \lambda \qquad l \in A. \tag{4}$$

So the mean-field Hamiltonian of the HF alloy system can be written as

$$H = H_0 + H_{\text{dis}}$$

$$H_0 = \sum_{k,\sigma} [c_{1k\sigma}^+ c_{2k\sigma}^+ f_{k\sigma}^+] \mathbf{E}(k) \begin{bmatrix} c_{1k\sigma} \\ c_{2k\sigma} \\ f_{k\sigma} \end{bmatrix}$$

$$H_{\text{dis}} = \sum_{l,\sigma} [c_{1l\sigma}^+ c_{2l\sigma}^+ f_{l\sigma}^+] [\xi_l \mathbf{V}_{\text{A}} + (1 - \xi_l)\mathbf{V}_{\text{B}}] \begin{bmatrix} c_{1l\sigma} \\ c_{2l\sigma} \\ f_{l\sigma} \end{bmatrix} \qquad \mathbf{E}(k) = \begin{bmatrix} \varepsilon_{1k} & 0 & 0 \\ 0 & \varepsilon_{2k} & 0 \\ 0 & 0 & \tilde{E}_f \end{bmatrix} \tag{5}$$

$$\mathbf{V}_{\text{A}} = \begin{bmatrix} 0 & 0 & Vb^* \\ 0 & 0 & Vb^* \\ Vb & Vb & 0 \end{bmatrix}$$

$$\mathbf{V}_{\text{B}} = \begin{bmatrix} 0 & 0 & 0 \\ 0 & 0 & 0 \\ 0 & 0 & \varepsilon_L \end{bmatrix}$$

in which b and λ can be determined from equation (3) after taking their averages over the disorder ensemble. $\tilde{E}_f\,(= E_f + \lambda)$ represents the renormalised f level on the A site and is generally quite small. As an approximation, we take $E_f = 0$, i.e. $\lambda = |E_f|\varepsilon_L = E_L - \tilde{E}_f$. This is equivalent to replacing constraint $\langle n_f^l \rangle = 1 - |b_l|^2$ with constraint $\langle n_f^l \rangle = 1$ on each A site. This is obviously reasonable in the Kondo (or near-integer) limit.

Now, we use the CPA method [26] to solve the disorder Hamiltonian. The CPA equations are given in [15], where the coherent potential can be written as a 3×3 matrix:

$$\Sigma = \begin{bmatrix} 0 & 0 & Vb^* \\ 0 & 0 & Vb^* \\ Vb & Vb & \Sigma_{ff} \end{bmatrix} \tag{6}$$

after taking $\varepsilon_L(E_L) \to \infty$. In equation (6), $\Sigma_{ff}(z)$ is determined by

$$\Sigma_{ff} F_{ff} = x - 1 \tag{7}$$

in which $\mathbf{F}(z) = (1/N)\Sigma_k \mathbf{G}(k\,z)$ and $\mathbf{G}(k\,z)$ is the matrix Green function in the effective

medium. One can easily see that $\Sigma_{ff}(z) = 0$ for the pure lattice case ($x = 1$). In the dilute limit ($x \to 0$), we have

$$x\Sigma_{ff} \to (1-x)z + V^2|b|^2[F_{11}^0(z) + F_{22}^0(z)] \tag{8}$$

where

$$F_{ii}^0(z) = \frac{1}{N}\sum_k \frac{1}{z - \varepsilon_{ik}} \qquad i = 1, 2.$$

The coherent potential for $0 < x < 1$ can be obtained by solving equation (7) numerically on $z = \omega + i0^+$. The Green functions in the effective medium can be used to determine the mean-field physical quantities of the alloy system. In the following sections, we shall calculate the mean-field electrical resistivity of the system using the above results of the Green functions and the coherent potential.

3. The zero-temperature resistivity of the alloy

The electrical conductivity of the system is mainly due to the conduction electrons because of the very large effective mass of the f electrons. There are two possible sources of scattering of electrons in the system. One is the scattering of the electrons from the impurities (i.e. B atoms). The other is the electrons scattered by the Gaussian fluctuations of the slave bosons on the A sites. First, we consider the former source. The latter will be considered in § 5. Here we assume that the Matthiesson rule is valid in the system so that the resistivities due to the above two scattering mechanisms can be added. Using the Green functions of the conduction electrons in the effective medium, the CPA conductivity of the system in the presence of the impurities only can be written as [14, 27]

$$\begin{aligned}\sigma_m &= \frac{2e^2v_F^2}{3\Omega}\int d\omega\left(-\frac{\partial f}{\partial\omega}\right)\frac{1}{2\pi}\sum_k [[\mathrm{Im}[\bar{G}_{11}(k\omega + i0^+)]^2 + \{\mathrm{Im}[\bar{G}_{22}(k\omega + i0^+)]\}^2]] \\ &= \frac{2e^2v_F^2}{3\Omega_c}\frac{1}{4\pi D}\int d\omega\left(-\frac{\partial f}{\partial\omega}\right)\int_{-D}^{D} d\varepsilon_k \\ &\quad \times \frac{\{\mathrm{Im}[\Sigma_{ff}(\omega + i0^+)]\}^2V^4|b|^4[(\omega - \varepsilon_{1k})^4 + (\omega - \varepsilon_{2k})^4]}{(X^2 + Y^2)^2}\end{aligned} \tag{9}$$

where

$$\begin{aligned}X &= (\omega - \varepsilon_{1k})(\omega - \varepsilon_{2k})\{\omega - \mathrm{Re}[\Sigma_{ff}(\omega + i0^+)]\} - V^2|b|^2[(\omega - \varepsilon_{1k}) + (\omega - \varepsilon_{2k})] \\ Y &= \mathrm{Im}[\Sigma_{ff}(\omega + i0^+)(\omega - \varepsilon_{1k})(\omega - \varepsilon_{2k})].\end{aligned} \tag{10}$$

In equation (9), $\bar{G}_{11}$ and $\bar{G}_{22}$ are Green functions of the conduction electrons in the effective medium [15]. Using the self-consistent equation for the coherent potential, the mean-field resistivity can be determined from equation (9). In the numerical calculation, one has to calculate the coherent potential on the real axis and determine the mean-field parameters b self-consistently before substituting them into equation (9). However, for a special case (i.e. at zero temperature), the calculation is quite simple. Since

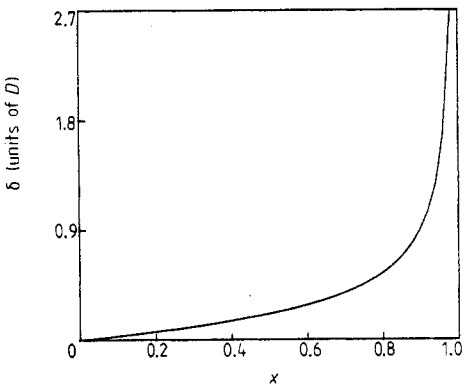

Figure 1. Variation in δ with concentration x; here $\varepsilon_0 = 0.2D$.

$(-\partial f/\partial\omega) = \delta(\omega)$ when $T = 0$ K, one can easily see from equation (9) that the zero-temperature conductivity depends only on the coherent potential at $z = \mathrm{i}0^+$. Because the electron–hole symmetry of the system, we can prove that

$$\mathrm{Re}[\Sigma_{ff}(-\omega + \mathrm{i}0^+)] = -\mathrm{Re}[\Sigma_{ff}(\omega + \mathrm{i}0^+)] \tag{11}$$

and

$$\mathrm{Im}[\Sigma_{ff}(-\omega + \mathrm{i}0^+)] = \mathrm{Im}[\Sigma_{ff}(\omega + \mathrm{i}0^+)]. \tag{12}$$

One can easily show from equation (11) that $\mathrm{Re}[\Sigma_{ff}(\mathrm{i}0^+)] = 0$. With this relation, after some algebra, the zero-temperature conductivity can be written as

$$\sigma_m(0) = \frac{2e^2 v_F^2}{3\Omega_c}\frac{1}{4\pi D}\int_{-D}^{D} d\varepsilon_k \frac{\delta^2[(\varepsilon_k - \varepsilon_0)^4 + (\varepsilon_k + \varepsilon_0)^4]}{[(\varepsilon_k^2 - \varepsilon_0^2)^2 + 4\delta^2\varepsilon_k^2]^2} \tag{13}$$

and thus the resistivity and the resistivity per mole of the magnetic ions are given, respectively, as

$$\begin{aligned} \rho_m(0) &= 1/\sigma_m \\ \bar{\rho}_m(0) &= 1/x\sigma_m. \end{aligned} \tag{14}$$

In equation (13), $\delta = V^2|b|^2/|\mathrm{Im}[\Sigma_{ff}(\mathrm{i}0^+)]|$ and can be determined by the self-consistent equation

$$x = \frac{1}{2D}\int_{-D}^{D} d\varepsilon_k \frac{4\delta^2\varepsilon_k^2}{(\varepsilon_k^2 - \varepsilon_0^2)^2 + 4\delta^2\varepsilon_k^2}. \tag{15}$$

Equation (15) can be easily derived from equation (7) using the relation $\mathrm{Re}[\Sigma_{ff}(\mathrm{i}0^+)] = 0$. Since the Gaussian fluctuation contribution to the resistivity is zero at zero temperature, the present result does represent the residual resistivity of the system. The numerical results of δ and the residual resistivity are shown in figure 1 and figure 2. Figure 2(a) shows that the curve of the residual resistivity against x differs greatly from the usual Nordheim law. Figure 2(b) indicates that the residual resistivity per mole decreases monotonically as the concentration increases from 0.3 to 1.0. This decrease in the resistivity is due to the formation of coherence in the system. In the concentrated limit, i.e. the pure lattice of the magnetic ions, the DC conductivity is infinite in the mean-field theory because of the periodicity. In this limit the only contribution to the resistivity of

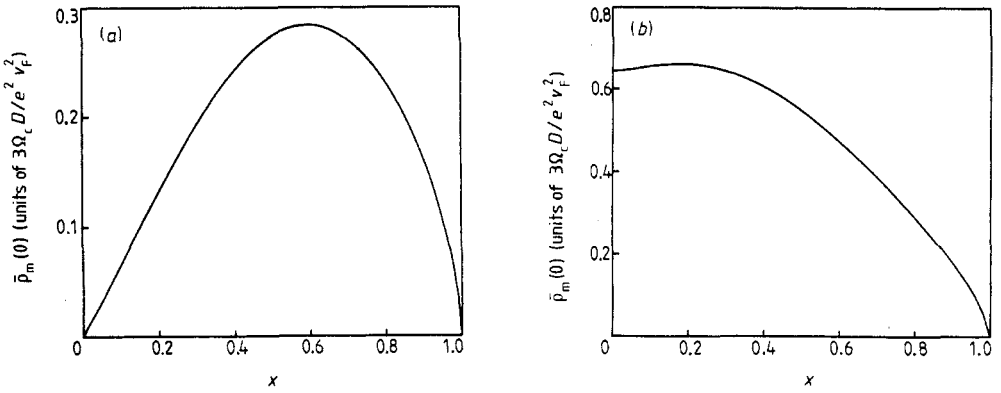

Figure 2. (*a*) Residual electrical resistivity as a function of x. (*b*) Residual electrical resistivity per mole of magnetic ions; here $\varepsilon_0 = 0.2D$.

the system comes from the electrons scattered by the Gaussian fluctuations of the boson fields and occurs only at a finite temperature $T > 0$.

4. The mean-field resistivity at a finite temperature

We shall present the numerical results of the mean-field electrical resistivity in this section. However, before we look at the numerical results, let us first analyse equation (9). Equation (9) can be rewritten as

$$\sigma_m(T) = \frac{2e^2 v_F^2}{3\Omega_c} \frac{1}{2D} \int d\omega \left(-\frac{\partial f}{\partial \omega}\right) \tau(\omega T) \tag{16}$$

where

$$\tau(\omega T) = \frac{x}{2\pi} \int_{-D}^{D} d\varepsilon_k \frac{\{\mathrm{Im}[\Sigma_{ff}(\omega + i0^+)]\}^2 V^4 |b|^4 [(\omega - \varepsilon_{1k})^4 + (\omega - \varepsilon_{2k})^4]}{(X^2 + Y^2)^2} \tag{17}$$

is a quantity which is similar to the transport relaxation time in the Wilson formula. The implicit temperature dependence of τ is through the temperature dependence of b which is nearly temperature independent at a sufficient low temperature. Therefore, at $T \ll T_K$, the temperature dependence of τ can be neglected. Using the electron–hole symmetry relations (11) and (12), one can easily show that $\tau(\omega) = \tau(-\omega)$. So $\tau(\omega)$ can be expanded in the vicinity of the Fermi level as

$$\tau(\omega) = \tau_0 + \tau_1 \omega^2 + \ldots \tag{18}$$

Therefore the low-temperature mean-field conductivity of the system can be written as

$$\sigma_m(T) = (2e^2 v_F^2 / 3\Omega_c)(1/2D)[\tau_0 + (\pi^2/3)\tau_1 T^2] \tag{19}$$

and the mean-field resistivity per mole is

$$\bar{\rho}_m(T) = 1/x\sigma_m(T)$$
$$= \rho_0/[1 + (B/\rho_0)T^2] = \rho_0 - BT^2 + O(T^4) \tag{20}$$

where ρ_0 is the residual resistivity:

$$\rho_0 = \bar{\rho}_m(0) = 3\Omega_\sigma D/e^2 v_F^2 \tau_0 \tag{21}$$

and

$$B = (\pi^2/3)(\tau_1/\tau_0)\rho_0 \tag{22}$$

is the coefficient of the T^2 term of the mean-field resistivity at low temperatures. The

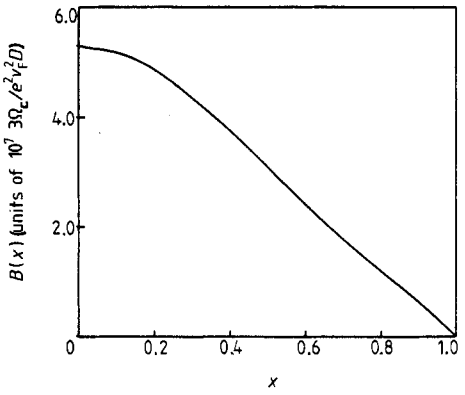

Figure 3. $B(x)$ against x.

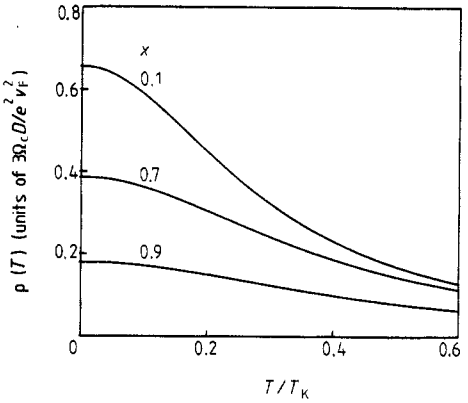

Figure 4. Temperature dependence of the mean-field resistivity per mole of magnetic ions for different concentrations, where we take $\varepsilon_0 = 0.2D$ and $J = V^2/|E_f| = 6.277 \times 10^{-2}D$.

numerical result for B at different x is shown in figure 3. Clearly, B decreases with increasing alloy concentration x, reflecting the growing coherence between magnetic ions. The development of coherence in the system is due to the interaction effect between the Kondo ions which has been included in our CPA scheme through the coherent potential of the effective medium. It is this interaction effect, which appears as the RKKY interaction in the magnetic moment model, this depresses the f–c exchange scattering near the Fermi level [28] and thus induces a decrease in B.

When $x \to 0$, using equation (8), we have

$$\tau(\omega T) = (\omega^2 + \Delta^2)/2V^2|b|^2\Delta \tag{23}$$

where $\Delta = V^2|b|^2/D$ represents the width of the f-electron resonance peak in the single-impurity problem. This result is simply the expression for the relaxation time of the dilute impurity problem.

The numerical results of the mean-field resistivity per mole against T for different concentrations are shown in figure 4. The total resistivity is composed of this mean-field part and a part from the Gaussian fluctuations. We shall discuss the latter part in the next section.

5. The resistivity maximum of the alloy system

From figure 4, we can see that the mean-field parts of the electrical resistivities of various concentrations decrease with increase in the temperature monotonically. At very low temperatures ($T \ll T_K$), it decreases as T^2. In the pure lattice case the mean-field resistivity is zero because our mean-field Hamiltonian can be diagonalised into three quasi-particle bands with the use of the Bogoliubov transformation. Beyond the mean-field approximation, there exists another contribution to the resistivity which comes from the Gaussian fluctuations of the boson fields. It has been shown both experimentally [9] and theoretically [5, 7] that the low-temperature contribution to the resistivity in the HF lattice is determined by the T^2 term and its coefficient A is inversely proportional to

the square of the Kondo temperature T_K within the Gaussian fluctuation approach. However, as pointed out in [29], the T^2 regime only persists at very low temperatures. It will begin to inflect or saturate as T approaches T_K. The inflection of the resistivity curve is due to saturation of the boson fluctuations. Here we give a simple explanation of the inflection in the resistivity in the two-conduction-band HF lattice system. In an earlier work [24], we have demonstrated that the Gaussian fluctuations of the slave boson field can cause the electrons to make a transition from the conduction bands into the f band. Therefore the transport relaxation rate due to the scattering of the electrons off the boson fluctuations is proportional not only to the 'strength' of boson fluctuations (i.e. the T^2 term in [5]), but also to the f-electron DOS near the Fermi level. In our two-conduction-band HF system, the f-electron DOS $N_f(\omega)$ near the Fermi level can be written as

$$N_f^{-1}(\omega) = 2\alpha D[1 + C(\omega/T_K)^2] \tag{24}$$

where $\alpha = m/m^*$ and m^* is the effective mass of the heavy electrons. C is a numerical parameter which can be adjusted by taking different value of ε_0 [24]. Using the Wilson formula and equation (24), we obtained the resistivity due to the boson fluctuations as

$$\rho_{fl}(T) = A_0 T^2/[1 + (T/T_s)^2] \tag{25}$$

where $T_s = (\sqrt{3}/\pi\sqrt{C})T_K$. Here we see that the resistivity deviates from the T^2 law through the denominator $[1 + (T/T_s)^2]$.

Since the boson fluctuation is local and incoherent, the resistivity due to the scattering of the electrons off the boson fluctuations in the alloys may have the same form as equation (25). As a preliminary estimate, we can assume that the resistivity per mole due to the fluctuations in the alloy system is the same as that of the lattice. Therefore the total resistivity of the system can be written as

$$\bar{\rho}(T) = \bar{\rho}_m(T) + \rho_{fl}(T) = \rho_0/[1 + (B/\rho_0)T^2] + A_0 T^2/[1 + (T/T_s)^2] \tag{26}$$

where A_0 and T_s are chosen suitably. The temperature of the resistivity maximum can be determined from

$$d\bar{\rho}(T)/dT = 0. \tag{27}$$

After a simple derivation, we obtain

$$T_{max} = \sqrt{(\sqrt{A_0} - \sqrt{B})/\sqrt{B}(1 - \sqrt{A_0 B}T_s^2/\rho_0)}\, T_s. \tag{28}$$

Obviously, the resistivity maximum appears at a finite temperature T_{max} (>0) for $B(x) < A_0$. There exists a critical concentration x_c satisfying the condition $B(x_c) = A_0$, where T_{max} decreases to zero as x decreases to x_c. Furthermore, when $T < T_K$, the coefficient of the T^2 term in the total resistivity can be expressed as

$$A(x) = A_0 - B(x). \tag{29}$$

$A(x)$ will change its sign from negative to positive with increasing x. The numerical results of T_{max} are shown in figure 5. The calculated resistivity curves of the HF alloys

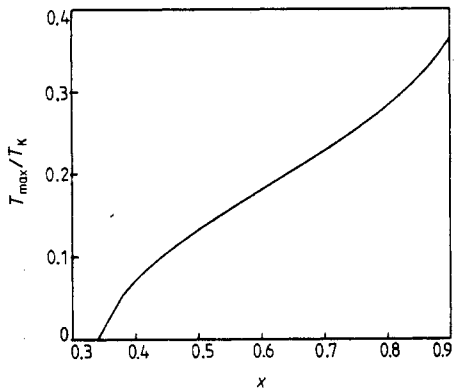

Figure 5. The resistivity maximum temperature T_{max} as a function of x, where we take $A_0 = 4.0 \times 10^7 (3\Omega_c/e^2 v_F^2 D)$ and $T_s = 0.25 T_K$.

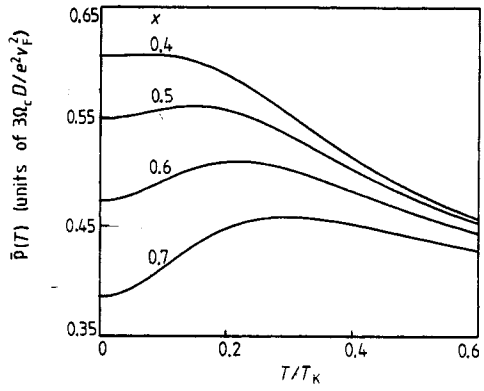

Figure 6. Variation in total electrical resistivities for different x with temperature, where we take $A_0 = 4.0 \times 10^7 (3\Omega_c/e^2 v_F^2 D)$ and $T_s = 0.25 T_K$.

below T_K are given in figure 6. Since in the slave boson theory an unphysical phase transition occurs at $T = T_K$ in the mean-field approximation as well as in the Gaussian fluctuation approach, our calculation is valid only in the regime $T < T_K$. Nevertheless, our low-temperature result can be compared with the experiments in [9] and [30].

6. Conclusions

In summary, we have investigated the electrical resistivity of the HF alloys through a two-conduction-band slave boson alloy Hamiltonian. In the calculation, we have assumed that the resistivity of the system is composed of two independent parts: impurity scattering part and the boson fluctuation part. In the low-temperature region when $T \ll T_K$, the total resistivity follows the $\rho_0 + AT^2$ law and the coefficient of the T^2 term increases and changes its sign from negative to positive with increasing concentration of the magnetic ions. The residual resistivity of the system is calculated in quite a simple way. The occurrence of the resistivity maximum at a finite temperature T_{max} (>0) is also obtained when the alloy concentration $x > x_c$ and contributions $\bar{\rho}_m(T)$ and $\rho_{fl}(T)$ are both taken into account. Our results are in agreement with the experiments in [9–11] qualitatively. We believe that the methods and results of the present paper offer a reasonable explanation for the crossover of the electrical resistivity of HF alloys from single-impurity behaviour to the coherence (concentrated) regime. Finally, we should mention that our discussion on the resistivity is valid only to the order of $1/N$ of the slave boson field, i.e. within the Gaussian fluctuation approximation approach. Beyond the Gaussian fluctuation approximation, as pointed out recently in [31], the RKKY interaction between the quasi-particles in the HF lattice limit exists. This is due to the spin fluctuations involving high-frequency ($\omega > T_K$) fluctuations of the slave boson field and is of great importance to the mechanisms of the superconducting pairing and magnetism for the HF systems. The high-frequency RKKY interaction may provide a small non-universal correction to the resistivity, but it is also important to recognise that the main features of the resistivity in the HF alloys can be understood qualitatively in the Gaussian fluctuation scheme without considering this high-frequency RKKY interaction.

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