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Renormalization of the hopping parameters in quasi-one-dimensional conductors in the presence of a magnetic field

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Abstract. We consider the competition between the one dimensionalization effect due to a magnetic field and the hopping parameters in quasi-one-dimensional conductors. Our study is based on a perturbative renormalization group method with three cut-off parameters, the bandwidth E_0, the 1D-2D crossover temperature T_1^*, which is related to the hopping process t_1, and the magnetic energy ω_c. We have found that the renormalized crossover temperatures T_1^* and T_2^*, at which the respectively hopping processes t_1 and t_2 become coherent, are reduced compared to the bare values as the field is increased. We discuss the consequences of these renormalization effects on the temperature-field phase diagram of the organic conductors.

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1 Introduction

The organic compounds $(TMTSF)_2X$ and $(TMTTF)_2X$ (noted $(TM)_2X$), where the anion X= PF_6, ClO_4..., exhibit a one dimensional character due to the anisotropy of the hopping integrals along the three directions. The estimated values of these integrals along the high conducting direction (**a** axis) and the two perpendicular ones (**b** and **c** axes) are in the ratios 100:10:1. Therefore, one can consider these materials as systems of one-dimensional chains coupled by small interchain hoppings [1]. The latter play a crucial role in the physics of the $(TM)_2X$ compounds as they span different dimensional regimes. At high temperature, $(TM)_2X$ systems behave as isolated 1D chains which may be described by a Luttinger liquid (Fig. 1). Indeed, there is a general consensus on the non-Fermi liquid character of the 1D normal state [1,2]. A large number of experimental and theoretical studies agree with a Luttinger liquid picture at high temperature and proved the strong character of the electron interactions in these compounds [1–3].

By decreasing the temperature, the interchain hopping becomes coherent and the dimensionality of the system is raised. Then, the system may, either undergoes a dimensional crossover from a Luttinger liquid to a 2D (or 3D) Fermi liquid at a critical temperature T_1^* or undertakes a phase transition to a long range order state. At very low temperature, the generic phase diagram of $(TM)_2X$ compounds shows various ordered phases

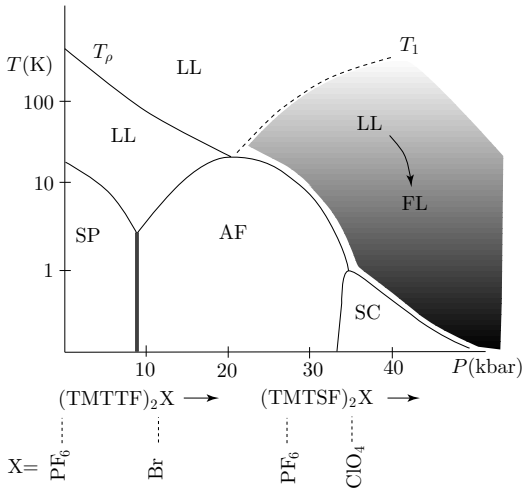

Fig. 1. The generic phase diagram of the $(TM)_2X$ as a function of pressure. LL (FL) denotes Luttinger (Fermi) liquid. The ordered states SP, AF and SC correspond respectively to the Spin-Peierls, the Antiferromagnetic and the superconducting states, after [1].

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(Fig. 1). By varying the anisotropy of these systems, the relative stability of the different phases may be substantially changed. Indeed, applying a magnetic field in the least conducting direction (**c** axis), induces a cascade of phase transitions between Field Induced Spin Density Wave (FISDW) states, which have been observed in the Bechgaard salts $(TMTSF)_2X$. These phases are characterized by the quantized Hall resistance $\rho_{xy} = h/2Ne^2$ in the sequence $N = ...5, 4, 3, 2, 1, 0$ as the field is increased. This phenomenon is well explained by the Quantized Nesting Model (QNM) [4–6]. According to this model, the orbital effect of the magnetic field destabilizes the metallic state by inducing a sort of Peierls instability to a SDW phase. This is possible, thanks to a variation of the nesting vector $\vec{Q}$, the longitudinal component of which, Q_x is quantized as $Q_x = 2k_F + NG$, where $G = ebH/hc$ is the magnetic wave vector, b is the interchain distance and H is the field magnitude.

Below T_1^*, nesting properties at zero magnetic field are based on the electron-hole symmetry relation $\epsilon(\vec{k}) = -\epsilon(\vec{k} + \vec{Q}_0)$ at $\vec{Q}_0 = (2k_F, \pi/b)$. This relation is vulnerable to small deviations in the dispersion relation $\epsilon(\vec{k})$. Such deviations are parametrized by the next-nearest-neighbor interchain hopping t_2 along the **b** direction which is a key issue in the formation of the FISDW phases. An applied magnetic field along the **c** direction may overcome the effect of t_2 by inducing, as we have noted, perfect nesting, provided the above quantization of $Q_x$is fulfilled. As the magnetic field is increased, longitudinal nesting conditions at $Q_x = 2k_F$ is improved.

Understanding how the t_2 hopping process is confined under a magnetic field is of particular interest for the physics of the FISDW phases and the high field phase diagram of the Bechgaard salts.

In this paper we focus on the competition between the t_2 hopping effect and the unidimensionalization of the electron motion induced by the magnetic field. Within a perturbative renormalization group approach, we propose a model to discuss how the t_2 parameter is renormalized by the magnetic field. By the way, we take the opportunity of also studying, at very high magnetic field, the renormalization of the first-nearest interchain hopping parameter t_1 along the **b** direction. This paper is organized as follows. In Section 2, we introduce the model for interacting particles under a magnetic field. In Section 3, the renormalization group equations for the hopping processes are derived by means of a renormalization procedure with three cut-off parameters, the bandwidth E_0, the crossover temperature T_1^* and the magnetic field energy $\omega_c = v_F G$. The results are discussed in Section 4. Finally we conclude our work in Section 5.

2 The model

We consider a system of an infinite number of chains weakly coupled *via* interchain one-particle hopping processes t_1 and t_2 respectively to the first and to the second-nearest neighbors along the **b** direction. The intrachain one-particle dispersion is linearized at the Fermi points $\pm k_F$ with the Fermi velocity v_F $(-v_F)$ for the right-moving (left-moving) electron. The dispersion relation reads then as:

$$\epsilon(\vec{k}) = v_F(|k| - k_F) - 2t_1 \cos k_\perp b - 2t_2 \cos 2k_\perp b$$

k is the longitudinal momentum whereas $k_\perp$ is the transverse momentum along the **b** direction. We will neglect the inter-plane hopping along the **c** axis since the electron motion along this axis is not affected by a transverse magnetic field.

Within the *g-ology* model [7], the intrachain interactions consist of backward scattering g_1, forward scattering g_2 and umklapp scattering g_3, which is due to the dimerization along the TM's chains. The g_i $(i = 1, 2, 3)$ constants are dimensionless and are expressed in units of πv_F.

Weakly coupled chains have been extensively studied in terms of weak coupling Renormalization Group (RG) approach [8–13]. Bourbonnais *et al.* [9,13,14] argued that the electron-electron interactions renormalize the effective interchain hopping t_1^*, which can be much smaller than the bare single particle interchain hopping t_1. On the other hand, it has been stressed that umklapp scattering strongly suppresses the t_1 process, which becomes irrelevant, inducing a deconfinement to confinement transition [10,12].

Here we focus on the case where the interchain hopping parameter t_1 is relevant and the electron motion is deconfined, since we are interested in the behavior of the t_2 process under a magnetic field. Such behavior should be studied in the deconfined 2D regime below the crossover temperature T_1^* as t_2 is a two-dimensional characteristic energy.

The methodology used in this paper is based on the perturbative renormalization group approach (PRG), as discussed by Bourbonnais *et al.* [8,9,13], which we generalize to take into account the presence of three cut-off's. It is worth noting that the PRG method with two cut-off parameters has been used in the case of one-dimensional electron system [15–18] and spin ladder materials [19]. The PRG approach has been modified to take care of the second energy scale involved in the problem, besides the bandwidth energy E_0.

Under a magnetic field $\vec{H} = (0, 0, H)$ applied along the least conducting direction **c**, and in the Landau gauge $\vec{A} = (0, Hx, 0)$, the interchain hopping Hamiltonian is given by:

$$\begin{aligned} H_\perp = & -t_1 \sum_{\langle l,l' \rangle} \int \mathrm{d}x \, \exp\left[\mathrm{i}e \int_{x,ld}^{x,l'\mathrm{d}} \vec{A}(\vec{s}) \cdot \mathrm{d}\vec{s}\right] \\ & \times \psi^+(x,l) \, \psi(x,l') \\ & - t_2 \sum_{\langle\langle l,l' \rangle\rangle} \int \mathrm{d}x \, \exp\left[\mathrm{i}e \int_{x,l\mathrm{d}}^{x,l'\mathrm{d}} \vec{A}(\vec{s}) \cdot \mathrm{d}\vec{s}\right] \\ & \times \psi^+(x,l) \, \psi(x,l') \end{aligned}$$

where $\psi(x,l)$ ($\psi^+(x,l)$) is the annihilation (creation) operator of a particle on the l^{th} chain and x is the coordinate along the chain direction. $\langle l,l'\rangle$ ($\langle\langle l,l'\rangle\rangle$) denotes the first (the second)-nearest neighbors.

Given our gauge choice, $H_\perp$ takes the form:

$$H_\perp = -t_1 \sum_{\langle l,l'\rangle} \int \mathrm{d}x\ \exp\left[\mathrm{i}Gx(l'-l)\right] \psi^+(x,l)\ \psi(x,l') - t_2 \sum_{\langle\langle l,l'\rangle\rangle} \int \mathrm{d}x\ \exp\left[\mathrm{i}Gx(l'-l)\right] \psi^+(x,l)\ \psi(x,l').$$

At this stage, it is useful to introduce the mixed representation by taking the Fourier transform with respect to x [5]. $H_\perp$ is then written as:

$$H_\perp = -t_1 \sum_{\langle l,l'\rangle,k} \psi^+(k,l)\ \psi(k+G(l'-l),l') - t_2 \sum_{\langle\langle l,l'\rangle\rangle,k} \psi^+(k,l)\ \psi(k+G(l'-l),l'). \quad (1)$$

The coupling between the magnetic wave vector $\vec{G} = (G,0,0)$ and the longitudinal momentum k expresses the orbital effect of the field.

Since t_2 is of the order of $t_1/10 \sim 10$ K, a magnetic field of 10 T–30 T is sufficient to bypass the effect of this hopping process. Such field values correspond to the domain where the FISDW phases of Bechgaard salts are formed.

3 Renormalization group formulation

The nature of the metallic 2D phase in the temperature range from 10 to 100 K is still an open question. Deviations from Fermi liquid behavior are clearly present in this regime although the compounds are not in a one-dimensional state [1,2]. Therefore, the influence of the high temperature 1D regime turns out to be crucial for the physical properties of the 2D phase. The 1D effects should be taken into account to investigate the low temperature metallic state.

Within the PRG, the *history* of the system in the 1D regime is conserved while it is relegated in mean filed theories.

On the other hand, in the 1D phase, the response function in either the Cooper or the Peierls channel gives logarithmic divergences at all orders of the perturbative expansion of the scattering amplitudes. To handle these divergences, the PRG is an appropriate approach within which it is possible to sum up divergences of the perturbative series.

Within the PRG, the temperature is parametrized as $T(l) = E_0 \mathrm{e}^{-l}$ where l is the scaling parameter and E_0 is the bandwidth cutoff, which is of the order of the Fermi energy. During the scaling procedure, one moves from the high-temperature scales corresponding to the 1D regime, to the low-temperature scales where interchain hopping processes are deconfined.

3.1 Renormalization group equations for the two-particle scattering processes

In the 1D regime where $T > T_1^*$, the scaling equations of the intrachain scatterings in a two-loop approximation are given by [9]:

$$\begin{aligned} \frac{\mathrm{d}g_1}{\mathrm{d}l} &= -g_1^2 - \frac{1}{2}g_1^3 \\ \frac{d\,(2g_2-g_1)}{\mathrm{d}l} &= g_3^2\left[1-\frac{1}{2}\,(2g_2-g_1)\right] \\ \frac{\mathrm{d}g_3}{\mathrm{d}l} &= g_3\,(2g_2-g_1)\left[1-\frac{1}{4}\,(2g_2-g_1)\right] - \frac{1}{4}g_3^3. \end{aligned} \quad (2)$$

It is worth to note that one may take into account, when deriving the RG equations of the g_i couplings, the effects of the dimensionality and those of the magnetic field. This is possible by including the transverse hopping terms in the determination of the corrections to the scaled coupling constants. Then, the first and the second terms in the r.h.s of the scaling equations in equation (2) have to be multiplied by dimensional crossover functions [20,21]. However, these effects are not relevant in the 1D regime since $\omega_c \ll T_1^*$. It should be stressed that the dimensional crossover functions may play an important role in the study of transient regime, due to thermal fluctuations, in the neighborhood of T_1^* [22]. We will not discuss further this effect, which is not the scope of our present work.

Below T_1^*, in the 2D regime, nesting properties become relevant. The corrections to the two-particle vertices corresponding to the g_i ($i = 1-3$) constants will then depend on the geometry of the Fermi surface (*i.e.* t_1 and t_2) and on the magnetic field. At one-loop level these corrections are simple to derive [20]. However, the two-loop level is more complicated and the corrections, which are not analytical, may be derived numerically. Since the main contribution to the scaling equations comes from the one-loop corrections [22], we will restrict our calculations, in the 2D regime, to the one-loop level. This will enable us to obtain the essential behavior of the renormalized parameters while handling analytical expressions. It should be noted that at the one loop level, the renormalization of the hopping parameters is, as we will show, independent of the g_i couplings. The scaling equations of the latter in the 2D regime, which we do not recall here, depend on a dimensional crossover function [20].

3.2 Renormalization group equations for the one-particle hopping processes

The diagrammatic representation, at the two-loop level, of the scaling equations of t_1 and t_2 are given in Figure 2.

In the 1D regime, we found that the corrections to the scaled hopping process t_m ($m = 1,2$) diverge as $\ln[(\omega - \delta\, m\, \omega_c)/E_0]$ where $\delta = \pm 1$. This logarithmic correction may be split in two terms [19]:

$$\ln\left(\frac{\omega - m\,\delta\,\omega_c}{E_0}\right) = \ln\left(\frac{\omega_c}{E_0}\right) + \ln\left(\frac{\omega}{\omega_c} - m\,\delta\right). \quad (3)$$

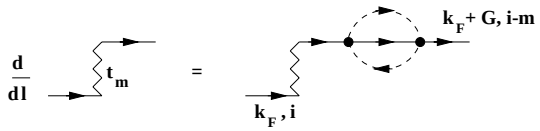

Fig. 2. Diagrammatic representation of the scaling equation of the one particle hopping processes t_m, $(m = 1, 2)$, denoted by a zigzag line. Black circle represents intrachain two particle scattering whereas the solid (broken) line corresponds to the right (left)-moving electrons. The i label is the chain index.

Since $\omega_c \ll E_0$, the problem gives rise to two logarithmic singularities, one in $\ln(\omega)$ (in the case of the two-particle sacttering) and the other in $\ln(\omega_c)$. As shown in reference [19], a double renormalization procedure with two cutoff parameters, E_0 and ω_c should be carried out.

The first step of the RG procedure corresponds to the logarithmic problem in ω_c/E_0 while ω/ω_c is kept constant. The scaling parameter ω_c/E_0 scales from 1 to its physical value. This renormalization step may be regarded as a field renormalization.

The scaling equation of the dimensionless hopping amplitude $\tilde{t}_m(l) \equiv t_m(l)/E_0$ $(m = 1, 2)$ is given by [9]:

$$\frac{\mathrm{d}\ln\tilde{t}_m(l)}{\mathrm{d}l} = 1 - \frac{1}{4}\left(g_1^2 + g_2^2 - g_1 g_2 + \frac{1}{2}g_3^2\right). \qquad (4)$$

The renormalization cutoff is $E(l) = E_0\mathrm{e}^{-l}$ where $l = 0, ..., l_{max1}$ and $l_{max1} = -\ln[\max(T,\ T_1^*,\ \omega_c)/E_0]$. l_{max1} may correspond to the thermal fluctuations (T), the 1D–2D dimensional crossover temperature (T_1^*) or the magnetic energy (ω_c).

At T_1^*, $\tilde{t}_1(l)$ reaches unity at the scaling parameter l_1^* defined by: $\tilde{t}_1(l^*) = 1$ [10]. Since the scaling parameter is identified with the temperature as $l = \ln E_0/T$ [9,10], the crossover temperature T_1^* is then given by:

$$T_1^* = E_0\mathrm{e}^{-l_1^*}.$$

For $l > l_1^*$, the scaling of t_1 becomes meaningless since t_1 cannot be regarded as a perturbation. At l_1^*, the one particle hopping t_1 becomes coherent and the system undergoes a crossover to a two-dimensional phase [10,12].

A point worth stressing here is that the energy value T_1^* at which the dimensional crossover takes place is still an open question [2]. The estimate of T_1^* of $\sim$100 K from transport experiments [23] disagrees with that deduced from early interpretations of NMR measurements [14]. In the following we will take t_1 $\sim$200 K which gives rise to $T_1^* \sim 150$ K within the PRG approach. It should be noted that, although the value of T_1^* is controversial, there is a consensus on the bare value of t_1.

The magnitude of the magnetic field used in experiments does not exceed 30 T. The magnetic energy ω_c is then lower than T_1^*. Therefore, the first step of the renormalization procedure should be stopped at l_1^* where the cutoff E_0 is scaled to T_1^*.

A second step, which is also a field renormalization, should be carried out in the two dimensional regime ($T < T_1^*$). As we have noted before, the scaling of t_1 is stopped whereas the scaling equation of t_2 reduces, at the one loop level, to [9]:

$$\frac{\mathrm{d}\ln\tilde{t}_2(l)}{\mathrm{d}l} = 1. \qquad (5)$$

The cutoff is parametrized, in this case, as $E(l) = T_1^*\mathrm{e}^{-l}$ where $l = 0, ..., l_{max2}$ and $l_{max2} = -\ln[\max(T, \omega_c)/T_1^*]$.

According to equation (5), the scaling of t_2 is independent of the coupling constants. In the 2D regime, and at the one loop level, one may, therefore, focus only on the renormalization of t_2.

The bare value of $\tilde{t}_2$ in the second step of the renormalization procedure is the renormalized value at l_1^* of the first step, *i.e.*:

$$\left[\tilde{t}_2(0)\right]_2 = \left[\tilde{t}_2(l_1^*)\right]_1$$

where the labels 1 and 2 correspond, respectively, to the first and to the second renormalization step procedure. To study the effect of the hopping parameter t_2 on the nesting properties of the Bechgaard salts, one may be interested in the low temperature regime $(T < \omega_c)$. The second step of the scaling procedure should then be stopped at l_c defined as:

$$E(l_c) \equiv T_1^*\mathrm{e}^{-l_c} = \omega_c$$

At this point, a third step of the renormalization procedure starts, where the scaling parameter ω/ω_c is varied from 1 to T/ω_c keeping ω_c/T_1^* constant. This step may be called a frequency renormalization.

Since the correction to the renormalized t_2 (Fig. 2) is proportional to $\ln(\omega/\omega_c - 2\delta)$, which is not divergent, the scaling equation of t_2 reduces then, in this step, to [19]:

$$\frac{\mathrm{d}\ln\tilde{t}_2(l)}{\mathrm{d}l} = 1. \qquad (6)$$

The scaling energy reads as $E(l) = \omega_c\mathrm{e}^{-l}$ where $l = 0, ..., l_{max3}$ and $l_{max3} = \ln[\max(T, T_2^*/\omega_c)]$. T_2^* is a sort of a crossover temperature given by:

$$T_2^* = \omega_c\mathrm{e}^{-l_2^*}$$

where l_2^* is specified as:

$$\tilde{t}_2(l_2^*) = 1.$$

At T_2^* the one particle hopping to the second nearest neighbors t_2 becomes coherent.

The initial condition of the third step of the scaling procedure is given by:

$$\left[\tilde{t}_2(0)\right]_3 = \frac{\left[\tilde{t}_2(l_c)\right]_2 \times T_2^*(H=0)}{\omega_c}$$

$T_2^*(H = 0)$ is the value of T_2^* at zero field. The different steps of the t_2 renormalization are summarized in Figure 3.

It is worth noting that to overcome the effect of the hopping parameter t_1, a magnetic field of $\sim$100 T is necessary to fulfill the condition $\omega_c > T_1^*$. The renormalization of t_1 will hence consist in a two step-procedure within

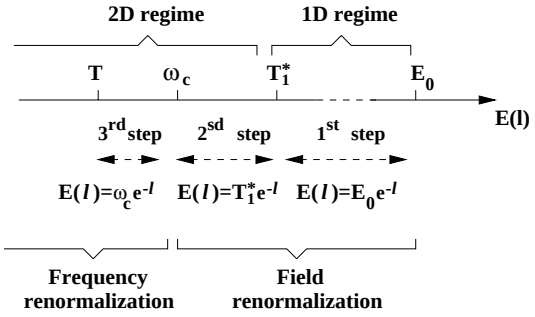

Fig. 3. Three-step renormalization procedure of the t_2 hopping parameter. The cutoff parameters are E_0, T_1^* and ω_c. $E(l)$ is the scaling energy.

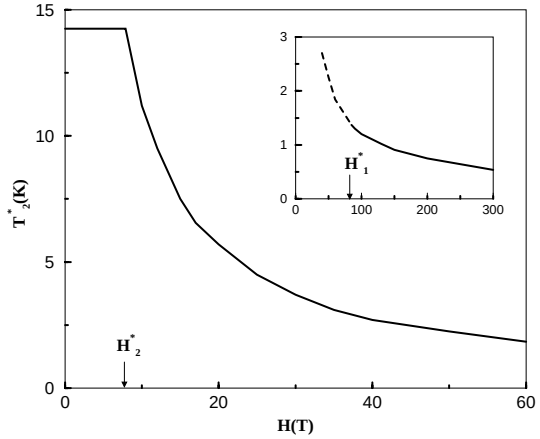

Fig. 4. Field renormalization of T_2^*. H_2^* is the critical field under which the hopping process t_2 is unaffected by the magnetic field. The inset shows the renormalization of T_2^* in the high field regime (solid line) for $\omega_c > T_1^*$. The meaning of H_1^* is given in Figure 5.

the 1D regime. The first one is carried out in the domain $[\omega_c, E_0]$ where the scaling parameter is ω_c/E_0 whereas the second one deals with the domain $T < \omega_c$ with ω/ω_c as a scaling parameter.

The scaling equation of t_1 in the first step of the renormalization procedure is given by equation (4). The bare value of t_1 in the second step is given by:

$$\left[\tilde{t}_1(0)\right]_2 = \frac{\left[\tilde{t}_1(l_c)\right]_1 \times T_1^*(H=0)}{\omega_c}$$

where $T_1^*(H=0)$ is the 1D-2D crossover temperature at zero field and l_c is defined as $l_c = \ln E_0/\omega_c$. As we have discussed in the case of the t_2 parameter, the scaling equation of t_1 during the frequency renormalization step (second step) is reduced to:

$$\frac{\mathrm{d}\ln \tilde{t}_1(l)}{\mathrm{d}l} = 1. \tag{7}$$

This second step is stopped at l_H^* defined by:

$$\tilde{t}_1(l_H^*) = 1.$$

The renormalized 1D-2D crossover temperature $T_1^*(H)$ in the presence of the magnetic field is then written as:

$$T_1^*(H) = \omega_c \mathrm{e}^{-l_H^*}.$$

One should expect, as in the one-dimensionalization picture, a decrease of $T_1^*(H)$ due to the enhancement of the 1D character with increasing magnetic field. The 1D regime will therefore expand at the expense of the 2D regime.

Since we are interested in the confinement effect of the magnetic field, we will, therefore, focus on the relevant parameters T_1^* and T_2^* at which the hopping processes t_1 and t_2 become coherent. The effective hopping value of t_1 and t_2, which could not be derived directly from the present PRG approach, may be determined from T_1^* and T_2^* [9].

4 Results and discussion

We have considered, as in reference [10], that the intrachain scattering strengths depend not only on the on-site coulomb repulsion but also on the nearest-neighbors one. For the first step of the renormalization procedure, we take as initial condition at $l = 0$ [10]:

$$\pi v_F g_1(0) = \pi v_F g_3(0) = 0.2 \text{ and } \pi v_F g_2(0) = 0.6$$

The bare values of the hopping processes are $t_1(0)$ =250 K and $t_2(0)$ =25 K, which are consistent with the experimental values. The magnetic energy is given as ω_c/H = 1.8 K/T [24], whereas the bandwidth is taken as E_0 = 3000 K.

Let us consider the case where $\omega_c \ll T_1^*$. Carrying out the first renormalization step procedure, $\tilde{t}_1(l)$ reaches unity at $T_1^* = E_0 \mathrm{e}^{-l_1^*}$ =143 K. During the second step of the scaling procedure, one may meet two situations:
(a) At $\omega_c < E(l) < T_1^*$, $\tilde{t}_2(l)$ reaches unity at l_2^*, which is the same as that obtained at zero field. The t_2 hopping process is then not affected by the magnetic field and its scaling should be stopped at this stage.
(b) At $\omega_c < E(l) < T_1^*$, $\tilde{t}_2(l)$ is smaller than unity. In this case, the third step of the renormalization should be carried out for $T < E(l) < \omega_c$. This step will be stopped when $\tilde{t}_2(l)$ reaches unity, which determine the temperature T_2^*.

At zero field, $\tilde{t}_2(l)$ attains an order of unity around the crossover temperature T_2^*=14 K. The ratio T_1^*/T_2^* = $t_1(0)/t_2(0)$ =10 is equal to that of the effective hopping parameters t_1 and t_2 obtained experimentally.

Carrying out the different steps of the renormalization procedure in the presence of a transverse magnetic field, we have obtained the dependence of the crossover temperature T_2^* on the magnetic field H as depicted in Figure 4. The latter shows that below a critical value H_2^*, T_2^* is field independent. The critical value H_2^* corresponds to the limit value $\omega_c\,(H_2^*) = T_2^*(H=0)$ below which the magnetic field cannot bypass the hopping process t_2. The experimental value of t_2, and therefore its bare value $t_2(0)$ (and, of course $t_1(0)$) is strongly dependent on the applied pressure. Since it considerably increases with pressure, we

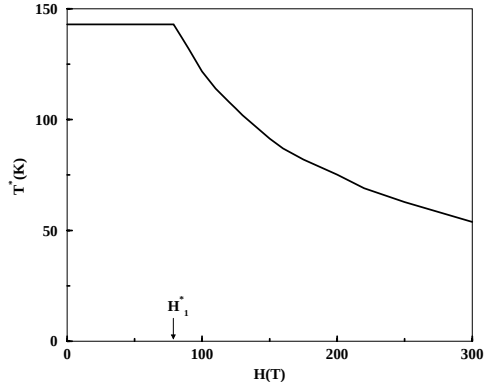

Fig. 5. Field renormalization of T_1^*. H_1^* is the critical field under which the hopping process t_1 is unaffected by the magnetic field.

can conclude that H_2^* also increases with increasing pressure. Therefore, the unidimensionalization effect will require, as expected, an increasing magnitude of the applied field as the pressure increases. For $H > H_2^*$, T_2^* decreases with increasing field as a power law. At H =30 T, T_2^* is reduced by a factor of 4 and for the high field regime T_2^* tends to a limit of ∼1 K. The effective parameter t_2 is then reduced as the field increases and the nesting properties are improved.

Let us now focus on the case where $\omega_c > T_1^*$. The first step of the renormalization procedure will be stopped at $l_c = \ln E_0/\omega_c$ where $\tilde{t}_1(l_c) < 1$. To determine the renormalized T_1^*, a second step should be carried out.

At a scaling parameter l_H^*, $\tilde{t}_1(l)$ reaches unity and the system undergoes a dimensional crossover at $T_1^* = \omega_c \mathrm{e}^{-l_H^*}$. The filed dependence of this crossover temperature is shown in Figure 5.

Figure 5 shows that T_1^* is considerably reduced for $H > H_1^*$ at which $\omega_c = T_1^*(H = 0)$. The effective dimensionality of the system is then decreased towards the 1D case. This effect may be called a magnetic-field-induced Luttinger liquid by analogy with the the so-called magnetic-field-induced Luttinger insulator [24]. The 1D regime will expand at the expense of the 2D phase, which enhances the anisotropy of the system. Such behavior may be observed in the transport measurements. In reference [25] an experimental evidence for a magnetic field induced one-dimensionalization effect has been provided based on magnetoresistance measurements in $(TMTSF)_2PF_6$. However, in reference [26], the authors reported, from a magnetoresistance study, that the anisotropy of the system is field independent, on the contrary to other experimental and theoretical studies [24,25,27]. The origin of this controversy may lie in the pressure regime at which were carried out the measurements. In reference [25], the experiments were done above the critical pressure P_c at which the SDW phase (phase AF in Fig. 1) collapses, namely in the metallic regime. Nevertheless, in reference [26], the magnetoresistance is studied under a pressure $P < P_c$, corresponding to the ordered SDW state. In this latter, and unlike the metallic regime, it may be more difficult to study the competition between the filed confinement effect and the interchain hopping process due to the presence of strong AF correlations.

We should note that the critical field H_1^* is high compared to the experimental values [25]. Within the present PRG approach we are not able to seek the field renormalization of t_1 in the 2D regime, since in this regime, the perturbative approach is no longer valid. However, the t_1 hopping process may be confined in the 2D phase, for $H < H_1^*$ as it is expected within the one-dimensionalization scenario [24,25,27,28].

To derive the field renormalization of t_2 for $\omega_c > T_1^*$, one should carry out a third scaling step in the 2D regime with the initial condition:

$$\left[\tilde{t}_2(0)\right]_3 = \left[\tilde{t}_2(l_H^*)\right]_2 .$$

The scaling equation is still given by equation (6). This renormalization procedure will be stopped at the scaling parameter at which $\tilde{t}_2(l)$ reaches unity. The inset of Figure 4 shows that the T_2^* line extends in the high field regime ($H > H_1^*$) (solid line) and continues to decrease smoothly. Therefore, the 1D renormalization procedure (for $\omega_c > T_1^*$) (solid line in the inset of Fig. 4) and the 2D one (for $\omega_c < T_1^*$) (dashed line in the inset of Fig. 4) are complementary and coherent.

We now discuss qualitatively the effect of the field renormalization of the imperfect nesting parameter t_2 on the temperature-field phase diagram of the Bechgaard salts. The field renormalization of t_2 should be taken into account when deriving, within the QNM, the transition temperatures of the FISDW phases. The physics of the QNM remains qualitatively valid, whereas the quantitative behavior may be affected. The field reduced effective t_2 may induce a shift of the FISDW cascade to lower fields.

Within a subphase of a quantum number $N \neq 0$, and at a fixed field, the transition temperature is expected to decrease as t_2 decreases. Indeed, a decrease of t_2 yields to a decrease of the well known coefficients I_N ($N \neq 0$) [5], which reduces the transition temperature. This behavior is consistent with the dependence on pressure of the transition temperature T_N^c of the $N \neq 0$ FISDW phases obtained experimentally in the case of $(TMTSF)_2PF_6$ compound [27]. T_N^c are found to decrease with decreasing pressure, which is equivalent to a decreasing t_2.

Concerning the $N = 0$ phase, as t_2 is lowered, the I_0 coefficient and then the transition temperature are raised, which furthers the $N = 0$ phase. This feature is in agreement with the increase of the metal-$N = 0$ phase transition temperature with decreasing pressure as found in the experimental phase diagram of the $(TMTSF)_2PF_6$ [27].

5 Conclusion

In this paper, we have discussed the effect of a transverse magnetic field on the $(TM)_2X$ systems. Using a perturbative renormalization group (PRG) approach with three

cut-off parameters, the bandwidth E_0, the magnetic energy ω_c and the 1D-2D crossover temperature T_1^*, we found that the interchain hopping parameters t_1 and t_2 are renormalized by applying a magnetic field in the **c** direction.

Above a critical field H_1^*, the crossover temperature T_1^* is strongly reduced with increasing field, which decreases the effective dimensionality of the system towards the 1D state. This result is in agreement with the one-dimensionalization picture and with recent experimental studies.

In the 2D phase, we have shown that for $H > H_2^*$, the crossover temperature T_2^*, at which the t_2 process becomes coherent, decreases as the field increases, which improve the nesting properties. We have argued that the field renormalization of t_2 should be taken into account when deriving the high field part of the FISDW cascade (*i.e.* for $H > H_2^*$) in quasi-one dimensional organic conductors.

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