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K. VLADÁR,
G.T. ZIMÁNYI

ELECTRON-LOCALIZATION EFFECTS ON IMPURITY DYNAMICS II. TWO-LEVEL SYSTEMS

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II. TWO-LEVEL SYSTEMS

K. VLADÁR, G.T. ZIMÁNYI

Central Research Institute for Physics
H-1525 Budapest 114, P.O.B.49, Hungary

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ABSTRACT

The scattering of electrons on two-level-systems (TLS) in disordered medium is investigated. The obtained perturbative series is summed up in leading logarithmic order. It is found, that localization effects increase T_k, the crossover temperature. Furthermore T_k shows a maximum as a function of disorder. For certain range of parameters a stronger than logarithmic dependence is obtained ($\delta\rho^{d=2} \sim T^{-1}$, $\delta\rho^{d=3} \sim T^{-1/2}$).

АННОТАЦИЯ

Исследуется рассеяние электронов на двухуровневых системах (TLS) в неупорядоченных структурах. Полученный ряд возмущения суммируется в старшем логарифмическом порядке. Обнаружено, что температура кросс-овер T_K повышается за счет локализационных эффектов, а T_K показывает максимум в зависимости от неупорядоченности. В некоторых диапазонах параметров температурная зависимость сильнее логарифмической.

KIVONAT

Elektronoknak kétnivós rendszereken (TLS) való szóródását vizsgáljuk rendezetlen rendszerekben. A kapott perturbációs sort vezető logaritmikus rendben felösszegezzük. Azt találjuk, hogy T_k crossoverhőmérsékletet a lokalizációs effektusok növelik, továbbá T_k maximumot mutat a rendezetlenség függvényében. Bizonyos paraméter-tartományokban a logaritmusnál erősebb hőmérsékletfüggést kapunk.

1. INTRODUCTION

The two-level-systems (TLS) as special excitations of amorphous insulators were suggested by Anderson, Halperin and Varma (1972) and independently by Phillips (1972) to interpret some unusual low temperature behavior of these materials. The ultrasound measurements of Golding et al. (1978) called the attention to the presence of these type of excitations in metallic glasses also, practically in the same concentration as in insulators. Their contributions to different physical quantities can be measured: the resistivity correction by direct method, the TLS relaxation time in ultrasound-absorption experiments and the inelastic electron lifetime, which occurs in the localization theory. In all the three types of measurements several data are available (Black 1981).

It turned out, that the theory, proposed in an earlier paper (Vladár and Zawadowski 1983) is capable to account for these measurements. The model is based on a realistic angular dependent electron-TLS interaction. The applied field-theoretical method makes use of Abrikosov's pseudofermion description of the TLS and the arising logarithmic contributions are summed up with the help of a scaling theory. During the course of scaling the effective coupling becomes strong at a crossover temperature T_k. This enhanced interaction ensures the numerical agreement between theory and experiment. Here we also mention that below T_k a screening electron cloud builds up around the TLS. However the model handles the electrons as free ones, which is surely not the case in a high resistivity metallic

glass. So we address ourselves to the problem of the influence of static impurities on the above picture. We will focus on the change of T_k, as we did for the Kondo-problem in the previous paper (Vladár and Zimányi 1984) (hereafter referred as I.)

2. THE MODEL

The Hamiltonian of the TLS-static impurities problem is very similar to that we used in I. The conduction electrons interact with a single two level system and the randomly distributed static impurities:

$$\mathcal{H} = \mathcal{H}_{el} + \mathcal{H}_{TLS} + \mathcal{H}_{el-TLS} + \mathcal{H}_{imp-el} \tag{2.1}$$

The first and the fourth terms are identical with those in I.

The conduction electron band is represented by

$$\mathcal{H}_{el} = \sum_{\underset{\sim}{k},\sigma} \varepsilon_{\underset{\sim}{k},\sigma}\, a^{+}_{\underset{\sim}{k},\sigma}\, a_{\underset{\sim}{k},\sigma} \tag{2.2}$$

where $a^{+}_{\underset{\sim}{k},\sigma}$ and $a_{\underset{\sim}{k},\sigma}$ are the creation and the annihilation operators of an electron with momentum $\underset{\sim}{k}$ and spin σ.

A simplified density of states

$$\varrho(\varepsilon) = \begin{cases} \varrho & , \text{ if } -D < \varepsilon < D \\ 0 & \text{otherwise} \end{cases} \tag{2.3}$$

is used again.

The interaction between the electrons and the static impurities is described by

$$\mathcal{H}_{imp-el} = u \cdot \sum_{a} \sum_{\underset{\sim}{k},\underset{\sim}{k}'} a^{+}_{\underset{\sim}{k}',\sigma}\, a_{\underset{\sim}{k},\sigma}\, e^{i(\underset{\sim}{k}-\underset{\sim}{k}')\underset{\sim}{r}_a} \tag{2.4}$$

The scattering is elastic and its amplitude u is uniform and isotropic. The random positions of the impurities are denoted by $\underset{\sim}{r}_a$.

The TLS Hamiltonian in second quantized form is:

$$\mathcal{H}_{TLS} = -\lambda \cdot \sum_{\alpha} b_\alpha^+ b_\alpha \tag{2.5}$$

b_α^+ and b_α are the creation and annihilation operators of the pseudofermions (see I), $\alpha = 1,2$. λ is their chemical potential. The energy splitting of the two level is neglected assuming that it is smaller than the temperature. This simplification is also supported by the observation (Vladár and Zawadowski 1983), that the splitting is neither affected itself, nor influences the vertex corrections in leading logarithmic order perturbation theory, provided $\Delta \ll T$.

The electron-TLS interaction is given by

$$\mathcal{H}_{el-TLS} = \sum_{\alpha,\beta,i} \sum_{\underset{\sim}{k},\underset{\sim}{k}'} a_{\underset{\sim}{k}}^+ b_\alpha^+ V^i_{\underset{\sim}{k}\underset{\sim}{k}'} \sigma^i_{\alpha\beta} b_\beta a_{\underset{\sim}{k}'} \tag{2.6}$$

where $\underset{\sim}{\sigma}^i$ -s are the Pauli-matrices. The properties of $\underset{\sim}{V}^i$ matrices were investigated in detail by Vladár and Zawadowski (1983). The most important features are:

1.) $$V^z_{\underset{\sim}{k},\underset{\sim}{k}'} \gg V^x_{\underset{\sim}{k},\underset{\sim}{k}'} \tag{2.7}$$

This means, that when an electron scatters on the TLS, in most

of the cases the latter will stay in its initial state, i.e. electron assisted tunneling occurs only with small probability.

2.)
$$V^{n}_{\underset{\sim}{k},\underset{\sim}{k}'} = V^{n}_{\underset{\sim}{k}',\underset{\sim}{k}} \qquad n = x, z$$
$$V^{y}_{\underset{\sim}{k},\underset{\sim}{k}'} = -V^{y}_{\underset{\sim}{k}',\underset{\sim}{k}} \tag{2.8}$$

This is a consequence of the time reversal symmetry.

3.)
$$V^{y}_{\underset{\sim}{k},\underset{\sim}{k}'} = 0 \tag{2.9}$$

if the conduction electrons interact with the atom (or atoms) of the TLS by local interaction.

In the course of the calculations a sphere wave representation turned out to be the most convenient in which the matrix $\underset{\sim}{V}^{z}$ is diagonal.

3. PERTURBATION THEORY

3.1 The pure TLS problem

The applied perturbation technique is identical in principle with the one we used in I. We introduce the electron and pseudofermion Green-functions (I. Chap. 3.), and include the effects of the disorder through τ , the elastic electron lifetime and the pole contributions (I. Chap. 3.1).

Without the static impurities for the electron-TLS problem a summation of the parquet vertex diagrams can be performed, too. An equivalent but more simple method is to solve the first order scaling equations (Vladár and Zawadowski 1983). The scaling equations can be written from the first order vertex contributions (I. fig. 6.)

$$\frac{\partial V^{i}_{mn}}{\partial x} = -2i\varrho\, \varepsilon^{ijk}\, V^{j}_{mm'} V^{k}_{m'n} \qquad (3.1)$$

where $x = \ln \frac{D}{T}$

Integrating the equations the couplings become relevant in a two dimensional subspace and from here the problem will be equivalent with the well-known case of magnetic impurities. In complete accordance with the anisotropic Kondo-problem the resulting temperature dependence of the vertex shows a drastic increase at the crossover temperature

$$T_c = D \cdot \left(\frac{V_\circ^x}{4V_\circ^z}\right)^{\frac{1}{4\varrho V_\circ^z}} \qquad (3.2)$$

In the combined problem the two-dimensionality of the relevant subspace will not be proved: we assume that the substantial properties of the system can be modeled by two electron sphere-wave in this case also. The same assumption was used by Zawadowski (1980). Then the electron-TLS interaction takes the form:

$$\mathcal{H}_{el\text{-}TLS} = \sum_{\substack{\alpha,\beta \\ m,n}} \sum_{i,j} a_m^+ b_\alpha^+ V_j^i \sigma_{\alpha\beta}^i \sigma_{mn}^j b_\beta a_n \qquad (3.3)$$

Applying a suitable orthogonal transformation and using the condition (2.8) this turns into (Vladár and Zawadowski 1983):

$$\mathcal{H}_{el\text{-}TLS} = \sum_{\substack{\alpha,\beta \\ m,n}} \sum_{i=x,z} a_m^+ b_\alpha^+ V^i \sigma_{\alpha\beta}^i \sigma_{mn}^i b_\beta a_n \qquad (3.4)$$

In this model the (3.1) scaling equations are:

$$\begin{aligned} \frac{\partial V^x}{\partial x} &= 4\rho V^y V^z \\ \frac{\partial V^y}{\partial x} &= 4\rho V^x V^z \\ \frac{\partial V^z}{\partial x} &= 4\rho V^x V^y \end{aligned} \qquad (3.5)$$

If $V^x = V_0^x$, $V^y = 0$, $V^z = V_0^z$ at $x = 0$, the solution is (Shiba 1970, Vladár and Zawadowski 1983):

$$\begin{aligned} x &= \frac{1}{4\rho V_0^z} F(\alpha, q) \\ \alpha(x) &= \operatorname{arc\,tg} \frac{V^y(x)}{V_0^x} \quad , \quad q = \frac{\left(V_0^{z\,2} - V_0^{x\,2}\right)^{1/2}}{V_0^z} \end{aligned} \qquad (3.6)$$

where $F(\alpha,q)$ is the elliptic integral of the first kind. The other two couplings are obtained as:

$$V^x(x) = \left(V^y(x)^2 + V_0^{x\,2} \right)^{1/2}$$
$$V^z(x) = \left(V^y(x)^2 + V_0^{z\,2} \right)^{1/2} \qquad (3.7)$$

3.2 Unified problem

Turning to the combined problem one has to evaluate again all the parquet contributions with finite lifetime Green--functions; furthermore diffusive and maxcross poles must be inserted between any two electron lines.

The finite electron lifetime τ in the Green-function causes only a

$$g \to g' = g(1-g') \quad , \quad g' = \frac{1}{\pi D \tau} \sim g = \frac{1}{2E_F \tau} \qquad (3.8)$$

decrease in the electron density of states at the Fermi surface (see I.). A ladder or a maxcross pole between two electron lines results a $\frac{3}{\pi}\, g^2$ multiplicative factor after the frequency summation except the case when the pole connects the incoming and the outgoing electron lines: then the maxcross pole results $\frac{6}{\pi}\, g^2$ and the ladder pole does not give contribution (see I.). However concerning the summation of the electron states the methods presented in I. no longer can be applied directly. The poles restrict the momenta of the connected

electron lines since in the integral (I. 3.15.) contributions from the region $q \ll \frac{1}{v_F \tau}$ are important only. The restrictions are represented in fig. 1.

The product of the coupling matrices originally had the form:

$$\int d\underset{\sim}{k}_2 \dots d\underset{\sim}{k}_n \; V^{i_1}_{\underset{\sim}{k}_1 \underset{\sim}{k}_2} \cdot V^{i_2}_{\underset{\sim}{k}_2 \underset{\sim}{k}_3} \cdot \dots \cdot V^{i_n}_{\underset{\sim}{k}_n \underset{\sim}{k}_{n+1}} \equiv V^{(n)}_{\underset{\sim}{k}_1, \underset{\sim}{k}_{n+1}} \tag{3.9}$$

Turning to sphere-wave representation and using the simplified matrix notation the product is

$$\underset{\sim}{V}^{(n)} = \underset{\sim}{V}^{i_1} \cdot \underset{\sim}{V}^{i_2} \cdot \dots \cdot \underset{\sim}{V}^{i_n} \tag{3.10}$$

If two of the electron lines are connected with a ladder pole, then the matrix product assumes the form:

$$\int d\underset{\sim}{k}_2 \dots d\underset{\sim}{k}_n \; V^{i_1}_{\underset{\sim}{k}_1 \underset{\sim}{k}_2} \cdot \dots \cdot V^{i_p}_{\underset{\sim}{k}_p \underset{\sim}{k}_{p+1}} \cdot V^{i_{p+1}}_{\underset{\sim}{k}_{q+1} \underset{\sim}{k}_{p+2}} \cdot \dots \cdot V^{i_q}_{\underset{\sim}{k}_q \underset{\sim}{k}_{q+1}} \cdot \tag{3.11}$$

$$\cdot V^{i_{q+1}}_{\underset{\sim}{k}_{p+1} \underset{\sim}{k}_{q+2}} \cdot \dots \cdot V^{i_n}_{\underset{\sim}{k}_n \underset{\sim}{k}_{n+1}} = V^{(n)}_{\underset{\sim}{k}_1 \underset{\sim}{k}_{n+1}, \text{ladder}}$$

so that

$$\underset{\sim}{V}^{(n)}_{\text{ladder}} = \underset{\sim}{V}^{i_1} \cdot \underset{\sim}{V}^{i_2} \cdot \dots \cdot \underset{\sim}{V}^{i_p} \cdot \underset{\sim}{V}^{i_{q+1}} \cdot \dots \cdot \underset{\sim}{V}^{i_n} \cdot \mathrm{Tr}\left(\underset{\sim}{V}^{i_{p+1}} \cdot \dots \cdot \underset{\sim}{V}^{i_q}\right) \tag{3.12}$$

If two of the electron lines are connected with a maxcross pole, then the matrix product is equal to:

$$\int d\underset{\sim}{k}_2 \dots d\underset{\sim}{k}_n \, V^{i_1}_{\underset{\sim}{k}_1 \underset{\sim}{k}_2} \cdot \dots \cdot V^{i_p}_{\underset{\sim}{k}_p \underset{\sim}{k}_{p+1}} \cdot V^{i_{p+1}}_{-\underset{\sim}{k}_{q+1} \underset{\sim}{k}_{p+2}} \cdot \dots$$

$$\cdot V^{i_q}_{\underset{\sim}{k}_q, -\underset{\sim}{k}_{p+1}} \cdot V^{i_{q+1}}_{\underset{\sim}{k}_{q+1} \underset{\sim}{k}_{q+2}} \cdot \dots \cdot V^{i_n}_{\underset{\sim}{k}_n \underset{\sim}{k}_{n+1}} = V^{(n)}_{\underset{\sim}{k}_1 \underset{\sim}{k}_{n+1}\ maxcross} \qquad (3.13)$$

Using the (2.8-9) properties of the coupling matrices this turns into

$$\underset{\sim}{V}^{(n)}_{maxcross} = \underset{\sim}{V}^{i_1} \cdot \underset{\sim}{V}^{i_2} \cdot \dots \cdot \underset{\sim}{V}^{i_p} \cdot \underset{\sim}{V}^{i_q} \cdot \underset{\sim}{V}^{i_{q-1}} \cdot \dots \cdot \underset{\sim}{V}^{i_{p+1}} \cdot \underset{\sim}{V}^{i_{q+1}} \cdot \dots \cdot \underset{\sim}{V}^{i_n}$$

As $\underset{\sim}{V}^i = V^i \underset{\sim}{\sigma}^i$ (see 3.4) we can connect the three types of the matrix products with the use of the following relation:

$$\underset{\sim}{\sigma}^{i_1} \cdot \dots \cdot \underset{\sim}{\sigma}^{i_n} + (-1)^{q-p} \underset{\sim}{\sigma}^{i_1} \cdot \dots \cdot \underset{\sim}{\sigma}^{i_p} \cdot \underset{\sim}{\sigma}^{i_q} \cdot \dots \cdot \underset{\sim}{\sigma}^{i_{p+1}} \cdot$$

$$\cdot \underset{\sim}{\sigma}^{i_{q+1}} \cdot \dots \cdot \underset{\sim}{\sigma}^{i_n} = \underset{\sim}{\sigma}^{i_1} \cdot \dots \cdot \underset{\sim}{\sigma}^{i_p} \cdot \underset{\sim}{\sigma}^{i_{q+1}} \cdot \dots \cdot \sigma^{i_n} \cdot \qquad (3.14)$$

$$\cdot Tr\left(\underset{\sim}{\sigma}^{i_{p+1}} \cdot \dots \cdot \underset{\sim}{\sigma}^{i_q}\right)$$

This relation is proved in Appendix A. The equation (3.14) gives the sum of contributions coming from the two kind of poles for odd (q-p) (the number of interaction points along the electron Green functions, between the two ends of the pole) and the difference for even (q-p). Unfortunately we didn't manage to determine the sum for even (q-p). In Appendix B we argue in a reasonable way, that the total contribution of parquet-graphs containing one maxcross pole for even (q-p) can be neglected, so it is allowed to change the sign of all these terms. By this

approximation we can reduce the problem to the anisotropic Kondo model in disordered medium and the results of I. can be directly applied. On the basis of the arguments explained there the vertex can be obtained from the pure Kondo vertex by the following operation:

$$\underset{\sim}{V}^{i}(x,g) = \left(1 + \frac{3}{\pi} g^2 \frac{\partial^2}{\partial x^2}\right) \underset{\sim}{V}^{i}\left((1-g')x, g\right) \qquad (3.15)$$

$$x = \ln \frac{D}{T}$$

and the right hand side must be expanded up to g'^2 .

The vertex function was expressed by Vladár and Zawadowski (1983) in an implicit equation. Near T_k this expression can be transformed to explicit form. Furthermore in this "critical" region the three V^i -s are close to one another and the scaling equations are identical with those in the anisotropic Kondo problem. In complete accordance with that case, the cross-over temperature T_k where the dimensionless couplings $\varrho \cdot V^i$ grow to unity is modified by the disorder (cf. formula I(3.38)):

$$T_k(g) = T_k(0) \cdot \exp\left(\frac{3\alpha}{4\pi} \cdot \frac{g^2 \left(\ln \frac{V_0^x}{4V_0^z}\right)^2}{\varrho^2 \cdot V_0^{z\,2}} - \right. \qquad (3.16)$$

$$\left. - \frac{g' \cdot \ln \frac{V_0^x}{4V_0^z}}{4\varrho V^2} \right)$$

where $T_k(0) = D \cdot \left(\frac{V_0^x}{4V_0^z}\right)^{\frac{1}{4V_0^z \varrho}}$

and $\alpha = 1$. If one performs an analogous calculation for the pseudofermion lifetime, the same expression is recovered for T_k with the value of $\alpha = 3/2$.

When calculating the lifetime of electrons, besides the one pole contributions considered above, certain two-pole-processes are also important (fig. 2.). (These are the analogues of the Fukuyama-contributions in the Kondo case (Ohkawa et al. 1983)) The arising energy-integrals again yield $\sim T^{-1/2}$ or $\sim T^{-1}$ contributions in 3 and 2 dimensions, respectively. The summation of the parquet-type processes is performable and in accordance with (I. 3.32.) we find:

$$\frac{1}{\tau_{el}} = \frac{3\pi}{2} \cdot \frac{1}{\rho} \cdot \left[\frac{1-g'-\frac{6}{\pi}g^2}{(x_c-x)^3} - \frac{2g'(1-g')x}{(x_c-x)^3} + \right.$$

$$\left. +3\,\frac{g'^2x^2+\frac{6}{\pi}g^2x_c^2}{(x_c-x)^4} + g^2\,\frac{x_c}{(x_c-x)^3}\,\frac{1}{\sqrt{T\tau}} \right] \qquad (3.17)$$

where

$$x_c = \frac{1}{4V_0^2}\ln\frac{4V_0^2}{V_0^x}$$

and we assumed to have reached the isotropic region.

In two dimensions these processes are more singular, so the summation of logarithms is not necessary anymore:

$$\delta\left(\frac{1}{\tau_{el}}\right) = \frac{\pi}{\rho}\,V_0^z \cdot V_0^{x^2} \cdot g \cdot \frac{1}{T\tau} \qquad (3.18)$$

3.3 DISCUSSION

In conclusion we have found, that besides some mathematical differences the two-level-systems in disordered medium behave analogously to the Kondo impurity.

i.) in 3 dimensions the crossover temperature shows the competition of two effects in principle, however - since the metallic glasses are strongly disordered systems, possessing large residual resistivity - the localization effects are expected to dominate, thus increasing T_k.

A consequence of this result is, that for materials of the $(A_x B_{1-x}) C$ type T_k is a nonmonotonic function of x: concerning the whole composition range ($0 < x < 1$) it shows a maximum, since the degree of disorder is expected to be maximal at an intermediate value of x.

ii.) when conferring the Fukuyama-type and normal processes, the ratio $V_0^2/\sqrt{T_k \cdot \tau}$ determines the importance of the former. Since in metallic glasses the mean time between elastic scatterings τ is much shorter than in dilute alloys; on the other hand V_0^2 can be relatively large ($Pd_{30}Zr_{70}$: $V_0^2 = 0.16$, $Nb_{20}Zr_{80}$: $V_0^2 = 0.12$) (G. Weis et al. 1980, N. Thomas et al. 1980), this ratio might exceed unity considerably. In this case we expect these processes to surpass the logarithmic contributions. Furthermore in two dimensions the T^{-1} behavior is surely the most important one. Unfortunately we do not know about such experiments, which would have been performed in the whole composition range in reliably non-magnetic materials.

There are measurements, where $T_k(x)$ increased for small values of x (Mizutani 1983), however - as noticed in I. - this monotonic change can be accounted for in other ways, too.

Finally we mention, that the same problem has been attacked by Stone, Joanopoulos and Thouless (1983) in lowest order of perturbation theory. In calculating the electron lifetime they apply a phase-space restriction for the arguments of the scattering amplitude $V_{\underset{\sim}{k}_1 \underset{\sim}{k}_2}$. We see no back up for this assumption. Indeed, abandoning this restriction we obtain equally important contributions to the integrals from the neglected regions.

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Appendix A

We prove here the following relation

$$\underset{\sim}{\sigma}^{i_1}\cdot\underset{\sim}{\sigma}^{i_2}\cdot\ldots\cdot\underset{\sim}{\sigma}^{i_p}+(-1)^p\,\underset{\sim}{\sigma}^{i_p}\underset{\sim}{\sigma}^{i_{p-1}}\ldots\underset{\sim}{\sigma}^{i_1}=\mathrm{Tr}(\underset{\sim}{\sigma}^{i_1}\cdot\underset{\sim}{\sigma}^{i_2}\cdot\ldots\underset{\sim}{\sigma}^{i_p})\cdot\underset{\sim}{1} \qquad (A.1)$$

Let the number of $\underset{\sim}{\sigma}^x, \underset{\sim}{\sigma}^y, \underset{\sim}{\sigma}^z$ matrices be denoted by p_x, p_y and p_z in one of the above products. Eight different cases must be considered: when 0, 1, 2 or 3 of p_x, p_y and p_z are even. For example when only p_x is even, the product under the trace is $i\sigma^x$ or $-i\sigma^x$ so the r.h.s. is zero. The second term of the l.h.s. is the hermitean adjoint of the first terms, so they cancel and the l.h.s. is also zero.

The validity of (A.1) for the further cases can be seen by applying similarly simple considerations.

Appendix B

Besides the formula (A.1) we know an additional connection between the two terms of its l.h.s.

$$\underset{\sim}{\sigma}^{i_p}\underset{\sim}{\sigma}^{i_{p-1}}\cdot\ldots\cdot\underset{\sim}{\sigma}^{i_1}=(-1)^{\tilde{p}}\cdot\underset{\sim}{\sigma}^{i_1}\underset{\sim}{\sigma}^{i_2}\cdot\ldots\cdot\underset{\sim}{\sigma}^{i_p} \qquad (B.1)$$

where

$$\tilde{p}=p_x p_y+p_x p_z+p_y p_z \qquad (B.2)$$

p_x, p_y and p_z are the number of $\underset{\sim}{\sigma}^x$, $\underset{\sim}{\sigma}^y$ and $\underset{\sim}{\sigma}^z$ matrices in the product.

In our problem, since $V^{y}_{\underset{\sim}{k}\,\underset{\sim}{k}'}=0$ (cf. 2.9), only $\underset{\sim}{\sigma}^x$ and $\underset{\sim}{\sigma}^z$ matrices occur. In this case:

$$\tilde{p}=p_x p_z \qquad (B.3)$$

Here is the important difference between the p odd and p even situations. In the former the parity of $\tilde{p}$ is always even, whereas for the latter the parity of $\tilde{p}$ is even in half of the cases and odd in the remaining ones.
With this at hand, we examine the summation of the terms, containing a maxcross pole. Let us fix the arrangement of the matrices in a given graph-contribution and then sum up for the possible insertions of the maxcross pole (p held constant). Then the parity of $\tilde{p}$ will change from term to term. This yields an oscillating sum. Here we remark, that as long as finite order contributions are considered, the perturbative series will

diverge only at T=0, so the non-zero characteristic temperature is the result of the high order terms. Regarding them, the vanishing of the oscillating sum seems to be a good approximation. Put in an other way, an other estimation for this oscillating contribution shows, that it is less divergent around T_k, than those taken into account.

This procedure is also supported by the fact, that the previously observed deep analogy with the Kondo-problem (see also Vladár and Zawadowski 1983) is recovered in this way.

REFERENCES

Anderson P W, Varma C M, Halperin B 1972 Phil. Mag. 25 1

Black J L, 1981 Glassy Metals I, Springer

Golding B et al., 1978 Phys. Rev. Lett. 41 1487

Mizutani U, 1983 Progr. in Mat. Sci. 28 97

Ohkawa J, et al. 1983 J. Phys. Soc. Jpn. 52 1701

Philips W A, 1972 J. Low Temp. Phys. 7 351

Shiba H, 1970 Progr. Theor. Phys. 43 601

Stone A D et al., 1983 preprint

Thomas N et al., 1980 Solid State Comm. 33 523

Vladár K, Zawadowski A 1983 Phys. Rev. B28, 1564, 1582, 1596

Vladár K, Zimányi G T this issue, previous article

Weis G et al., 1980 Solid State Comm. 33 111

Zawadowski A 1980 Phys. Rev. Lett. 45 211

FIGURE CAPTIONS

Fig. 1. The phase space restrictions on electron Green--functions caused by the diffusive poles

Fig. 2. Contributions to the electron lifetime

Figure 1.

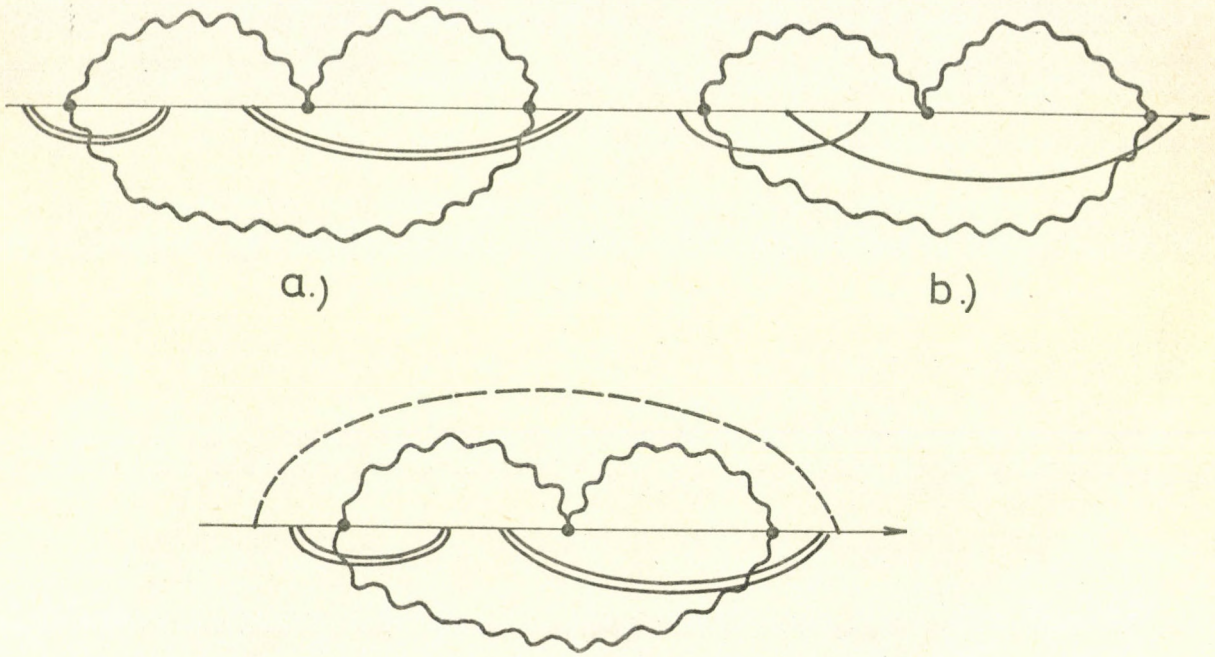
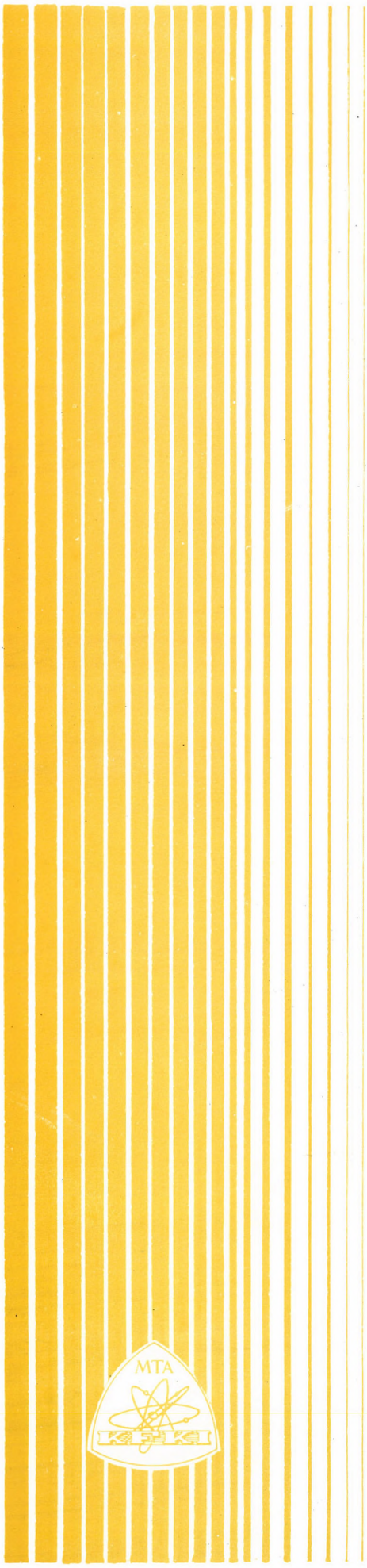


Figure 2.



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