Magnetoresistance of a Wigner liquid in a parallel magnetic field

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In recent years, low-density systems of two-dimensional electrons (holes) are intensively studied (see, e.g., reviews\cite{1}). That sort of systems exhibits a number of unusual properties. One of the main properties is a metal-insulator transition which occurs upon a decreasing of the carrier density. Moreover, with an increasing of a parallel magnetic field, a growth of resistivity and a metal-insulator transition are observed. The present paper is devoted to the theory of such systems in a parallel magnetic field.

The above-mentioned phenomena may have various reasons. For example, theoretically the effective mass $m^*$ could increase with spin polarization, but the experiments apparently do not manifest that. Further, interaction of carriers with impurities can increase with spin polarization. This effect was considered in the work\cite{2} where the interaction with charged impurities was discussed. However, the concrete realization of the idea in this work gives rise to doubts. The point is that, for the screened interaction with charged impurity, the used in the work formula is applicable only in the limit $r_S << 1$, while, on the contrary, the limit $r_S >> 1$ takes place for the considered systems ($r_S$ is the dimensionless average distance between the particles). In the work\cite{2} another mechanism connected with the presence of hole traps near the SiO$_2$/Si interface and their possible recharging was discussed. Though this mechanism is possible but its effect depends on a sample and changes from specimen to specimen, whereas experiments show more or less universal behavior for any utilized sample. Seemingly, it is necessary to find some universal mechanism.

Apparently, the universal reason exists because of some universal property of considered systems. Namely, a low-density system of two-dimensional electrons (under $r_S >> 1$) is the system with the strong correlations, i.e. it has a near order in particle arrangements (as in a Wigner crystal and, therefore, such a system is called a Wigner liquid sometimes). The use of this property is the main point of the present work.

First of all, it is necessary to do some remarks. In the limit $r_S >> 1$, there is the following hierarchy of energies. 1) The Coulomb energy is the greatest one and is taken as the unity. After all, it is this energy that results in a Wigner crystal and gives a near order in a Wigner liquid. 2) The energy of zero oscillations (or equivalently of typical plasma oscillations) is of the order of $1/\sqrt{r_S}$. Apparently, because of the zero oscillations the liquid do not crystallize right up to $r_S \approx 37$ (see the work\cite{3}). 3) At last, the Fermi energy is of the order of $1/r_S$ if this energy is estimated by using of band effective mass $m_b$ as in a Fermi gas. In reality, the Fermi energy is even smaller because the effective mass $m^*$ renormalized due to interaction is greater than $m_b$. Further, we shall deal with the energies of the order of the Fermi energy that is smallest one among the all energies and is of the order of exchange energy (see below).

The main idea of the present work is connected with the conception accepted in the work\cite{2}. In this work, an estimation of the exchange interaction of particles (and, accordingly, of the effective mass $m^*$) was done. For that, two neighboring particles interacting with external potential produced by environment and with each other via Coulomb were considered. Either of the two takes up its position mainly in its potential minimum, and from time to time they interchange their positions. With the help of this model, the exchange splitting of energy levels $(E_A - E_S)$ (its dependence on $r_S$) has been found. The exchange Hamiltonian that gives the splitting may be presented in the form:

$$H_{ex} = \lambda_S \left\{ (S_1S_2) + (Q_1Q_2) + \frac{4(S_1S_2)(Q_1Q_2) - 1}{4} \right\},$$

where the indices 1,2 are particle numbers, $S$ is the spin operator, and $Q$ is the quasi-spin operator that corresponds to the two-valley case ($Q = 1/2$). Further, the two-valley case as at SiO$_2$/Si inversion layer is examined. The constant $(-1/4)$ in the brackets is added for convenience. The value $\lambda_S$ and the effective mass $m^*$ are connected with the energy splitting in the following way:

$$2\lambda_S = E_A - E_S \sim \frac{n}{m^*},$$
where $n$ is the carrier concentration. The equation \textbf{(2)} gives estimation of $n^*$. Note that here the right-hand part is of the order of the Fermi energy of quasi-particles.

**Starting idea.** Further, one would think, it is necessary to operate in accordance with the Landau theory of the Fermi liquid (by using the effective mass estimation \textbf{(2)}). However, the question arises about adequacy of similar approach in our case. Actually, one may imagine a situation when, for example, a pair energy is smaller than the double Fermi energy of quasi-particles, and so a number of the pairs would arise together with the quasi-particles.

In order to elucidate the situation, let us turn to the physical picture which was implied in the works. In this work, the exchange energy of two nearest neighbors was calculated. If take into account all jumps of a particle in liquid, then one can see that the minimal energy of a particle can be even lower than in a pair, and this decrease is of the order of the exchange energy. Obviously, this minimal energy can be attributed to quasi-particle with zero momentum. So, as long as the all relevant to our problem energies are of the same order, then, in particular, it is permissible to compare a pair energy and the Fermi energy of quasi-particles. Therefore, the above-mentioned effect can exist and then influence of the pairs on the properties of the system should be taken into account.

One can consider instead of the pairs some other formations. For near order as in a Wigner crystal (i.e. as in triangular lattice), a cluster composed of three nearest neighbors is the most appropriate object. Obviously, energy of such a cluster (per a particle) is smaller than energy of a pair, and so, probably, it is necessary to do a preference for a three-particle cluster.

It is unlikely that one should consider more complicated objects in liquid because for that the ordered arrangement of particles is necessary not only for nearest neighbors but for next nearest neighbors as well. So further we shall consider the pairs and the threes only.

Evidently, the number of the clusters depends on a magnetic field (since clusters have spin) and temperature and the same relates to the number of mobile carriers of the Fermi type (fermions). It is natural to consider that the clusters are slow-moving formations because their jumping amplitude is smaller than for a particle. It means that they can be pinned even by small inhomogeneities and be excluded from the electric current. Moreover, there is an additional contribution to a resistance due to fermion scattering by the pinned clusters. Therefore, the resistivity of the system should depend on a magnetic field and temperature as well. The above picture is a basis for proposed description of the properties of the system.

In some extent, our approach is partly similar to the model used in the work. However, unlike this work, we appeal directly to the intrinsic properties of the strong correlated system rather than to the external factors.

**Clusters.** First of all, it is necessary to describe the properties of the clusters. The states and energies of the clusters are meant. Let $S$ and $Q$ be the total spin and quasi-spin of a cluster, accordingly. We start with a pair. The quantum numbers corresponding to its minimal energy are $S = 1$, $Q = 0$ or $S = 0$, $Q = 1$. With the help of \textbf{(4)}, it is easy to calculate the minimal energy $E_2$ of a pair with a result:

\[
E_2/2 = -3\lambda_S/8
\]

(here the energy per a particle is indicated).

Now consider a cluster of three particles. The state with $S = 3/2$, $Q = 3/2$ has the maximum energy and is of no interest for us. Passing to consideration of the states with $S = 3/2$, $Q = 1/2$ or $S = 1/2$, $Q = 3/2$ with the same energy $E_3'$, we obtain:

\[
E_3'/3 = -7\lambda_S/12
\]

(per a particle as well). These states are doubly degenerated for the fixed quantum numbers $S$, $S_3$, $Q$, $Q_3$. At last, for the energy $E_3''$ of a state with $S = 1/2$, $Q = 1/2$ we have:

\[
E_3''/3 = -3\lambda_S/4
\]

This state is not degenerated for the fixed quantum numbers $S$, $S_3$, $Q$, $Q_3$.

**Model.** Here we restrict ourselves to the simplest model. For the determination of the equilibrium properties, pinned clusters and a gas of mobile carriers (fermions) are considered without any interactions. An interaction of fermions with clusters is taken into account only for estimation of resistivity. The energy of the system is written as:

\[
E = \sum_{p,\sigma} \left[ \epsilon(p) + H\sigma \right] n_{\sigma}(p) + \sum_{\Sigma,\nu} \kappa(\Sigma) \left[ \gamma(\Sigma)E_0(\Sigma,\nu) + HS_3 \right] N(\Sigma,\nu) .
\]

Here the first sum relates to the fermions (the summations take places over the momentum $p$ and spin projection $\sigma = \pm 1/2$). The second sum relates to the clusters. The following notation for the set of quantum numbers is used:

\[
(S, S_3, Q, Q_3) \rightarrow \Sigma .
\]

In the second sum, the summations occur over these quantum numbers and state numbers $\nu$. The value $\gamma(\Sigma)E_0(\Sigma,\nu)$ is a cluster energy (with the quantum numbers $\Sigma$ and in the state $\nu$) and $\gamma(\Sigma)$ is number of particles in the cluster. The constant $\kappa(\Sigma)$ takes into account additional degeneracy of a level ($\kappa = 2$ for the states with the energy \textbf{(4)} and $\kappa = 1$ for all other states). Finally,

\[
H \equiv g^* \mu_B B .
\]

Here $B$ is the magnetic field, $g^*$ is the effective $g$ - factor (supposed to be the same for the fermions and clusters), $\mu_B$ is the Bohr magneton.
Further it is meant that only one cluster can be located in the every state \( \nu \). Apparently, the total number of these states is of the order of the total number of the particles. One would think, that in a liquid it is necessary to consider a cluster as a particle with enough big effective mass and characterize it by momentum. However, in our opinion, it is not unreasonable to describe the clusters by localized states because of imperfections.

**Equilibrium properties.** The cluster characteristics can be obtained with the help of statistical sum \( Z(\nu) \) for a given state:

\[
Z(\nu) - 1 = \sum_\Sigma \kappa(\Sigma) \exp \left\{ - \gamma(\Sigma) \left[ E_0(\Sigma, \nu) - \mu \right] + H S_\Sigma \right\}, \tag{7}
\]

where \( \mu \) is the chemical potential of the system.

The fermion distribution function has the usual form:

\[
n_\sigma(p) = \left\{ \frac{\epsilon(p) - \mu_\sigma}{T} + 1 \right\}^{-1}, \quad \mu_\sigma \equiv \mu - H \sigma. \tag{8}
\]

For quadratic dependence of the fermion energy on the momentum \( (\epsilon(p) = p^2/(2m^*) ) \), we have the following expression for the fermion concentration:

\[
n_\sigma = \frac{1}{V} \sum_p n_\sigma(p) = \frac{m^* T}{2\pi} \ln \left\{ 1 + \exp(\mu_\sigma/T) \right\} \tag{9}
\]

(for a fixed spin projection and for a fixed valley).

As usual, the chemical potential \( \mu \) is defined by the condition that the total number of the particles is equal to the sum of the numbers of fermions and electrons inside the clusters. In the case of two valleys, this condition looks as:

\[
n = n_F + n_\Gamma, \quad n_F = 2 \sum_\sigma n_\sigma, \quad n_\Gamma = \sum_{\Sigma, \nu} \frac{T}{Z(\nu)} \frac{\delta Z(\nu)}{\delta E_0(\Sigma, \nu)}. \tag{10}
\]

Here \( n \) is the total concentration of electrons, \( n_F \) is the concentration of mobile carriers (fermions), and \( n_\Gamma \) is the concentration of electrons inside the clusters (it is determined by a functional derivative of the function \( Z(\nu) \), see the expression (7)).

Up to now, the various expressions were written in general form that took into account the clusters containing the different numbers of electrons. Further, we restrict ourselves to clusters with three electrons for which the more low energies are obtained than for pairs (see (4) - (6)). For these clusters, it is necessary to take into account not only the states with lowest energy (4) but also the states with energies (4) because, in a sufficiently high magnetic field, a contribution of clusters with spin 3/2 is able the main. Furthermore, for simplicity, we neglect with dispersion of the cluster levels, i.e. we consider that every energy \( E_0 \) does not depend on the state number \( \nu \).

Instead of the designation \( E_0(\Sigma) \) for which it is necessary all the time to indicate a set of the quantum numbers, further it is convenient to use some other designations, namely:

\[
E_0(1/2, 1/2) \rightarrow \epsilon_0, \quad E_0(1/2, 3/2) = E_0(3/2, 1/2) \rightarrow \epsilon_1 \tag{11}
\]

(the arguments correspond to the spin and quasi-spin values). As appears from the expressions (11) and (10), \( \epsilon_0 < \epsilon_1 \).

After these remarks, the statistical sum (7) is written in the form:

\[
Z = 1 + 4 \exp \left[ -3(\epsilon_0 - \mu)/T \right] \cosh \frac{H}{2T} + 8 \exp \left[ -3(\epsilon_1 - \mu)/T \right] \left\{ \cosh \frac{3H}{2T} + 3 \cosh \frac{H}{2T} \right\}. \tag{12}
\]

For the concentration of electrons inside the clusters, we have:

\[
n_\Gamma = 3n_0 \frac{Z - 1}{Z}, \tag{13}
\]

where \( n_0 \) is the concentration of the cluster states.

Now let us find the spin polarization \( M \) depending on a magnetic field. Its ratio to the maximum possible value, \( M_m = n/2 \), can be written in the form:

\[
\frac{M}{M_m} = \frac{\eta}{3Z} \frac{2T}{\partial Z/\partial H} \mu + \frac{T}{2\epsilon_F} \left\{ \frac{H}{2T} + \ln \frac{\cosh \left[ \left( \mu + H/2 \right)/(2T) \right]}{\cosh \left[ \left( \mu - H/2 \right)/(2T) \right]} \right\}. \tag{14}
\]

Here the first term in the right hand side is the cluster contribution, and the second term is the fermion contribution. The constant \( \eta \) corresponds to the maximum part of the electrons which can be located inside the cluster states (apparently, one can think that \( \eta \) is less and of order of unity), \( \epsilon_F \) is the Fermi energy for the two-valley case at \( T = 0 \) and without taking into account any clusters:

\[
\eta = \frac{3n_0}{n}, \quad \epsilon_F = \frac{\pi n}{2m^*}. \tag{15}
\]

**Zero temperature.** Let \( \epsilon_0 < \epsilon_F \). In a magnetic field and at \( T = 0 \), it is necessary to take into account only the lowest levels of the clusters, namely:

\[
\epsilon_0 \rightarrow \epsilon_0 - (H/2)/3, \quad \epsilon_1 \rightarrow \epsilon_1 - H/2.
\]

The energy shifts for electrons are \( \delta \epsilon_{\uparrow, \downarrow} = \pm H/2 \). As a result, we have the following expressions for concentrations of fermions with the different spin projections and electrons inside the clusters (with minimal spin projections):

\[
n_\uparrow = \frac{m^*}{\pi} (\epsilon_0 - 2H/3),
\]

\[
n_\downarrow = \frac{m^*}{\pi} (\epsilon_1 - 2H/3).
\]
\[ n_\downarrow = \frac{m^*}{\pi} (\epsilon_0 + H/3) , \]
\[ n_\uparrow = n - \frac{m^*}{\pi} (2\epsilon_0 - H/3) ; \]
\[ \left( 0 < H < 3\epsilon_0/2 \right) . \]

In the greater magnetic fields,
\[ 3\epsilon_0/2 < H < 3(2\epsilon_F - \epsilon_0) , \]  
(a devastation up to zero of the cluster levels and an increase of \( n_\downarrow \rightarrow n \) take places (according to the linear laws). These behaviors are without taking into account the levels with the energy \( \epsilon_1 \) that is in the case \( \epsilon_1 > 2\epsilon_F \).

Now, let \( \epsilon_0 < \epsilon_1 < 2\epsilon_F \). The level with energy \( \epsilon_1 \) comes into the play at magnetic field \( H/3 > \epsilon_1 - \epsilon_0 \). There are the two cases. Firstly, this level can be connected up under \( n_\uparrow = 0 \), i.e. when the magnetic field falls into the interval \( [17] \). In this case, the cluster levels are devastated not completely (at the left border of the interval \( [17] \), i.e. under \( \epsilon_1 = 3\epsilon_0/2 \), these levels are not devastated at all). Secondly, this level comes into the play under \( n_\uparrow > 0 \), i.e. under \( \epsilon_1 < 3\epsilon_0/2 \). Let us discuss the latter case.

In the range of magnetic field \( 0 < H < 3(\epsilon_1 - \epsilon_0) \), there are the previous expressions \([10]\). In the point \( H/3 = \epsilon_1 - \epsilon_0 \), the states with initial energy \( \epsilon_1 \) are filled up instead of the states \( \epsilon_0 \) and, without taking into account a dispersion of the cluster levels, that takes place by a jump. In the range of magnetic field \( 3(\epsilon_1 - \epsilon_0) < H < \epsilon_1 \), we have:
\[ n_\uparrow = \frac{m^*}{\pi} (\epsilon_1 - H) , \]
\[ n_\downarrow = \frac{m^*}{\pi} \epsilon_1 , \]
\[ n_\uparrow = n - \frac{m^*}{\pi} (2\epsilon_1 - H) . \]

That continues up to the magnetic field \( H = \epsilon_1 \), i.e. up to the complete polarization, when the value of \( n_\uparrow \) equals to zero and the values of \( n_\downarrow \) and \( n_\uparrow \) become saturated. The dependence of any quantity, including a resistivity, on a magnetic field has a form of a broken curve, and for a magnetic moment even with jumps. It is clear that the curves are smoothed by temperature.

We note that the above expressions are valid until the obtained number of electrons inside the clusters is smaller than a maximal possible one, i.e., until \( n_\uparrow < 3n_0 \). That takes place under the condition
\[ 1 - \frac{\epsilon_1}{2\epsilon_F} < \eta . \]

Now about the magnetic moment. Let us consider the case \( \epsilon_0 < \epsilon_F \), \( \epsilon_0 < \epsilon_1 < 3\epsilon_0/2 \) (see \([16],[18]\) for the two intervals of a magnetic field). We obtain for the low magnetic fields:
\[ \frac{M}{M_m} = 1 + \frac{H - \epsilon_1}{\epsilon_F} , \]
\[ 3(\epsilon_1 - \epsilon_0) < H < \epsilon_1 . \]

The jumps of the magnetic moment take places at the points \( H = 0 \) and \( H = 3(\epsilon_1 - \epsilon_0) \). The ratio of susceptibilities in the two ranges of low and high magnetic field is
\[ \frac{\chi_1}{\chi_2} = \frac{5}{9} . \]

\[ 3\epsilon_0/2 < H < 3(2\epsilon_F - \epsilon_0) , \]

In considered case, the complete polarization takes place under \( H = \epsilon_1 < 3\epsilon_0/2 < 3\epsilon_F/2 \), while without the clusters under \( H = 2\epsilon_F \).

In connection with the jump of magnetic moment near the \( H = 0 \), it is necessary to note the following. In the work \([7]\), the direct measurements of the thermodynamic spin susceptibility were produced. It is revealed that magnetization behaves by nonlinear manner, the spin susceptibility becoming greater under magnetic field lowering. Possibly, that is the effect of the clusters.

**Resistivity.** For the resistivity \( \rho \), we use the ordinary expression:
\[ \rho = \frac{m}{n_\uparrow e^2 \tau} , \]

where \( \tau \) is the relaxation time. The question is what is necessary to take as the \( \tau \)? If an interaction of the mobile carriers with impurities is weak (not so with the clusters) or the number of the impurities is sufficiently small, then the relaxation time is proportional to an impurity concentration, i.e., for the scattering by clusters, \( 1/\tau \sim n_\uparrow \). That is the main for the sufficiently pure specimens. Perhaps, it is necessary to take into account the other scatterers too, then, in the simplest case, one can write:
\[ 1/\tau \sim n_\uparrow + \alpha n_i , \]

where \( n_i \) is the concentration of the other scatterers and the factor \( \alpha \) takes into account a difference of the scatterers. As a result, taking into account the expressions \([13]\) and \([10]\), one can write:
\[ \rho \sim \frac{\eta(Z - 1) + AZ}{(1 - \eta)Z + \eta} , \]

where \( A \equiv \frac{\alpha n_i}{n} \).

Let us consider the low temperatures \( T << \epsilon_F \). Under \( H = 0 \), from the expressions \([10]\) and \([13]\), for determination of the chemical potential, we have the equation:
\[ n - \frac{2m^*\mu}{\pi} = \frac{3\epsilon_0}{1(4/3) \exp [3(\epsilon_0 - \mu)/T] + 1} \]

(excluding the cluster states with the energy \( \epsilon_1 > \epsilon_0 \)). If one looks for the chemical potential in the form
\[ \mu = \epsilon_0 + \mu_1 , \]
then one can write the equation \[21\] as follows:

\[
1 - \frac{\varepsilon_0}{\varepsilon_F} = \frac{\mu_1}{\varepsilon_F} + \frac{\eta}{(1/4) \exp(-3\mu_1/T) + 1}. \tag{25}
\]

Under a partial filling of the cluster levels (for that it is necessary \(\varepsilon_0 < \varepsilon_F\)) and at \(T = 0\), \(\mu = \varepsilon_0\). And that is right also at \(T \neq 0\) for certain values of parameters, namely:

\[
1 - \frac{\varepsilon_0}{\varepsilon_F} = \frac{4\eta}{5}, \quad \mu_1 = 0. \tag{26}
\]

If that is the case then the cluster number is a constant. In this case, a resistivity does not change while the expression \[24\] is valid. So, we have the condition for a separatrix.

For small changes of the parameters, one can ignore the first term in the right-hand side of the \[25\], so that one obtains approximately:

\[
\frac{3\mu_1}{T} \approx -\ln\left\{\frac{5\beta_R}{\beta_L} - 4\right\}, \tag{27}
\]

where, for convenience, the designations \(\beta_L\) and \(\beta_R\) are used for the left-hand and right-hand sides of the \[25\], accordingly. From the \[27\], one can see that, under the condition \(\beta_R/\beta_L > 1\) \((\mu_1 < 0)\), the fermion number decreases and the number of scatterers increases, therefore the resistivity increases with temperature. Such behavior corresponds to a metallic phase. Otherwise, we have an insulator phase. In the both cases, the resistivity behaves in accordance with a linear law.

Resistance behavior with temperature depends on a value of a magnetic field. Let us consider a magnetic field in the interval given for the expression \[18\]. In this case, it is enough to take into account the cluster state only with the spin 3/2 and quasi-spin 1/2. The previous expressions \[24\] - \[26\], after the replacement

\[
\varepsilon_0 \rightarrow (\varepsilon_1 - H/2),
\]

remain valid. For example, instead of the \[26\] we come to:

\[
1 - \frac{\varepsilon_1 - H_S/2}{\varepsilon_F} = \frac{4\eta}{5}, \quad \mu_1 = 0, \tag{28}
\]

where \(H_S\) is the magnetic field corresponding to a separatrix. In view of \(H_S < \varepsilon_1\), we obtain the following condition for origin of separatrix:

\[
\frac{\varepsilon_1}{2\varepsilon_F} < 1 - \frac{4\eta}{5}.
\]

Under deviation from the \[28\], the expression for \(\mu_1\) coincides with the \[27\] but, in place of the previous left-hand and right-hand parts, it is necessary to substitute the corresponding parts of the \[25\]. With magnetic field, the left-hand part grows whereas the right-hand part is constant. Then, as follows from the \[27\], for the lower magnetic field, \(H < H_S\), we have the metallic phase (the resistivity grows with temperature) and for the greater magnetic field, \(H > H_S\), we have the insulator phase (the resistivity decreases with temperature increasing).

Physical meaning of the appearance of separatrix is simple. When a filling of the cluster states is small then it increases with temperature. Otherwise, for a large filling, it decreases with temperature increasing. The cluster filling grows with the magnetic field. At some magnetic field, an intermediate case can take place and the cluster filling is constant. So, a separatrix appears. Thus, when one draws a conclusion from the temperature dependance of resistivity then the metal - insulator transition can happen. Such a behavior of resistivity corresponds qualitatively to the experiments.

In the figure, the resistivity dependence on magnetic field is shown for different temperatures, namely: \(T/\varepsilon_F = 0.05\) (dots), \(T/\varepsilon_F = 0.1\) (solid line), and \(T/\varepsilon_F = 0.15\) (dashed line). The following parameters are taken: \(\varepsilon_0/\varepsilon_F = 0.95\), \(\varepsilon_1/\varepsilon_F = 1.1\), \(\eta = 0\), and \(A = 0\). The \(\rho_0\) is resistivity under \(T = H = 0\). The point of intersection of the curves corresponds to separatrix.

The conclusion about enhancement of the resistivity under influence of the magnetic field is the direct consequence of our model accounting for the presence of the clusters. The validity of the expressions \[24\] and \[25\] is assumed. That would be the case if the interaction of the mobile carriers with the clusters could be considered as weak. For strong interaction, the expression \[24\] can be used only in the limit of small number of the clusters, \(n_T \ll n\), possibly up to \(n_T \sim n\). However, for some values of the \(n_T\) (of the order of \(n\)), that simple picture can be incorrect because it is necessary to take into account the more complicated phenomena such as, e.g., localization of the fermions.

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FIG. 1: Resistivity dependence on magnetic field for different temperatures (see text)