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Abstract. — We argue that magnetic and angular oscillations observed in quasi-one-dimen-
sional (Q1D) organic conductors represent a new type of many-body phenomena. The physical
nature of such effects as “magic angles” in (TMTSF)$_2$ClO$_4$, (TMTSF)$_2$PF$_6$, and (DMET-TSeF)$_2$
AuCl$_2$ as well as “rapid magnetic oscillations” in (TMTSF)$_2$ClO$_4$ is shown to be beyond the
standard theory of metals. Below we discuss an explanation of these phenomena which utilizes
unusual many-body effect — a change of an effective dimensionality of electron-electron (“e-e”)
interactions with changing both a magnitude and a direction of a magnetic field. We show that
some exotic transport properties of a metallic state can be interpreted in terms of these dimen-
sional crossovers. We also demonstrate that magnetic field dependence of “e-e” interactions has
to break Fermi liquid description of quasiparticles at high magnetic fields, $H \geq 25 - 30$ T. This
leads to the appearance of strong forbidden oscillations of magnetic susceptibility, $\delta x/\chi_0 \sim 1-10$,
and magnetic moment, $\delta M/M_0 \sim 0.1$. All of the above mentioned unique properties of a metallic
phase in (TMTSF)$_2$X and (DMET-TSeF)AuCl$_2$ allow us to call it an anomalous metallic phase.

1. Introduction

Numerous families of organic conductors which were synthesized during last two decades
demonstrate a wide variety of properties in a magnetic field (for a review, see [1]). Some of
them, quasi-two-dimensional (Q2D) compounds (ET)$_2$X $(X = I_3, IBr_2, (ET)$_2$Mg(SeCN)$_4$,
(ET)$_2$Mg(SeCN)$_4$ $(M = NH_4, K, Rb, Tl)$, and (ET)$_2$Cu(NCS)$_2$, contain closed quasiparticle
orbits in their electron spectra. These organic metals exhibit well known Shubnikov-de Haas
(SdH) oscillations, de Haas-van Alphen (dHvA) oscillations, and magneto-breakdown oscillations
of resistivity. The basic properties of SdH, dHvA, and magneto-breakdown oscillations in
(ET)$_2$X, (ET)$_2$Cu(NCS)$_2$, and (ET)$_2$Mg(SeCN)$_4$ compounds are found to be in a good agree-
ment with standard Lifshits-Kosevich (LK) formula in moderate magnetic fields [2]. At higher
magnetic fields, $H \approx 25 - 30$ T, experimental data are becoming more complicated. Some
measurements on (ET)$_2$Mg(SeCN)$_4$ materials indicate that an effective mass is becoming

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magnetic field dependent at high fields [3], but at the moment it is not clear if this corresponds to some new physics or not.

From our point of view, the major physical question regarding the magnetic oscillations in Q2D materials could be the following: “Which deviations from the LK formula are there appear at high fields when cyclotronic frequency, \( \omega_c \), is of the order of electron bandwidth in the direction perpendicular to plane, \( \omega_c \sim t_\perp [4,5] \), or of the order of Fermi energy?” The latter seems to be possible in \((\text{ET})_2\text{NH}_4\text{Hg(SCN)}_4\) material under pressure where small pockets of quasiparticles are expected to exist [6].

In a tilted magnetic field Q2D compounds demonstrate nontrivial angular resistivity oscillations which are well described in terms of standard theory of anisotropic Q2D metals in the case of \((\text{ET})_2X\) \((X = \text{I}_3, \text{IBr}_2)\) materials [7]. As to angular peaks and dips of resistivity recently discovered in so-called “reconstructed” phase of \((\text{ET})_2\text{MHg(SCN)}_4\) compounds [8], they are believed to be understood in terms of Yamaji’s and Osada’s effects [7,9] within some more complicated variant of “fermiology” (see, for example, [10]).

The physical origin of magnetic and angular oscillations observed in quasi-one-dimensional (Q1D) organic conductors \((\text{TMTSF})_2X\) \((X = \text{ClO}_4, \text{PF}_6, \text{etc.})\) and \((\text{DMET-TSeF})_2\text{AuCl}_2\) is not so simple. In a metallic state these compounds (which have only open Fermi surfaces (FS’s)) demonstrate such unusual phenomena as “magic angles” (MA’s) (i.e., nontrivial angular dependences of both transport [11–18] and thermodynamic [15] properties), “rapid magnetic oscillations” (RMO) of resistivity [19–26], Danner-Chaikin’s (DC) angular resonance [27], and Osada-Kagoshima’s (OK) angular effect [28]. By the present moment there has been done a lot of attempts to understand the origin of MA, RMO, DC, and OK phenomena. All existing theories of magnetic properties in a metallic state of Q1D conductors can be subdivided into three groups:

A) “Fermiological” theories (FL) [9,27–32] (i.e., refined variants of the standard theory of metals).

B) Perturbative many-body (MB) theories [33–42] (these theories consider the changes of many-body effects in a magnetic field, particularly, the dependence of “e-e” interactions on both a magnitude and a direction of a magnetic field).

C) “Luttinger liquid” (LL) approach (the authors of Refs. [43,44] suppose that Tomonaga-Luttinger liquid is created in large enough magnetic fields and claim that magnetic properties cannot be understood within the framework of perturbative MB theories).

Let us discuss in brief the relationship between these different groups of theories. Important methods of an investigation of open FS’s, DC and OK angular resonances, were suggested [27,28] and explained [27,32] within the standard “fermiology”. These angular oscillations are observed in a magnetic field which is almost parallel to the chains of Q1D compounds. For such directions of a magnetic field, as it is shown in references [27,28,32], FL approach works well.

It is important that, unlike DC and OK effects, MA and RMO phenomena are observed in magnetic fields with large enough projection perpendicular to the chains, \( H_\perp \geq 1 - 2 \) T. From papers on spin-density-wave (SDW) formation in \((\text{TMTSF})_2X\) compounds [45–50], it is known that perpendicular projection of a magnetic field, \( H_\perp \), increases “e-e” interactions due to the phenomenon of “one-dimensionalization” of an electron spectrum [45,46,50]. At low temperatures this effect results in the appearance of a cascade of field-induced SDW subphases [19,51]. A variant of perturbative MB approach to magnetic phenomena developed by us and other authors [35–42] is based on the similar idea but it claims more. We argue that perpendicular component of a magnetic field changes an effective dimensionality of “e-e” interactions in Q1D metals and that these changes of the dimensionality are responsible for exotic magnetic properties. From this point of view, a dramatic increase of “e-e” interactions...
in the presence of large enough perpendicular magnetic field \([33-36, 45-47]\) seems to be the main reason why FL theories \([9, 29, 30]\) meet with serious difficulties \([31]\) while describing MA and RMO phenomena.

Below, we present our subjective perturbative MB view of the physical nature of magnetic and angular oscillations in Q1D metals. In 1986, thermodynamic MA phenomena were predicted to exist in field-induced SDW state of \((\text{TMTSF})_2X\) conductors \([50]\). A few years later, the similar effect was suggested for a resistivity in a metallic state \([35]\) and then it was experimentally discovered \([11-18]\). Although at first there were serious discrepancies between original prediction \([35]\) and experimental data \([11-18]\), all of the qualitative discrepancies have been eliminated in a different variant of this model \([36]\). To meet experimental data \([11-18]\), a crossover between 1D “e-e” interactions (at arbitrary directions of a magnetic field) and 2D “e-e” interactions (at MA directions of a magnetic field) was suggested in reference \([36]\) (for alternative point of view, see Ref. \([37]\)). Several years ago perturbative MB theories \([39, 40]\) succeeded in explanation of the thermodynamical MA phenomenon observed in a torque in \((\text{TMTSF})_2\text{ClO}_4\) \([15]\). And very recently, there was appeared perturbative MB theory \([38]\) that describes RMO observed in a metallic state of \((\text{TMTSF})_2\text{ClO}_4\) \([19-26]\). At the end of the review, we discuss a possibility of a break of the Fermi liquid description at high magnetic fields within a framework of the perturbative MB approach \([39-42]\). On the contrary to Fermi liquid theory, we predict \([42]\) the existence of strong forbidden thermodynamic oscillations of magnetic susceptibility and magnetic moment in \((\text{TMTSF})_2\text{ClO}_4\) conductor which has only open FS’s.

To summarize, the major qualitative features of MA and RMO phenomena (at least in \((\text{TMTSF})_2\text{ClO}_4\) and \((\text{DMET-TSeF})_2\text{AuCl}_2\)) seem to be understood within perturbative MB theories \([35-41]\), although new experiments on \((\text{TMTSF})_2\text{PF}_6\) \([44]\) demonstrate that a real physical picture of magnetic phenomena is more complicated.

Among experiments which are still far from the understanding are:
A) Nonmonotonous temperature dependence of a resistivity and unusual temperature dependence of a magnetic susceptibility in moderate magnetic fields \([52]\).
B) Non-analytical dependence of resistivity on a magnetic field in \((\text{TMTSF})_2\text{PF}_6\) \([44]\).
C) Disappearance of a quasiclassical magnetoresistance background and FS-effects in \((\text{TMTSF})_2\text{PF}_6\) when perpendicular magnetic field is applied \([44]\).

From our opinion, the nonmonotonous temperature dependence of resistivity \([52]\) might be prescribed to a modification of “e-e” interactions in the vicinity of SDW transition \([53]\) or to some localization effects in a field \([52, 54]\). As to the other phenomena mentioned above, it seems that they represent some new non-Fermi-liquid physics. Note that non-analytical magnetic field dependence of resistivity \([32]\) observed in experiments on OK angular resonance \([28]\) has been explained by FL approach \([32]\). Nevertheless, the non-analytical magnetic field dependence of resistivity under more general conditions of experiment \([44]\) as well as the disappearance of FS effects at high enough perpendicular magnetic fields \([44]\) still have no explanation within both FL and perturbative MB approaches. The authors of reference \([43]\) speculate that these phenomena are a manifestation of Luttinger liquid formation.

2. Experiments on “Magic Angles” and “Rapid Magnetic Oscillations”

In this section we present a brief review to some experimental aspects of MA and RMO phenomena in a metallic state of Q1D materials. MA phenomena are usually observed in a magnetic field \(H = (0, H \sin \theta, H \cos \theta)\) which is rotated in \((b, c)\)-plane perpendicular to the direction of the chains, \(a\). If magnetic field is applied along one of the possible vectors of a
"Magic" directions of a magnetic field are defined by vectors of the crystalline lattice \( m (b, c) \) plane. One of the main MA's is shown by arrow (see Eq. (1)).

\[
\tan(\theta_{k,m}) = \left( \frac{k}{m} \right) \left( \frac{b^*}{c^*} \right)
\]

(see Fig. 1), there appear nontrivial peculiarities on angular dependences of resistivity and magnetic torque [11-18]. (Here, \( \theta \) is the angle between \( \mathbf{H} \) and \( c \) axis in an orthorombic model of a unit cell; \( b^* \) and \( c^* \) are the lattice constants; \( k \) and \( m \) are integers).

A typical experimental plot of the second derivative of a magnetoresistance with respect to angle \( \theta \) is presented in Figure 2 where a number of maxima is observed. As it follows from Figure 2, the most common MA's are the integer ones which correspond to \( m = 1 \) in equation (1). Note that two major MA's are observed at \( \theta_{\infty,1} = \pi/2 \) and \( \theta_{0,1} = 0 \) when magnetic field is applied parallel to \( b \) and \( c \) axes, correspondingly. The third MA with \( \theta_{11} = \arctan(b^*/c^*) \) is usually also pronounced whereas the higher order MA's correspond to local minima (which have much smaller magnitudes) and so they are visible only on a plot of the second derivative of magnetoresistance (see Fig. 2).
We point out that MA phenomena are very common in a low-dimensional conductors. MA’s are observed in experiments on transport properties [11–18] as well as in experiments on thermodynamic ones [15,16] in a number of Q1D compounds in both a metallic and a SDW phases. The other unusual magnetic phenomenon, RMO, is very common in a SDW state [20–26,55–58]. As to metallic state, RMO have been revealed only in (TMTSF)2ClO4 compound by the present moment [19–26]. A typical experimental geometry of an observation of the resistive RMO corresponds to $H \parallel c$ and $I \parallel a$.

3. “Fermiological” Approach to “Magic Angles” and “Rapid Magnetic Oscillations” Difficulties

Electron spectrum of (TMTSF)$_2$X (X = ClO$_4$, PF$_6$, AsF$_6$, etc.) and (DMET-TSeF)$_2$AuCl$_2$ compounds corresponds to open slightly deformed sheets of the FS:

$$\epsilon(p) = \pm v_F(p_a \mp p_F) + 2t_b \cos(p_b b^*) + 2t_c \cos(p_c c^*) + \sum_{k>0, m=1} 2t_{k,m} \cos(kp_b b^* + mp_c c^*)$$

(2)

(see Fig. 3).

In equation (2) the first term represents a free motion of electrons along the chains on the right (+) and left (−) sheets of the FS, with $p_F$ and $v_F$ being the Fermi momentum and the Fermi velocity, correspondingly ($p_F v_F \approx 2.5 \times 10^8$ K). The second and the third terms correspond to the hopping of electrons in the perpendicular directions, b and c ($t_b \approx 200$ K, $t_c \approx 5$ K). The summation over $k$ and $m$ in equation (2) represents a small effective hopping of electrons in (b, c) plane. Note that, due to the extremely small overlapping of the wave functions along c axis, we retain only the terms with $m = 1$. It is known that in (TMTSF)$_2$PF$_6$ and (DMET-TSeF)AuCl$_2$ compounds $t_{1,1} \gg t_{2,1} \gg t_{3,1}$ and so we may retain only a few terms (say, with $k = 1$ and $k = 2$) in the summation in equation (2). On the contrary, in the case of (TMTSF)$_2$ClO$_4$ we have to consider all higher harmonics, $t_{k,1} \sim 1$ K ($k = 1, 2, 3, \ldots$) due to a corrugation of FS resulted from an anion ordering (AO) gap. $\square$ [1] (see Fig. 4).

In this section we show that in Q1D conductors with electron spectrum (2) negligible magnetoresistance effects are expected within the standard “fermiology” if inchain current, $I \parallel a$, 

\[
p_b \quad \pi \quad b \quad \pi \quad a \quad \pi \quad c \quad \pi \quad
\]

Fig. 3. — A typical Q1D electron spectrum.
Due to the AO phenomena, an electron spectrum in (TMTSF)$_2$ClO$_4$ corresponds to four open sheets of FS.

is measured in a magnetic field perpendicular to the chains, $\mathbf{H} \perp \mathbf{a}$. Indeed, let us consider an electron motion along open sheets of the FS in perpendicular magnetic field (see Figs. 3, 4). Due to small values of the parameters $t_b, t_c$, and $\Box$, longitudinal component of an electron velocity, $v_a$, is almost independent on the position of an electron on the FS, $v_a(p) \simeq \pm v_F$. This means that magnetic field does not disturb an electron motion along the chains and thus no any magnetic effects are expected.

The same is directly seen from Boltzmann kinetic equation:

$$e\mathbf{E} \frac{d[f(p)]}{dp} + \left( \frac{e}{c} \right) [\mathbf{v} \times \mathbf{H}] \frac{d[f(p)]}{dp} = -\left( \frac{1}{\tau} \right) [f(p) - f_0(p)], \quad (3)$$

where the Lorentz force, $\mathbf{F} = (e/c)[\mathbf{v} \times \mathbf{H}] \simeq (e/c)v_FH(b/b^*)$, is almost perpendicular to the electric field, $\mathbf{F} \perp \mathbf{E}(\mathbf{a})$; $e$ is the electron charge and $c$ is the velocity of light.

A simple analysis of equation (3) shows that all magnetic effects in longitudinal resistivity have to be of the order of $\delta \rho_a(\mathbf{H})/\rho_a(0) \sim (t_b/p_Fv_F)^2 \sim 10^{-2} - 10^{-3}$. This conclusion of the standard theory of metals is in a sharp contradiction with experimental data, since the experimental relative magnitudes of MA [11,13,14,16-18] and RMO [19-26] phenomena are of the order of 1 and $10^{-1}$, correspondingly.

To avoid the above mentioned contradiction, an interesting phenomenological "hot spots" model was recently proposed [30]. This model suggests the existence of some spots on the Q1D FS where electron relaxation time is much less than one on the rest part of FS. Although a subsequent "microscopic" theory [31] confirms the existence of "hot lines" on the FS, nevertheless, it is not able to explain the appearance of MA's in a magnetic field. There exist also two other FL theories of MA's [9,29], but both of them are inconsistent with large experimental magnitudes of MA effects and their main experimental features.

To summarize, we conclude that an observation of MA's and RMO in a Q1D conductors is in a dramatic contradiction with the predictions of "fermiology" based on the standard theory of metals.

Before proceeding to the formulation of a perturbative MB approach to the magnetic properties in Q1D metals, let us consider a phenomenon of “one-dimensionalization” of a Q1D electron spectrum in a tilted magnetic field [50]. In this section we ignore a small last term in equation (2) and discuss a quasiclassical equations of electron motion in the case of a simplified electron spectrum:

\[ \epsilon(p) = \pm v_F (p_a \mp p_F) + 2t_b \cos(p_b b^*) + 2t_c \cos(p_c c^*) \] (4)

in a magnetic field \( \mathbf{H} = (0, H \sin \theta, H \cos \theta) \).

Due to small corrugations of the FS (4), there are basically two projections of the Lorentz force which lie in \((b, c)\) plane:

\[ \frac{dp_b}{dt} = e \frac{v_F}{c} H \cos \theta, \quad \frac{dp_c}{dt} = -e \frac{v_F}{c} H \sin \theta \] (5)

Using well known quasiclassical equations [59]:

\[ v_b = \frac{d\epsilon(p)}{dp_b}, \quad v_c = \frac{d\epsilon(p)}{dp_c}, \]

it is easy to obtain electron trajectories in a magnetic field:

\[ x_b^H(t) = \frac{2t_bb^*}{\omega_b} \cos[\omega_b t], \quad x_c^H(t) = -\frac{2t_cc^*}{\omega_c} \cos[\omega_c t], \] (7)

where \( x_b^H \) and \( x_c^H \) are electron coordinates in \((b, c)\) plane, \( t \) is a time: \( \omega_b = eH v_F b^* \cos \theta / c, \omega_c = eH v_F c^* \sin \theta / c. \)

From equations (7), it is followed that in the plane perpendicular to the chains electron motion is a quasiperiodic and localized within a tube which cross section area, \( S_\perp(H) = (2t_b b^*/\omega_b) \times (2t_c c^*/\omega_c) \), is decreasing with increasing of a magnetic field. Using a quantum mechanical language we can say that electron wave functions are localized in \((b, c)\) plane and are extended only along a axis.

Indeed, let us consider a quantum mechanical problem. In a tilted magnetic field in the Landau gauge, \( \mathbf{A} = (0, H \cos \theta x, -H \sin \theta x) \), the Schrodinger equations are given by Peierls substitution \( p \rightarrow p - (e/c)A \) [59]:

\[ \left[ \mp i \frac{v_F}{d} - 2t_b \cos(p_bb^* - \frac{\omega_b x}{v_F}) + 2t_c \cos(p_cc^* + \frac{\omega_c x}{v_F}) \right] \psi_{1D}^{\pm}(x, p_b, p_c) = \epsilon \psi_{1D}^{\pm}(x, p_b, p_c) \] (8)

From equation (8), it is easy to derive electron wave functions in \((x, p_b, p_c)\) representation:

\[ \psi_{1D}^{\pm}(x, p_b, p_c) = \exp(\pm i \frac{\epsilon x}{v_F}) \exp[\pm i \frac{\lambda_b}{2} \sin(p_bb^* - \frac{\omega_b x}{v_F})] \exp[\mp i \frac{\lambda_c}{2} \sin(p_cc^* + \frac{\omega_c x}{v_F})] \] (9)

and to convert them into Wannier representation:

\[ \psi_{L,N}^{\pm}(x, y = lb^*, z = nc^*) = \exp \left\{ \pm i \left( \frac{\epsilon}{v_F} (l - L)\omega_b \mp (n - N)\omega_c \right) x \right\} \]

\[ \times J_{\pm(L-l)}(\lambda_b/2)J_{\mp(N-n)}(\lambda_c/2), \]

where \( \lambda_b = 4t_b/\omega_b, \lambda_c = 4t_c/\omega_c. \)
Using properties of Bessel functions, \( J_n(x) [60] \), it is possible to make sure that wave functions (10) exponentially decay in \((b, c)\) plane on the scales \( x_b^H \sim \lambda_b b^* \) and \( x_c^H \sim \lambda_c c^* \). On the contrary, wave functions are extended along the chains and thus energy depends on momentum \( p_a \) and two quantum numbers:

\[
\epsilon(p) = \pm v_F (p_a \mp p_F) \mp l'\omega_b \pm n'\omega_c \quad (11)
\]

To summarize, in this section we demonstrated that at arbitrary direction of a magnetic field electron motion is becoming effectively 1D if we take a simplified electron spectrum (4).


Below, we consider a more realistic electron spectrum (2) to describe an important for further development effect of “two-dimensionalization” of an electron spectrum at MA directions of a magnetic field [36].

Since electron hoppings, \( t_{k,1} \), are small comparable with \( t_b \) and \( t_c \), in the vicinity of each MA only one term is important in the summation in equation (2). Suppose we retain the term \( t_{k,1} \) in equation (2). Let us see what is happened if a magnetic field is applied strictly along the MA \( \theta_{k,1} \) (see Eq. (1)). Due to a conservation of momentum projection along a magnetic field, electrons can move free along the direction \((k, 1)\) if \( t_{k,1} \neq 0 \). This motion obviously destroys the “one-dimensionalization” of electron spectrum (9)-(11) and results in the appearance of 2D properties since a free electron motion is now allowed along the chains and along the direction of a magnetic field.

Let us demonstrate the above mentioned crossover between 1D and 2D electron motion by using a quantum mechanical language. In the case of electron spectrum (2) the Schrodinger equations can be rewritten as follows:

\[
\Psi_{1D}^\pm (x, p_b, p_c) = \exp [\mp i \lambda_{k,1} \sin (\frac{\omega_{k,1} x}{2v_F}) \frac{\cos (kpb^* + pcc^*) - \omega_{k,1}}{2v_F}] \Psi_{1D}^\pm (x, p_b, p_c),
\]

where \( \Psi_{1D}^\pm (x, p_b, p_c) \) are defined by equation (9); \( \omega_{k,1} = k \omega_b - \omega_c = e(v_F/c)Hc^* \cos \theta (\tan \theta_{k,1} - \tan \theta) \), \( \lambda_{k,1} = 4t_{k,1}/\omega_{k,1} \).

A change of an effective dimensionality at MA’s is directly followed from equations (13). Indeed, if the direction of a magnetic field tends to the angle \( \theta = \theta_{k,1} \) then \( \omega_{k,1} \to 0 \) and \( \lambda_{k,1} \to 2t_{k,1}x/v_F \). As a result, at \( \theta = \theta_{k,1} \) we have essentially 2D electron wave functions:

\[
\Psi_{2D}^\pm (x, p_b, p_c) = \exp [\mp i 2t_{k,1}x \cos (kpb^* + pcc^*)/v_F] \Psi_{2D}^\pm (x, p_b, p_c)
\]

and 2D electron spectrum:

\[
\epsilon(p) = \pm v_F (p_a \mp p_F) + 2t_{k,1} \cos (kpb^* + pcc^*) \mp (l' - n'k)\omega_b
\]

Unlike 1D spectrum (11), electron energy at “magic” directions of a magnetic field (15) depends on two momenta: momentum along the chains, \( p_a \), and momentum along a magnetic field, \( p_H = kpb + p_c \).
To summarize this part, we can make a remarkable conclusion. At arbitrary irrational directions of a magnetic field electron spectrum is effectively 1D (see Eq. (11)), whereas at rational directions (1) (i.e., at MA's) electron spectrum is becoming 2D (15). Since \( t_{k,1} \) is quickly decreasing with the number \( k \), in practice only a few MA's are important. Therefore, if a magnetic field is far enough from the main MA's we can regard electrons as effectively 1D.


Let us call MA's which correspond to negligible values of \( t_{k,m} \) in equation (2) the secondary MA's to distinguish them from the main MA's. As it is shown in the next section, the secondary MA's can be defined by a simple inequality \( t_{k,m} \ll T \). From equation (15), it follows that for negligible values of \( t_{k,m} \) electron spectrum is again becoming effectively 1D. It is important that at “magic” directions of a magnetic field \( k\omega_b = m\omega_c \) in equation (11) and thus at the secondary “magic” directions of a field 1D electron spectrum (11, 15) depends only on one quantum number \( (l'-n'k) \). This degeneration means that cyclotronic frequencies of an electron motion along \( b \) and \( c \) axes in a real space are commensurate and that electron trajectories in \((b,c)\) plane are periodic (see Eq. (7)). This crossover from a quasiperiodic motion (at arbitrary directions of a field) to a periodic one (at the secondary MA directions of a field [35, 50]) is particularly important in a SDW phase where the secondary MA's, \( \theta_{1,2} \) and \( \theta_{1,3} \), are sometimes more pronounced than the main MA's [12, 16]. As to metallic phase, the secondary MA are visible only as a local maxima of resistivity on a plot of the second derivative of magnetoresistance (see Fig. 2).

7. Model

In the previous sections we described unique phenomena of a change of the effective dimensionality of electrons in a tilted magnetic field. It is known that “e-e” interactions are dramatically increasing with the decreasing of the dimensionality. In 1D case “e-e” interactions are becoming singular that leads to Peierls instability at low temperatures and high magnetic fields [19, 45-51].

In this review we discuss properties of a metallic phase well above SDW phase transition. Our message is that the changes of “e-e” interactions with changing of an effective dimensionality define some magnetic properties in a metallic phase. We show that such nontrivial experimental properties as MA's and RMO can be prescribed to the dependence of many-body effects on effective dimensionality in a magnetic field.

From experiments on SDW formation in a magnetic field [19, 51, 61, 62], it is known that basic properties of the field-induced SDW state are understood within the framework of a weak coupling model [45-50]. As it is shown in the previous sections, the effective dimensionality of electrons is decreasing in a magnetic field, therefore an applicability of the perturbative MB approach is becoming not clear. Nevertheless, even in strong magnetic fields, \( H \approx 20 - 50 \) T, there exists an important large parameter which is the area of the localization of electron wave functions in \((b,c)\) plane expressed in terms of the area of unit cell: \( S_1(H) = S_\perp(H)/(b^*c^*) \approx 10 - 20 \). It is possible to make sure that the existence of a small parameter \( 1/S_\perp(H) \leq 10^{-1} \) allows us to consider “e-e” interactions within the perturbative approach.

8. Many-Body Theory of “Magic Angles” [36]

It is known that in the absence of a magnetic field longitudinal resistivity obeys a \( T^2 \) law in \((TMTSF)_2X\) materials in a broad temperature region. This indicates that “e-e” scattering is
likely a dominant mechanism of resistivity in these Q1D materials. Moreover, (TMTSF)$_2$X conductors can be considered as 1:2 commensurate ones that provides the existence of an effective mechanism of resistivity — the so-called Umklapp “$e$-$e$” scattering with nonconservation of the total momentum: $p_1^2 + p_2^2 = p_3^2 + p_4^2 + 4p_F$ [1]. Below, we suggest that Umklapp scattering is a major mechanism of resistivity and consider the influence of changes of an effective dimensionality in a magnetic field on the probability of Umklapp scattering.

Let us take a realistic electron spectrum (2) and calculate a longitudinal resistivity in a magnetic field by means of variational principle for “$e$-$e$” scattering [63]. According to [63], to estimate a longitudinal resistivity it is necessary to average a probability of Umklapp scattering, $\langle U(p_1, p_2; p_3, p_4) \rangle$, over electron momenta using Fermi factors, $n(e)$:

$$\rho_a \sim 1/\tau = \int_{-\infty}^{+\infty} dp_1^2 dp_2^2 dp_3^2 dp_4^2 \ U(p_1, p_2; p_3, p_4) \delta(p_1 + p_2 - p_3 - p_4)$$

$$\times \delta(e(p_1) + e(p_2) - e(p_3) - e(p_4)) n(e(p_1)) n(e(p_2)) [1 - n(e(p_3))] [1 - n(e(p_4))], \quad (16)$$

where $\rho_a \sim 1/\tau$ in a Q1D case; $1/\tau$ is an inverse “$e$-$e$” relaxation time.

Substituting electron wave functions (13) into equation (16) we get the following expression for resistivity [36]:

$$\rho_a(H, \theta) \sim g_U^2 \ T \int_{-\infty}^{+\infty} \ dx \ \frac{2\pi T/v_F}{\sinh(2\pi T x/v_F)} \left[ \frac{2\pi T|x|/v_F}{\exp(4\pi T|x|/v_F) - 1} + \frac{2\pi T|x|/v_F - 1}{2} \right] \langle ABC \rangle,$$

$$A = \exp[-2i\lambda_b \sin(\frac{\omega_b x}{2}) \cos(\alpha \sin \alpha_1 + \sin \alpha_2)], \quad (17)$$

$$B = \exp[-2i\lambda_c \sin(\frac{\omega_c x}{2}) \cos(\beta \sin(\beta_1 - k(1) - \frac{\pi k_2}{2}) + \sin(\beta_2 - k(2) - \frac{\pi k_2}{2}))], \quad (19)$$

$$C = \exp[-2i\lambda_{k,1} \sin(\frac{\omega_{k,1} x}{2}) \cos(\alpha k + \beta)(\sin \beta_1 + \sin \beta_2)], \quad (20)$$

where $g_U$ is a dimensionless constant of Umklapp “$e$-$e$” scattering, $\langle . . \rangle$ means the averaging over all variables $\alpha, \beta, \alpha_1$, and $\beta_1$ ($i = 1, 2$). (Note that Zeeman splitting is not important while calculating a probability of Umklapp scattering in a magnetic field [35,36], therefore, we omit electron spins everywhere).

In an experimentally interesting case, when $\omega_b, \omega_c \gg T$, it is possible to derive some analytical expressions both for resistivity at the main MA, $\theta = \theta_{k,1}$ ($t_{k,1} \gg T$):

$$\rho_a(H, \theta_{k,1}) \sim \frac{1}{\tau}(H, \theta_{k,1}) = \frac{g_U^2 T^2}{\lambda_b \lambda_c t_{k,1}} \ln^2(\lambda_b), \quad \omega_c \ll t_c,$$

$$\rho_a(H, \theta_{k,1}) \sim \frac{1}{\tau}(H, \theta_{k,1}) = \frac{g_U^2 T^2}{\lambda_b t_{k,1}} \ln^2(\lambda_b), \quad \omega_c \gg t_c. \quad (21)$$

and for magnetoresistance background (when the direction of a magnetic field is far enough from the main MA, $\delta \theta = |\theta - \theta_{k,1}| \gg (T/\omega_c) \ll 1$):

$$\rho_a(H, \theta) \sim \frac{1}{\tau}(H, \theta) = \frac{g_U^2 T}{\lambda_b \lambda_c} \ln^2(\lambda_b) \sin(2\theta), \quad \omega_c \ll t_c;$$

$$\rho_a(H, \theta) \sim \frac{1}{\tau}(H, \theta) = \frac{g_U^2 T}{\lambda_b} \ln^2(\lambda_b) \cos \theta, \quad \omega_c \gg t_c. \quad (22)$$
Fig. 5. — A few main MA’s are shown in the plot of the calculated resistive MA phenomenon (see Eq. (17)).

Here, we also present the expressions for resistivity in the vicinity of MA, \( \delta \theta \ll (T/\omega_c) \ll 1 \), in the case when \( \theta_{k,1} \) is the main MA:

\[
\frac{\rho_a(H, \theta) - \rho_a(H, \theta_{k,1})}{\rho_a(H, \theta_{k,1})} = A_1(\delta \theta)^2 \frac{\omega_c^2}{(2\pi T)^2}, \quad t_{k,1} \gg T;
\]

and in the case when \( \theta_{k,1} \) is the secondary MA:

\[
\frac{\rho_a(H, \theta) - \rho_a(H, \theta_{k,1})}{\rho_a(H, \theta_{k,1})} = -A_2(\delta \theta)^2 \frac{\omega_c^2}{(2\pi T)^2}, \quad t_{k,1} \ll T,
\]

where \( A_1, A_2 > 0 \) are dimensionless parameters; \( A_1, A_2 \sim 1 \).

From equations (21) and (22), it follows that in the case \( t_{k,1} \geq T \) an effective “two-dimensionalization” of an electron spectrum (see Eqs. (14, 15)) leads to the appearance of 2D \( (T^2) \) temperature dependence of resistivity at the main “magic” directions of a magnetic field whereas at the directions of a field which are far enough from the main MA’s resistivity possess 1D \( (T^1) \) temperature dependence. This phenomenon results in the appearance of angular minima of resistivity at the main MA’s (see Eq. (23) and Fig. 5). In the opposite case, \( T \geq t_{k,1} \), a secondary MA phenomenon (see Sect. 6) leads to the appearance of maxima of resistivity (24). The prediction of the main and the secondary MA’s [36] is in a qualitative agreement with experimental data on \((\text{TMTSF})_2\text{ClO}_4\) [12,15] (see Fig. 2 in Ref. [36]).

Below, we also present the expressions for resistivity at two major “magic” directions of a magnetic field, \( \theta_{\infty,1} = \pi/2 \) and \( \theta_{0,1} = 0 \), that are slightly different from the common equation (21):

\[
\rho(H, \theta_{\infty,1}) \sim \frac{1}{\tau_c}(H, \theta_{\infty,1}) = \frac{T^2}{t_b} \ln^2(t_b/T)
\]

(25)
\[
\rho(H, \theta_{0,1}) \sim \frac{1}{\tau}(H, \theta_{0,1}) = g^2 \frac{T^2}{\lambda_b t_c} \ln^2(\lambda_b) \ln(t_c/T) (26)
\]

Note that at $\theta = \pi/2$ magnetoresistance is negligible in accordance with the experimental data [11–18], whereas at other main MA’s (including $\theta_{0,1}$) magnetoresistance obeys an approximately $H^1$ law.

9. Many-Body Theory of “Rapid Magnetic Oscillations” [38]

In this section we extend perturbative MB approach to describe another exotic phenomenon — RMO of resistivity experimentally observed in (TMTSF)$_2$ClO$_4$. As it is mentioned in Section 3, RMO of longitudinal resistivity, $\rho_L$, are completely unexpected within a framework of standard theory of metals. Below, we prescribe the existence of RMO in (TMTSF)$_2$ClO$_4$ to unusual many-body effect — the oscillations of a probability of “e-e” Umklapp scattering in a magnetic field.

For further development it is important that, unlike (TMTSF)$_2$PF$_6$, (TMTSF)$_2$ClO$_4$ has AO gap, $\Box$, in its electron spectrum (see Fig. 4). By present moment, it has been established that when AO gap is suppressed strong RMO of resistivity disappear [24,64]. This means that the appearance of RMO in (TMTSF)$_2$ClO$_4$ is closely related to the existence of AO gap. Thus to describe RMO we need to find a solution of quantum mechanical problem in a magnetic field for some more complicated electron spectrum.

The AO introduces a periodic potential, $\Box(y) = \Box \cos(\pi y / b^*)$, which leads to a doubling of the crystalline lattice in $y(b)$ direction. Therefore, in the presence of the AO gap there are four components of an electron wave function which obey the following Schrodinger equations:

\[
[\pm v_F(p_x \mp p_F) + 2t_b \cos(p_y b^*)] \psi_{\pm}^\chi(p_b) + \Box \psi_{\pm}^\chi(p_b + \pi/b^*) = \epsilon \psi_{\pm}^\chi(p_y), (27)
\]

\[
[\pm v_F(p_x \mp p_F) - 2t_b \cos(p_y b^*)] \psi_{\pm}^\chi(p_b + \pi/b^*) + \Box \psi_{\pm}^\chi(p_b) = \epsilon \psi_{\pm}^\chi(p_y + \pi/b^*) (28)
\]

As a result, electron spectrum splits into four open Q1D branches:

\[
\epsilon_n(p) = \pm v_F(p_x \mp p_F) + (-1)^n \sqrt{[2t_b \cos(p_y b^*)]^2 + \Box^2}, (29)
\]

where $n = 1, 2, 3, 4$ (see Fig. 4).

(All small parameters, $t_c \approx 5$ K and $t_{k,m} \sim 1$ K, are omitted in equations (27-29) since we study RMO in a high field metallic state ($H \geq 15 – 30$ T) which is stable at temperatures $\pi T \geq 17$ K $\geq 2t_c, 2t_{k,m}$).

In a magnetic field $\mathbf{H} \parallel \mathbf{c}$ the Schrodinger equations in the Landau gauge, $\mathbf{A} = (0, H x, 0)$, take a form [38]:

\[
[\mp i v_F \frac{d}{dx} + 2t_b \cos(p_y b^*) - \omega y x v_F] \psi_{\pm}^\chi(x, p_b) + \Box \psi_{\pm}^\chi(x, p_b + \pi/b^*) = \epsilon \psi_{\pm}^\chi(x, p_b), (30)
\]

\[
[\mp i v_F \frac{d}{dx} - 2t_b \cos(p_y b^*) - \omega y x v_F] \psi_{\pm}^\chi(x, p_b + \pi/b^*) + \Box \psi_{\pm}^\chi(x, p_b) = \epsilon \psi_{\pm}^\chi(x, p_b + \pi/b^*) (31)
\]

where the amplitudes $\psi_{\pm}^\chi(x, p_b)$ and $\psi_{\pm}^\chi(x, p_b + \pi/b^*)$ are related to electron Bloch functions, $\Psi_{\pm, p_b}^\chi(x, y)$:

\[
\Psi_{\pm, p_b}^\chi(x, y) = \exp(i p_b y) [\psi_{\pm}^\chi(x, p_b) + \exp(i \pi y/b^*) \psi_{\pm}^\chi(x, p_b + \pi/b^*)] (32)
\]

(Note that equations (30-32) demonstrate an effective “one-dimensionalization” of electrons which have Q1D electron spectrum (29), since electron energy, $\epsilon$, does not depend on momenta perpendicular to the chains, $p_y$ and $p_z$).
From experimental works on RMO [21–26], it is known that magnetic breakdown phenomenon occurs between two open sheets of FS denoted by \( n = 1 \) and \( n = 2 \) in equation (29) (see also Fig. 4). Experimental estimation of magnetic breakdown field [23] gives the following value: \( H_{MB} \approx 15 \) T. If we take a theoretical expression for \( H_{MB} \) [65] calculated for electron spectrum (29):

\[
H_{MB} = \pi \Box^2 / 2e v_F b^* ,
\]
we can conclude that \( \Box \simeq 50 – 70 \) \( K < 2t_b \simeq 400 \) \( K \).

From general theory [59], it is known that perturbative approach to Schrodinger equations (30, 31) is valid if \( H > H_{MB} \). In this case the zeroth-order wave functions are a symmetrical (S) and an antisymmetrical (A) combinations of two solutions of equations (30, 31) in the absence of the AO gap [38,66,67]:

\[
\begin{align*}
\psi^+_S(x, p_b) & = \exp \left[ \pm i \left( \epsilon - \Box^* \right) x \right] \left( + \exp \left[ \pm \frac{i \lambda}{2} \sin \left( \frac{p_b b^*}{v F} \right) \right] \right) \sin \left( \frac{\omega b x}{v F} \right) , \\
\psi^+_S(x, p_b + \pi \over b^*) & = \exp \left[ \pm i \left( \epsilon + \Box^* \right) x \right] \left( + \exp \left[ \pm \frac{i \lambda}{2} \sin \left( \frac{p_b b^*}{v F} \right) \right] \right) \sin \left( \frac{\omega b x}{v F} \right) .
\end{align*}
\]

The S and A wave functions are characterized by the following 1D branches of an electron spectrum [38,66,67]:

\[
\begin{align*}
\epsilon^+_S(p) & = \pm v_F \left( p_a \mp p_F \right) + \Box^* , \\
\epsilon^+_A(p) & = \pm v_F \left( p_a \mp p_F \right) - \Box^* ,
\end{align*}
\]

where

\[
\Box^* \simeq \Box \sqrt{\omega_b / 2 \pi t_b \cos (4 t_b c / e v_F H b^*)}
\]
is a difference between the energies of S and A wave functions. The absolute value of this energy splitting, \( |\Box^*(H)| \), rapidly oscillates with inverse magnetic field:

\[
\delta \left( \frac{1}{H} \right) = \frac{\pi e v_F b^*}{4 t_b c}
\]

At this point we are ready to calculate inverse “e-e” relaxation time by substituting known electron wave functions (34, 35) and electron spectrum (36, 37) into equation (16). As a result, we have the following formula for a longitudinal resistivity:

\[
\rho_a(H) \sim \frac{1}{\tau(H)} = \frac{1}{\tau} \int_{-\infty}^{+\infty} dx \frac{2 \pi T / v_F}{\sinh^2 (2 \pi T x / v_F)} \left[ \frac{2 \pi T |x| / v_F}{\exp (4 \pi T |x| / v_F) - 1} + \frac{2 \pi T |x| / v_F - 1}{2} \right] \int_0^{2 \pi} dk_1 J_0^2 (2 \lambda_b \sin \left( \frac{\omega b x}{2 v_F} \right) \sin (k_1)) \cos^4 \left( \frac{\Box^* x}{v_F} \right),
\]

where we have used inequality \( \lambda_b \gg 1 \).

Since equation (40) contains the oscillatory function, \( \Box^*(H) \), longitudinal resistivity have to possess a strong magnetic oscillations. We note that frequency of the resistive oscillations is given by equation (39) and corresponds to area between two open orbits (see Fig. 4). The calculated magnitude of these oscillations possess unusual magnetic field and temperature
dependences. It is important that at high enough temperatures resistive oscillations (39, 40) should demonstrate a power law temperature decay [38] due to their many-body nature:

\[ \frac{\delta \rho_a(H)}{\rho_a(H)} = B_1 \frac{\Box^2 \omega_c}{\pi^3 T^2 t_b} \cos(8t_b c/\pi e v_F b^*) + B_2 \frac{\Box^4 \omega_c^2}{\pi^6 T^4 t_b^2} \cos(16t_b c/\pi e v_F b^*) \]  

where \( B_1, B_2 \sim 1 \).

A power law temperature decay (41) clearly distinguishes the many-body oscillations (39-41) from all known kinds of one-body oscillations (i.e., SdH oscillations, so-called, Stark oscillations, and magneto-breakdown ones). Indeed, as it is known [59], SdH and magneto-breakdown oscillations decay exponentially with temperature whereas Stark oscillations do not depend on temperature per se. We would like to point out that Stark oscillations and many-body oscillations suggested in this section can coexist in metals with open FS’s such as (TMTSF)\(_2\)X conductors. Therefore, we stress that, besides different temperature dependences, there exists another important difference between these two kinds of oscillations. It is possible to make sure that equations (40) and (41) are qualitatively applicable at arbitrary direction of an electron current. I. This means that many-body oscillations (39-41) have to exist at arbitrary directions of a current, while Stark oscillations have to be observed only when \( \mathbf{I} \parallel \mathbf{b} \) (see, for example, [59]). The above mentioned features of many-body RMO are in accordance with preliminary experimental data [68].

In conclusion of this section we note that equation (40) exhibits a weak oscillations, \( \frac{\delta \rho_a(H)}{\rho_a(H)} \leq 10^{-2} \), even in the absence of the AO gap (i.e., when \( \Box^* = 0 \)). This type of oscillations was early proposed in reference [33]. Recent experiments [24] show that strong RMO (39-41) may coexist with weak oscillations [33] in a metallic state of (TMTSF)\(_2\)ClO\(_4\).

10. Forbidden Magnetic Oscillations: a Break of Fermi Liquid Description [40, 42]

In this section we show that perturbative MB approach to magnetic properties of a 1D metals predicts some phenomenon which is in disagreement with Fermi liquid theory. It is well known that Fermi liquid picture does not allow any thermodynamic oscillations to exist in metals which have only open quasiparticle orbits [59]. Below, we demonstrate that a magnetic field dependence of “e-e” interactions in Peierls channel dramatically changes this statement at high magnetic fields, \( H \geq 20 - 30 \) T. We show that, due to the oscillatory nature of magneto-breakdown electron spectrum (see Eqs. (36, 37) and Fig. 6), a magnetic field dependence of Peierls interaction results in the appearance of a strong forbidden oscillations of magnetic susceptibility and magnetic moment [42].

A few years ago it was shown [39, 40] that magnetic field dependence of the contribution to free energy due to “e-e” interactions in a metallic state can explain thermodynamic MA phenomenon [15]. In addition, in reference [40] a new mechanism of thermodynamical magnetic oscillations (which is not related to the presence of any closed orbits) was suggested. Unfortunately, the calculated magnitudes of the oscillations were occured too small to be observed, \( \delta M/M_0 \sim 10^{-3} - 10^{-4} \) [40].

In this section we discuss another mechanism of thermodynamic magnetic oscillations in 1D metals [42] which results in the appearance of strong forbidden oscillations. According to reference [40], we consider a second order correction to free energy due to “e-e” interactions in a magnetic field \( \mathbf{H} \parallel \mathbf{c} \) (see Fig. 7). Unlike reference [40], we study oscillations in (TMTSF)\(_2\)ClO\(_4\) compound which has a remarkable oscillatory magneto-breakdown spectrum (36-38).

Below, we describe a method of the calculations of reference [42] and then proceed to the discussion of the final results. Once wave functions and energy spectrum are known (see
Fig. 6. — Electron spectrum of (TMTSF)$_2$ClO$_4$ in a magnetic field consists of four 1D branches of FS (see Eqs. (36-38)) with energy splitting, $\Box^*$, being an oscillatory functions of $1/H$.

Fig. 7. — A second order correction to free energy due to “e-e” interactions; arrows stand for electron spins.

Eqs. (34-38)), one can derive Green functions by using a standard procedure [69]:

$$G^\pm(x, x': p_b, \epsilon_b) = \sum_{\epsilon_1 \in \{\pm A\}} \frac{[\psi_1^\pm(x, p_b)]^* \psi_1^\pm(x', p_b)}{\omega_n - \epsilon_1(p)},$$

(42)
\[ G^\pm(x, x'; p_b, p_b + \frac{\pi}{b'}) = \sum_{i=1}^{4} \left[ \psi_i^\pm(x, p_b) \psi_i^\pm(x', p_b + \frac{\pi}{b'}) \right] \frac{i \omega_n - \epsilon_i^\pm(p)}{\epsilon_i^\pm(p)} , \tag{43} \]

where \( \omega_n = 2\pi T(n + 1/2) \) is a fermion Matsubara frequency.

To estimate a second order correction to free energy it is necessary to make a summation of the Feynman diagrams corresponding to all possible combinations of \( G^\pm(x, x'; p_b, p_b) \) and \( G^\pm(x, x'; p_b, p_b + \pi/b') \) functions which conserve a total electron momentum (see Fig. 7). As a result of this procedure, we get the following correction to the free energy, \( \Delta F(H) \), calculated per one electron [42]:

\[ \Delta F(H) = -\frac{\pi^3 g_{\text{eff}}^2 T^3}{p_F v_F} \int_0^\infty \frac{dx}{2\pi} \int_0^\infty \frac{dF}{2\pi v_F} \sinh^2 \left( \frac{\omega_b x}{2v_F} \right) \sin \left( \frac{\omega_b}{v_F} \right) \cos^2 \left( \frac{\omega_b}{v_F} \right) , \tag{44} \]

where we have used inequality \( \lambda_b \gg 1 \) and have retained only term which does not depend on Zeeman splitting (see Fig. 7); \( g_{\text{eff}} \) is an effective constant of “e-e” interactions.

Although integral (44) is divergent at small \( x \), all measurable physical quantities such as magnetic moment, \( \Delta M = -d(\Delta F)/dH \), and magnetic susceptibility, \( \Delta \chi = d(\Delta M)/dH \), possess only a weak logarithmic divergence. This allows us to estimate these quantities with logarithmic accuracy in the experimentally interesting case when \( \omega_b \gg T \).

Below, we present the final expressions for relative magnitudes of the forbidden oscillations:

\[ \frac{\delta M}{M_0} \approx \frac{g_{\text{eff}}^2}{4\pi^5} \left( \frac{\omega_b}{\mu_B H} \right)^2 \left( \frac{H_{\text{MB}}}{H} \right) \ln^2 \left( \frac{8\pi^2 t_b}{\omega_b} \right) \ln \left( \frac{\omega_b}{\max[T, \Delta^*]} \right) \sin \left( \frac{8t_b c}{e\mu_F H b^*} \right) \]

\[ \frac{\delta \chi}{\chi_0} = \frac{g_{\text{eff}}^2}{\pi^5} \left( \frac{\omega_b}{\mu_B H} \right)^2 \left( \frac{H_{\text{MB}}}{H} \right) \ln^2 \left( \frac{8\pi^2 t_b}{\omega_b} \right) \ln \left( \frac{\omega_b}{\max[T, \Delta^*]} \right) \cos \left( \frac{8t_b c}{e\mu_F H b^*} \right) , \tag{46} \]

where \( M_0 \) and \( \chi_0 \) stand for magnetic moment and magnetic susceptibility of noninteracting electrons.

Note that constants of “e-e” interactions in a Q1D conductor, \( g_1 \), are strongly renormalized above SDW(CDW) phase transition (see, for example, [53]) and \( |g_{\text{eff}}| \approx |g_1| \leq 1 \) at temperatures \( T \geq T_{\text{SDW(CDW)}} \). If we take experimental values of the parameters \( \omega_b, t_b, \) and \( H_{\text{MB}} (\omega_b \approx 3\mu_B H, t_b \approx 200 K, H_{\text{MB}} \approx 15 T) \), we can estimate the magnitudes of the forbidden oscillations at \( T \approx 6 \sim 10 K \) and \( H \approx 25 \sim 35 T \) as follows:

\[ \left| \frac{\delta M}{M_0} \right| \sim 0.1 , \quad \left| \frac{\delta \chi}{\chi_0} \right| \sim 1 \sim 10 \tag{47} \]

As it is seen from equation (47), the magnitudes of the forbidden oscillations are unexpectedly large. A period of the forbidden oscillations is given by equation (39). It corresponds to an area between open sheets of FS (see Fig. 4) that is in a sharp contradiction with Fermi liquid picture. We suggest a high field measurements of magnetic susceptibility and magnetic moment in a metallic state of (TMTSF)\(_2\)ClO\(_4\) to reveal this possible deviation from Fermi liquid behavior.

11. Conclusion: an Anomalous Metallic State

In the previous sections we showed that effective dimensionality of an electron gas is decreasing at the main “magic” directions of a magnetic field in Q1D conductors. We demonstrated that “e-e” interactions are dramatically changing with changing of the effective dimensionality. The explanations of such phenomena as MA’s and RMO were discussed in terms of these
unusual dimensional crossovers. Following similar ideas, we predicted the appearance of a strong forbidden magnetic oscillations. On the basis of the above mentioned unique properties of Q1D compounds, we formulate a concept of “anomalous metal” — a metal which magnetic properties are defined by magnetic field dependent many-body correlations.

Acknowledgments


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