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**Magnetic oscillations and field-induced spin-density waves in (TMTSF)<sub>2</sub>ClO<sub>4</sub>**

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We report an analysis of the effects of magnetic field on a quasi-one-dimensional band of interacting electrons with a transverse dimerizing potential. One-particle problem in bond-antibond representation is solved exactly. The resulting propagator is used to calculate the spin-density-wave (SDW) response of the interacting system within the matrix random-phase approximation for the SDW susceptibility. We find that the value of the anion potential fitting experiments in relaxed (TMTSF)<sub>2</sub>ClO<sub>4</sub> is large, of the order of interchain hopping. In particular we predict the magnetic-field-induced transition of the first order between interband SDW<sub>0</sub> and intraband SDW<sub>±</sub> phases, we reproduce the rapid oscillations with a period of 260 T and the overall profile of the (TMTSF)<sub>2</sub>ClO<sub>4</sub> phase diagram.

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**I. INTRODUCTION**

Investigations of quasi-one-dimensional electronic systems at high magnetic fields and at low temperatures continue to give an important insight into the one-particle properties and interaction-induced phases such as spin- and charge-density-wave, superconductivity, and Mott localization.<sup>1</sup> One of most spectacular phases of this kind are field-induced spin-density wave (FISDW), found in Bechgaard salts<sup>2</sup> and in some other low-dimensional compounds.<sup>3</sup> The phenomenon of the FISDW is well understood in the Bechgaard salt (TMTSF)<sub>2</sub>PF<sub>6</sub> where the cascade of SDW phases with quantized wave vector is induced by orbital effects of magnetic field to the quasi-one-dimensional (Q1D) orbits of band electrons. Theory based on the mechanism of quantized nesting<sup>4</sup> reproduces satisfactorily main experimental data for this salt.

In this paper we concentrate on (TMTSF)<sub>2</sub>ClO<sub>4</sub>, a Bechgaard salt which after a slow cooling<sup>5,6</sup> enters into a qualitatively different type of FISDW phase at low temperatures, with a phase diagram that is still, after more than 10 years of intensive studies,<sup>1,2</sup> a matter of both experimental and theoretic controversies. In particular for magnetic field  $B > 8$  T the nature of the ordering in the relaxed material is not a simple FISDW with some low integer quantum number  $N$ , but a qualitatively different state containing several puzzling subphases.<sup>2,7,8</sup> This phase is at 8 T separated by a line of first-order transition from a cascade of FISDW phases which very much resembles that of the standard model. Another characteristic phenomenon, the rapid oscillations (RO) in  $1/B$  with a frequency of 260 T, is visible in transport properties in both metallic and FISDW state.<sup>2,8,9</sup> Similar RO are seen also in thermodynamic quantities such as torque, magnetization, sound velocity, and specific heat, but only in the ordered phase.<sup>1,2</sup> The highest value of  $T_c$  in the  $T_c(B)$  dependence is 5.5 K, instead of 12 K as expected from analogy with the (TMTSF)<sub>2</sub>PF<sub>6</sub> salt.

The incompatibility of above facts with the quantum nesting model (QNM) for a single quasi-1D band is believed to

stem from the particular ordering of ClO<sub>4</sub> anions.<sup>2,10</sup> This ordering introduces the new modulation with the wave vector  $(0, \pi/b, 0)$ , i.e., a dimerization in the low-conducting direction with the interchain distance  $b$ . The magnitude of the dimerizing potential can be tuned to some extent by varying the cooling rate.<sup>5,6</sup> Thus, anions presumably remain disordered in the rapidly quenched samples. Then there is no dimerization gap in the band, and the system shows properties of a *single* quasi-1D imperfectly nested band with a SDW order appearing already in the zero magnetic field.<sup>5,6,11</sup> The anion ordering in slowly relaxed samples is at about 24 K, and coincides with the onset of rapid oscillations in the magnetoresistance.<sup>12</sup> The RO in (TMTSF)<sub>2</sub>ClO<sub>4</sub> have been theoretically explained in two limiting cases. The limit of strong anion potential  $V \gg t_b$ ,  $t_b$  being the interchain hopping integral, was calculated by Brazovskii and Yakovenko,<sup>10</sup> while the opposite limit  $V \ll t_b$  was solved by Lebed and Bak.<sup>13</sup> In this paper we solve *exactly*, i.e., for any  $V$  and  $B$ , the one-particle problem, which determines the RO phenomenon.

The dimerized band has two pairs of Fermi sheets in the new Brillouin zone, as shown in Fig. 1. Already simple geometric arguments<sup>7</sup> suggest three possible nesting wave vectors favoring various SDW phases. The interband nesting  $\mathbf{Q}_0$  leads to SDW<sub>0</sub> that is the two-band version of the standard FISDW phase. Other two nesting vectors  $\mathbf{Q}_+$  and  $\mathbf{Q}_-$  relate Fermi sheets within the same band. They give SDW<sub>+</sub> for bond nesting and SDW<sub>-</sub> for antibond nesting. However the interplay between SDW<sub>0</sub> and SDW<sub>±</sub> is not only a geometric question of the choice of the nesting vector. Due to a finite anion potential  $V$  in the kinetic part of the Hamiltonian an off-diagonal term appears in the SDW response, making necessary an appropriate matrix approach<sup>14,15</sup> in the calculation of the critical susceptibilities. This matrix aspect of the problem was ignored in all former theoretical approaches.<sup>16–19,21</sup> We formulate the response matrix in the space of two order parameters  $\Delta_h$  (homogeneous) and  $\Delta_a$  (alternating) determining the magnetic pattern,

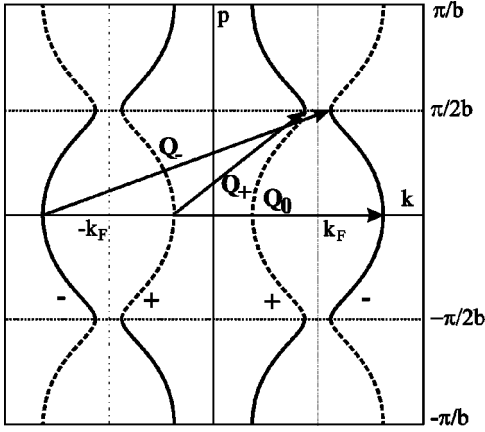


FIG. 1. Two Fermi surfaces of  $(\text{TMTSF})_2\text{ClO}_4$  (wrapping is highly exaggerated). Wave vectors  $\mathbf{Q}_0$ ,  $\mathbf{Q}_+$ , and  $\mathbf{Q}_-$  correspond, respectively, to  $\text{SDW}_0$ ,  $\text{SDW}_+$ , and  $\text{SDW}_-$ .

$$m_z(x, R_\perp) = (\Delta_h \pm \Delta_a) \cos[(2k_F + k)x + pnd]. \quad (1)$$

Here  $d \equiv 2b$  and the upper and lower sign stay for even ( $R_\perp = nd$ ) and odd ( $R_\perp = nd + d/2$ ) chains, respectively. As it is shown in Refs. 14,15, three types of SDW modulations with wave vectors shown in Fig. 1 are candidates for ordering at the phase transition from the metallic state.  $\text{SDW}_0$  is stabilized for low values of  $V$  (providing the imperfect nesting parameter  $t'_b$  allows for SDW stabilization), while  $\text{SDW}_+$  or  $\text{SDW}_-$  get stable for  $V/t_b > 1.6$  irrespective of the value of  $t'_b$ . The slowly relaxed  $(\text{TMTSF})_2\text{ClO}_4$  samples are expected to lie in the range of intermediate values of  $V$  in which there is no SDW ordering at  $B=0$  down to  $T=0$ .

Indeed, as it will be shown below,  $V/t_b$  fitting the experiments is close to unity, which is also in agreement with recent detailed x-ray data.<sup>20</sup> Still, Haddad *et al.*<sup>21</sup> recently put forward some arguments in favor of the small  $V$ . In order to construct the phase diagram with dominant  $\text{SDW}_\pm$  already in the range of small  $V$  these authors put larger coupling constant for  $\text{SDW}_\pm$  than for  $\text{SDW}_0$ . This assertion is not plausible because the renormalization group for the quasi-one-dimensional interacting fermions suggests that one obtains difference between intraband and interband effective coupling constants only if  $V$  is of the order or larger than  $t_b$ .<sup>22,23</sup>

In the range  $V/t_b \geq 1$  it is not allowed to use the quasi-classical approximation of Gor'kov and Lebed,<sup>17</sup> which consists in making Peierls substitution  $p \rightarrow p - eA$  in each subband separately and including the anions' effects only via magnetic breakdown (MB) junctions near the zone boundary. While this approximation is sufficient for  $V/t_b \leq 1$ , here one has to solve the whole quantum-mechanical problem instead.

It was pointed out several times<sup>8,12,17</sup> that a mechanism of coherent interband tunneling, very similar to Stark overgap quantum interference (QI) in magnesium,<sup>24</sup> is essential for high-field physics in  $(\text{TMTSF})_2\text{ClO}_4$ . In particular, RO in metallic state can be explained only in terms of QI mechanism because no closed orbits exist. On the contrary, in the SDW state both closed orbits *and* Stark interference contribute to RO. Oscillating behavior periodic in  $1/B$  can be seen already at the level of one-particle spectrum. This is the topic

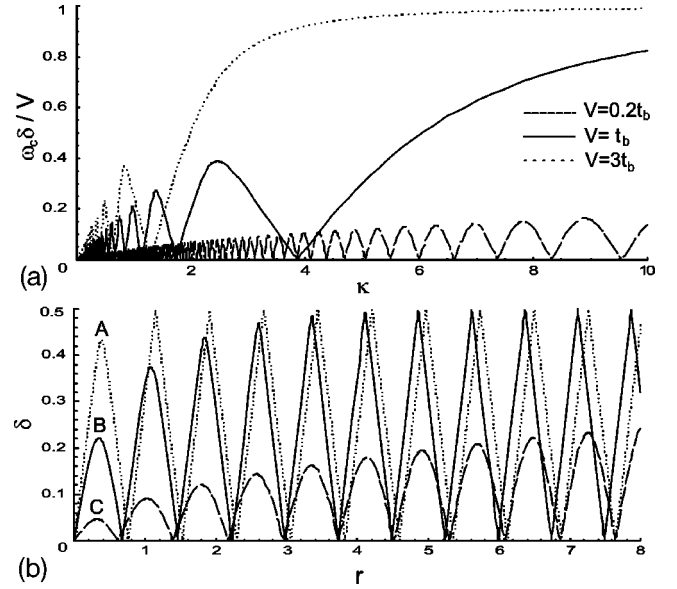


FIG. 2. (a) Energy ratio  $\omega_c \delta / V$  as a function of the magnetic breakdown parameter  $\kappa$  for several values of  $V/t_b$ . (b) Dependence of  $\delta$  on  $r$  for  $\theta = 10^\circ$  (A),  $45^\circ$  (B), and  $80^\circ$  (C).

of the following section. In Sec. III we include interactions via the matrix random-phase approximation (RPA) for the *two-component* SDW order parameter and construct the phase diagram. The last section contains conclusions.

## II. EXACT SOLUTION OF THE STARK INTERFEROMETER

We solved the one-particle problem of the Q1D band with anion potential  $V$ . Resulting electronic propagator with longitudinal momentum  $k$  has poles at

$$E_f = v_F [f(k - k_F) + GN] \pm v_F G \delta, \quad (2)$$

where  $f$  is left-right index,  $G \equiv eBb/\hbar$  is the magnetic wave number, and  $N$  is integer number. The first term in Eq. (2) is the standard QNM dispersion and the last term is the splitting due to anions. Overgap resonances are present in  $\delta(B)$  as will be discussed below (see Fig. 2). The expression for the spectrum (2) is common to perturbation calculations,<sup>16</sup> to quasiclassical tunnelling analysis,<sup>17</sup> and to our exact solution as well. What change from one approach to another are the dependence  $\delta(B)$  and the result for electronic wave function. In order to obtain them exactly we start from the effective one-particle Hamiltonian for electronic operators  $\Psi_f(x, p)$ ,

$$H_0 = iv_F \rho_3 \partial_x + \tau_3 \mathcal{T}(pb - Gx) + \bar{\mathcal{T}}(pb - Gx) - V\tau_1, \quad (3)$$

where  $\rho$ 's and  $\tau$ 's are Pauli matrices in left-right and bond-antibond indices, respectively. The most general transverse dispersion was split into two parts,

$$\mathcal{T}(pb) \equiv 2 \sum_{j=1}^{\infty} t_j \cos[(2j-1)pb], \quad (4)$$

$$\tilde{T}(pb) \equiv 2 \sum_{j=1}^{\infty} t'_j \cos[2jpb],$$

corresponding to effective hoppings<sup>25</sup> between odd and even neighbors, respectively. We diagonalize  $H_0$  by the unitary transform

$$\Psi_f = \begin{pmatrix} \alpha_f & \beta_f \\ -\beta_f^* & \alpha_f^* \end{pmatrix} e^{if\theta} \Phi_f, \quad (5)$$

with  $|\alpha|^2 + |\beta|^2 = 1$ , and functions  $\alpha$ ,  $\beta$ , and  $\theta$  depending on  $x$  and  $p$  only through the combination  $z = pb - Gx$ . From the requirement that the effective Hamiltonian for field  $\Phi$  be only  $ifv_F\partial_x$  we get  $\theta(z) = v_F^{-1} \int^z dz' \tilde{T}(z')$  and a system of differential equations for functions  $\alpha$  and  $\beta$ ,

$$\begin{aligned} ifv_F\alpha'_f(z) &= -\mathcal{T}(z)\alpha_f(z) - V\beta_f^*(z), \\ ifv_F\beta'_f(z) &= -\mathcal{T}(z)\beta_f(z) + V\alpha_f^*(z). \end{aligned} \quad (6)$$

This ‘‘one dimensionalization’’ of the effective Hamiltonian is the two-component generalization of the standard phase transformation procedure for Q1D systems in magnetic field.<sup>4</sup> Note that  $\theta(z+2\pi) = \theta(z)$  and that  $\alpha_+(z) = \alpha_-^*(z)$  and  $\beta_+(z) = \beta_-^*(z)$ , so that it suffices to follow, e.g., solutions  $\alpha_+(z), \beta_+(z)$  of the system (6). According to Floquet theory these solutions can be written in the form  $\alpha(z) = A(z)\exp(-iz\delta)$ ;  $\beta(z) = B(z)\exp(iz\delta)$ .  $A$  and  $B$  are periodic with the period  $2\pi$ , and closer inspection shows that the Floquet exponent  $\delta$  for the system (6) is real for all values of parameters, at least after keeping in  $\mathcal{T}(z)$  only the leading term  $j=1$ .

The Floquet exponent  $\delta$  and the functions  $A$  and  $B$  are calculated using the Hill's theory and the fundamental matrix method.<sup>26</sup> In the present work we limit our calculations only to first harmonics in Eq. (4), parametrized with  $t_1 = t_b$  and  $t'_1 = t'_b$ . Let us concentrate on the magnetic-field dependence of the Floquet exponent  $\delta$  that splits the QNM spectrum as given by Eq. (2). Figure 2(a) shows the energy  $\omega_c\delta$  (in units of  $V$ ) as a function of the magnetic breakdown parameter  $\kappa \equiv 2\omega_c t_b / V^2$ , where  $\omega_c = v_F G$  is the cyclotron frequency. In quasiclassical picture  $\kappa$  determines the probability of the overgap tunneling  $P = \exp(-\pi/2\kappa)$ .<sup>17</sup> One sees that the crossover from oscillating to saturating behavior does not coincide with the crossover from the weak ( $\kappa < 1$ ) to the strong ( $\kappa > 1$ ) MB. The position of the last zero of  $\delta$  is not universal in  $\kappa$ , but approximately in  $r \equiv [(\gamma V)^2 + t_b^2]^{1/2} / \omega_c$ , where the value of  $\gamma$  is 0.77. Figure 2(b) shows  $\delta(r)$  for several ‘‘polar angles’’ defined by  $\tan \theta \equiv t_b / \gamma V$ . Oscillations of  $\delta$  are approximately periodic in  $r$  with a period of 0.80. Choosing the parameters  $t_b = 300$  K,  $v_F = 2 \times 10^5$  m/s, and  $b = 7.7 \times 10^{-10}$  m we fit RO at 260 T by putting  $V \approx 0.8t_b$ .

Taking the limit of strong magnetic field  $\omega_c/t_b \gg 1$  and of weak anion potential  $V/t_b \ll 1$  we can easily reproduce the 1D spectrum of Osada *et al.*,<sup>16</sup>  $E_k \rightarrow f v_F(k - k_f) \pm \omega_c\delta$  with  $\delta \rightarrow (V/\omega_c) \mathcal{J}_0(4t_b/\omega_c)$ ,  $\mathcal{J}_0$  being the Bessel function. On the other hand the spectrum of Gor'kov and Lebed<sup>17</sup> is reproduced for weak anion potential,  $V/t_b \ll 1$ . The above fit,

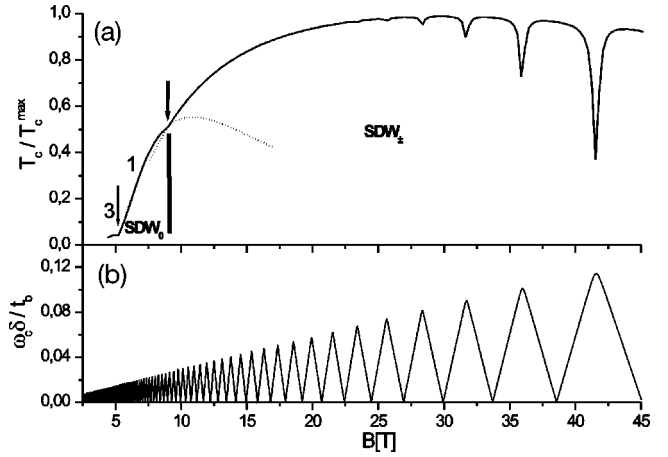


FIG. 3. (a) Phase diagram. (b) Energy ratio  $\omega_c\delta/t_b$  on the same magnetic scale as the phase diagram.

as well as other insights<sup>20,28</sup> however strongly suggest that  $V$  in  $(\text{TMTSF})_2\text{ClO}_4$  is rather large, i.e., comparable to  $t_b$ .

The rapid oscillations in observable response functions are related to the oscillations of  $\delta$  (Ref. 17), shown in Figs. 2 and 3(b). At 30 T the magnetic breakdown parameter has moderate value of  $\kappa \sim 0.5$ .

### III. MANY-BODY EFFECTS

We proceed with the solution of the interacting problem. Neglecting the absence of a presumably small umklapp scattering, the effective coupling for SDW is the forward-scattering amplitude  $g_2$ , here simply denoted by  $U$ . We employ the matrix RPA formalism developed in Ref. 14. The resulting relevant bare susceptibility is  $\chi_1(\mathbf{q}; T) = \frac{1}{2} \{ \chi_{aa} + \chi_{hh} + [(\chi_{aa} - \chi_{hh})^2 + 4(\chi_{ha})^2]^{1/2} \}$ , entering into the Stoner criterion

$$1 - U\chi_1(\mathbf{q}_c, T_c) = 0, \quad (7)$$

$\mathbf{q}_c$  being the wave vector at which  $\chi_1(\mathbf{q})$  has the maximum. The ratio of two SDW order parameters from Eq. (1) is also a function of bare correlators  $\chi_{aa}, \chi_{hh}, \chi_{ah}$  in the  $(a, h)$  basis (see Ref. 14). Their analytical expressions are

$$\begin{aligned} \chi_{hh} &= \sum_N \left[ \left| I_{h0} \right|^2 P_0 + \frac{1}{2} I_{h+}^2 P_+ + \frac{1}{2} I_{h-}^2 P_- \right], \\ \chi_{aa} &= \sum_N \left[ \left| I_{a0} \right|^2 P_0 + \frac{1}{2} I_{a+}^2 P_+ + \frac{1}{2} I_{a-}^2 P_- \right], \\ \chi_{ha} &= \sum_N \left[ \text{Re}(I_{h0} I_{a0}^*) P_0 + \frac{1}{2} I_{h+} I_{a+} P_+ - \frac{1}{2} I_{h-} I_{a-} P_- \right], \end{aligned} \quad (8)$$

where  $P_0, P_{\pm}$  stand for  $P(q_{\parallel} - NG, T)$  and  $P[q_{\parallel} - G(N \pm 2\delta), T]$ , respectively,  $P(k, T)$  being the familiar 1D Lindhard function at the wave number  $2k_F + k$ .  $P_0$  and  $P_{\pm}$  are the interband and the intraband susceptibilities of the  $N$ th split level of the spectrum (2). Coefficients  $a_N, b_N, \hat{a}_N$ , and  $\hat{b}_N$  are Fourier components of the products  $A \exp(i\theta)$ ,

$B \exp(i\theta)$ ,  $A^* \exp(i\theta)$ , and  $B^* \exp(i\theta)$ , respectively. The dependence on the transverse momentum is present in the amplitudes  $I(q_\perp, N)$ ,

$$\begin{aligned}
 I_{h_0}(q_\perp, N) &= \sum_n (a_n b_{N-n} - \hat{b}_n \hat{a}_{N-n}) e^{i(n-N/2)q_\perp}, \\
 I_{h_+}(q_\perp, N) &= \sum_n (\hat{a}_n \hat{a}_{N-n} + b_n b_{N-n}) e^{i(n-N/2)q_\perp}, \\
 I_{h_-}(q_\perp, N) &= \sum_n (a_n a_{N-n} + \hat{b}_n \hat{b}_{N-n}) e^{i(n-N/2)q_\perp}, \\
 I_{a_0}(q_\perp, N) &= \sum_n (a_n \hat{a}_{N-n} - \hat{b}_n b_{N-n}) e^{i(n-N/2)q_\perp}, \\
 I_{a_+}(q_\perp, N) &= \sum_n (\hat{a}_n b_{N-n} + b_n \hat{a}_{N-n}) e^{i(n-N/2)q_\perp}, \\
 I_{a_-}(q_\perp, N) &= \sum_n (a_n \hat{b}_{N-n} + \hat{b}_n a_{N-n}) e^{i(n-N/2)q_\perp}. \quad (9)
 \end{aligned}$$

There are two important selection rules for these amplitudes, namely, for  $N$  even,  $I_{h_0}(N) = I_{a_0}(N) = 0$  while for  $N$  odd,  $I_{h_\pm}(N) = I_{a_\pm}(N) = 0$ . Thus the interband processes contribute only to FISDW phases with odd  $N$  while the intraband processes contribute only to phases with even  $N$ . Consequently only phases with even  $N$  “see” the splitting by  $\delta$ .

According to Eq. (7) the maximum of  $\chi_1(\mathbf{q})$  attains the value  $1/U$  at  $T = T_c$ . Figure 3 shows the resulting phase diagram for a realistic choice of parameters,  $V = 0.85t_b$ ,  $t'_b = 0.03t_b$ , and  $T_c(V = t'_b = 0) = 13$  K. The resulting maximal critical temperature within the present field range is  $T_c^{\max} \approx 1.1$  K. The most obvious characteristic of the obtained phase diagram is the first-order transition from  $\text{SDW}_0$  to  $\text{SDW}_\pm$  at  $B_c \approx 9$  T. Dependence  $T_c(B)$  for  $B < B_c$  is similar to the FISDW cascade in  $\text{TMTSF}_2\text{PF}_6$ , with the difference that here only odd phases appear because the even ones are suppressed by splitting. We expect that at lower temperatures the first-order transition from  $\text{SDW}_0$  to  $\text{SDW}_\pm$  is driven by stabilization of soliton lattices with competing  $\text{SDW}_0$  and  $\text{SDW}_\pm$  domains.<sup>27</sup> For  $B > B_c$  the critical temperature increases towards the highest value  $T_c^{\max}$ . As the magnetic field further increases the critical temperature  $T_c(B)$  starts to oscillate, with the sharp dips corresponding to commensurability condition  $2G\delta = G$  between the Floquet wave number and the magnetic wave number. We can also estimate the quantum Hall effect in the phase  $\text{SDW}_+$  or  $\text{SDW}_-$ . The shift from the perfect nesting in these phases is  $v_F \Delta k = \sqrt{V^2 + 2t_b^2} - (\sqrt{V^2 + 4t_b^2} + V)/2$ .<sup>14</sup> For  $V \sim t_b$  this gives  $v_F \Delta k \sim t_b/10$  and the quantum number of the Hall effect,  $N_H \sim v_F \Delta k / \omega_c$ , takes values between 3 and 1 for magnetic fields between 10 and 30 T. However, the precise values of  $N_H$  and whether  $N_H$  is integer or not are the questions beyond the present analysis.

The result of the subtle interplay between two scales  $V$  and  $t'_b$  is that the realistic profile of the phase diagram is possible only within a rather restricted range of the  $(V, t'_b)$  space. We have calculated<sup>29</sup> the phase diagram for all values of  $V$  and concluded that  $V \approx 0.85t_b$  is indeed the only value fitting the phase diagram obtained in experiments. Namely, increasing  $V$  or  $t'_b$  by a few percent one reduces  $T_c(\text{SDW}_0)$  below  $T_c(\text{SDW}_\pm)$  in the whole  $B$  domain. On the other hand by decreasing  $V$  by a few percents one gets a hump in  $T_c(\text{SDW}_0)$  on the left of the transition  $\text{SDW}_0 - \text{SDW}_\pm$ .

The maximal value of the critical temperature in Fig. 3,  $T_c^{\max} \approx 1.1$  K, is considerably smaller than the experimental value of 5.5 K. In this respect we note that  $T_c^{\max}$  is essentially model dependent quantity, i.e., the Hamiltonian (3) represents a *minimal* model for understanding the interplay between two SDW phases in the magnetic field. Namely, recent experiments<sup>20</sup> suggest that the anion ordering in  $\text{TMTSF}_2\text{ClO}_4$  induces also, beside a strong dimerizing potential  $V$ , rather large changes in other band parameters.

The present treatment also does not include the quantitative analysis of the splitting of degeneracy of two intraband phases  $\text{SDW}_+$  and  $\text{SDW}_-$ . Physically the degeneracy is lifted because the realistic tight-binding dispersion along the chain is not strictly linear. Consequently the dominant instability will be that of  $\text{SDW}_-$ , as discussed in Ref. 14. Similar conclusions were obtained also by numerical calculations,<sup>19</sup> but without taking into account the two-component aspect of the order parameter (1). The critical temperature for the  $\text{SDW}_+$  subphase can be calculated within Landau theory as in Ref. 15, and by taking the nonlinearity of the band dispersion into account. The subphases of the high-field phase correspond to  $\text{SDW}_+$  phases within  $\text{SDW}_-$ , each one nesting its own pair of Fermi sheets. Such scenario is impossible for  $\text{SDW}_0$  since it proceeds through nesting of all four sheets at the *single* critical temperature. On this point our picture differs again from the one advanced in Ref. 21, where it was argued that  $\text{SDW}_+$  and  $\text{SDW}_-$  *must* order simultaneously because otherwise  $T_c$  would disappear exponentially. As far as we see this kind of locking of the two critical temperatures is not possible. The splitting of the single  $T_c$  to  $T_c(\text{SDW}_+)$  and  $T_c(\text{SDW}_-)$  is a smooth function of the appropriate band parameters, the simplest one being the effective third-neighbor interchain hopping  $t_3$ .<sup>15</sup>

#### IV. CONCLUSION

We solved exactly the one-particle problem of dimerized Q1D band of electrons in magnetic field. Observables contain characteristic periodicity in  $1/B$ , consistent with 260 T oscillations in normal and SDW phases of  $(\text{TMTSF})_2\text{ClO}_4$ . Using matrix RPA for SDW susceptibility we reproduce the overall profile of the experimental phase diagram, containing the first-order transition from the (low-field) interband  $\text{SDW}_0$  to the (high-field) intraband  $\text{SDW}_-$  (or  $\text{SDW}_+$ ). The value of the anion potential  $V$  fitting experiments in relaxed  $(\text{TMTSF})_2\text{ClO}_4$  is large, of the order of interchain hopping  $t_b$ .

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