Slow relaxation of magnetoresistance in AlGaAs–GaAs quantum well structures quenched in a magnetic field

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Abstract

We observed a slow relaxation of the magnetoresistance in response to an applied magnetic field in selectively doped p-GaAs–AlGaAs structures with a partially filled upper Hubbard band. We have paid special attention to excluding the effects related to temperature fluctuations. Although these effects are important, we have found that the general features of slow relaxation persist. This behavior is interpreted as related to the properties of the Coulomb glass formed by charged centers with account taken of spin correlations, which are sensitive to an external magnetic field. Variation of the magnetic field changes the numbers of the impurity complexes of different types. As a result, it affects the shape and depth of the polaron gap formed at the states belonging to the percolation cluster responsible for the conductance. The suggested model explains both the qualitative behavior and the order of magnitude of the slowly relaxing magnetoresistance.

1. Introduction

Recently [1] we reported observation of long time relaxation of (negative) magnetoresistance in p-type AlGaAs–GaAs quantum well structures where both wells and barriers were doped by Be. We argued that in these structures so-called A⁺-centers—doubly occupied acceptors belonging to the upper Hubbard band—are formed in the well. The observed long time behavior of magnetoresistance was explained by polaron effects involving the spin-correlated A⁺-centers.

In the present paper, we report more detailed studies of slow relaxations in the same samples. The main point requiring more detailed measurements is that in our previous paper [1] fluctuations of the sample temperature during measurements were out of proper control and account. Since temperature fluctuations may lead to a variation of the sample resistance our previous results and their interpretation need a proper verification. We are grateful to Ovadyahu who attracted our attention to this problem.

3 Throughout the paper, we assume that the charge carriers are holes. Then a A⁺-center is an acceptor atom occupied by two holes. Therefore it is charged positively. Similarly, a singly occupied (by holes) acceptor atom is neutral (A⁰-center), while a bare atom is negatively charged (A⁻-center).

Investigation of the role of the temperature fluctuations is the main goal of the present work. We will show that, although the temperature fluctuations indeed can produce a pronounced effect on the resistance, the observed long time relaxations of magnetoresistance have an independent source. More careful analysis of the experimental results has required, however, some modification of the theoretical model. Now we believe that the observed slow relaxation is due to a direct influence of the magnetic field on the aggregates responsible for formation of the polaron gap at the sites responsible for the conductance [2]. This mechanism turns out to be more important than the magnetic-field-induced shift of the chemical potential of the holes [1].

The paper is organized as follows. In section 2 we briefly describe the samples and experimental procedure and report the experimental results. These results are interpreted in section 3 where the theoretical models are considered and compared with experiment.

2. Samples, experimental procedure, and results

We used the structures containing 10 GaAs quantum wells separated by Al₀.₃Ga₀.₇As barriers. Thickness of both wells...
and barriers was 15 nm. Confining Al$_{0.3}$Ga$_{0.7}$As layers had a thickness of 20 nm. The growth procedure is described in detail in [1]. Central regions with a thickness of 5 nm of both wells and barriers were $p$-doped with Be (concentration $10^{17}$ atoms cm$^{-3}$). The contacts were made by 2 min burning at 450°C in deposited gold containing 3% of Zn. The samples were shaped as Hall bars. The resistance was determined from the voltage between the voltage probes at a fixed current of 1 nA. We studied several samples cut from the same wafer.

The samples were relatively low-Ohmic ($10^5$–$10^6$ Ohms/□ at 4 K) that is, in our opinion, due to the fact that the impurity band formed by $A^+$-centers is rather close to the valence band, the binding energy being $\sim 10$ meV [1]. We have checked that for all our measurements the $I$–$V$ curves were linear in the temperature domain 4.2–1.35 K. The temperature dependence of the resistance is compatible with the 2D Mott variable range hopping (see figure 1 in [1]), $R \propto \exp(T_0/T)^{1/3}$, with $T_0 = 1000$ K for the sample under investigation. The magnetoresistance curves were close to those reported in [1], namely, the linear negative magnetoresistance crossed over to the quadratic positive magnetoresistance at stronger magnetic fields. This behavior agrees with theoretical predictions [3].

Studies of the slow relaxations induced by controlled variation of external magnetic field are specifically difficult because our thermometers are sensitive to magnetic field. Therefore, to investigate the role of the temperature fluctuations we made special measurements of the sample resistance and its temperature in the absence of magnetic field. The corresponding curves are presented in figure 1.

One concludes that temperature fluctuations up to several mK indeed exist. We believe that these fluctuations are due to fluctuations of the pressure in the system pumped by a pre-evacuation pump without a sufficiently large damping volume, which should be at least by an order of magnitude larger than the Dewar volume. It is also seen that the resistance fluctuates keeping time with the temperature fluctuations—there is no lag between variations of the resistance and temperature. This fact allows us to conclude that at a given time the resistance of the sample is controlled by the temperature of the sample measured at the same instant of time. Accordingly, we can extract temperature derivative of the resistance and then estimate the corresponding fluctuation-induced contribution to the zero-field resistance at a known temperature as

$$\Delta R \approx \frac{\partial R}{\partial T} \Delta T. \quad (1)$$

From figure 1 it is seen that the variation of temperature by 1 mK leads to a variation of voltage and thus the variation of resistance is of the order of 0.2%.

Shown in figure 2 is variation of resistance (or rather of the voltage between the probe contacts) obtained by ramping of magnetic field in time with very slow sweep rate (about 30 Oe s$^{-1}$). Large circles show measured temperatures at $H = 0$, that correspond to three instants of time in figure 2. Note that the magnetoresistance curves of figure 2 were close to those reported in figure 1 of our previous paper [1], which were obtained at a fast magnetic field sweep of $\sim 200$ Oe s$^{-1}$. In particular, the linear negative magnetoresistance crossed over to the quadratic positive magnetoresistance at stronger magnetic fields. As it is seen from figure 2, the resistance at the successive instants of time when $H = 0$ is different—it gradually decreases with time (by $\approx 1\%$). At the same time, the temperature at those instants is almost the same showing an initially weak decrease and then a weak increase within the interval not exceeding $\approx 1$ mK. These deviations of temperature would lead only to deviations of the resistance by $\approx 0.2\%$ (upward and downward). These measurements demonstrate that, in addition to thermal effects, the resistance tends to decrease after application of the magnetic field, the effect at the times $\sim 1000$ s being at the level of 1%.

In figure 3, we present the behavior of the voltage across the probe contacts (and thus of the resistance) for two opposite

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$^4$ In [1] we erroneously estimated the sweep rate as 40 Oe s$^{-1}$. We apologize for the mistake.
current directions for the sample quenched during the time \( \tau_q = 1200 \) s in a steady magnetic field of 0.7 T. The resistance was measured after switching off the field. To prevent heating, the magnetic field was switched on and off at a relatively slow rate, during the time interval of \( \tau_{sw} = 250 \) s. The voltage drop at \( t = 0, U_0 \equiv U(t = 0) \), is subtracted. As is seen, the resistance gradually increases. One can check that it tends to its value at \( H = 0 \). In this case the temperature measurements are available during all the time of the measurement. Thus we are able to correct the resistance curve for the temperature fluctuations at all experimental points. It is seen that the corrected voltage, \( U^* \equiv U - U_0 - (\partial U/\partial T) \Delta T \), slowly relaxes to higher values and is saturated at times \( \sim 1000 \) s, which is close to \( \tau_q \). It is important that the saturation value is equal to the resistance of the sample before application of the magnetic field. In other words, after application of the magnetic field the resistance slowly decreases with time, but after switching off the field it is gradually restored to its initial value. Another important point is that the restoration time is nearly equal to the time at which the sample was subjected to the external magnetic field. Note that figure 2 contains both instantaneous reaction of the resistance to magnetic field and a slowly relaxing part. Contrary, figure 3 is recorded after switching off the magnetic field. Therefore, it contains only the slowly relaxing part, which is compatible with figure 2.

3. Discussion

We believe that the results presented above evidence a specific slow response of the sample resistance to applied magnetic field. This response cannot be explained by temperature fluctuations, which would not lead to the observed monotonous in time contribution to the resistance. As it follows both from experiment and our estimates, a steady heating due to eddy currents induced in the Cu substrate by variations in the magnetic field is negligibly small. At the same time, the observed resistance gradually decreases after application of magnetic field and tends to restore after switching it off. In addition, as it follows from direct measurements and our estimates, a possible effect of temperature fluctuations would be about an order of magnitude less than the observed variation of resistance.

3.1. Theoretical model

Hole–impurity complexes. Let us first review the previously suggested mechanism of slow relaxations [1], which is a generalization of the one developed in [2]. We assume that some localized states (including \( A^+ \)-centers—acceptor atoms doubly occupied by holes and located in the wells) form bistable aggregates. The low energy states of these aggregates are almost degenerate, the transitions between the states can take place only due to many-electron processes that results in very long transition times. Although the above aggregates do not belong to the percolation cluster, they can be polarized by the electrical charges located at the hopping sites. The polaron clouds formed by the aggregates for typical hopping sites lead to formation of a polaron gap at the Fermi level. As a result, conductance decreases.

An important feature of the material under consideration is that, in addition to doubly and singly occupied acceptors in the well (\( A^+ \)-centers and \( A^0 \)-centers, respectively), there can also exist neutral complexes consisting of a negatively charged acceptor in the barrier and a valence-band hole located in the well, which is coupled to the acceptor due to Coulomb interaction, figure 4. We will call such complexes \( A^0 \)-centers. Analog of such complexes in a n-type conductor were introduced in [4], and in what follows we will exploit the scheme suggested in that paper. In the material under consideration, one has to consider acceptor complexes of four types: (i) positively charged \( A^+ \)-centers—doubly occupied (by holes) acceptors located in the wells and forming the upper Hubbard band; (ii) neutral \( A^0 \)-centers—singly occupied (by holes) acceptors located in the wells; (iii) negative \( A^- \)-centers—acceptors in the barriers; and (iv) neutral \( A^0 \)-complexes consisting of an acceptor in the barrier coupled to a valence-band hole in the well. Note that the localization length of the \( A^+ \)-centers is around 10 nm, which is less than the typical distance between the centers (about 30 nm for relevant concentrations); the localization length of the \( A^0 \)-centers having relatively small binding energy is expected to be close to that of the \( A^+ \)-centers.
Assuming that the dopant concentrations, $N_A$, within the barrier and within the well are equal one has

$$N_{A^+} + N_{0^+} = N_{A^-} + N_{0^-} = N_A. \quad (2)$$

Due to charge conservation we also have

$$N_{A^+} = N_{A^-}. \quad (3)$$

In combination with equation (2) we get

$$N_{0^+} = N_{0^-}. \quad (4)$$

Since there are three unknown concentrations, $N_i$, one needs one more relation in order to calculate the chemical potential, $\mu$. To derive this relation let us take into account the fact that a hole provided by an acceptor within the barrier is inevitably captured within the well. However, it can belong either to an $A^+$ or it can exist as a valence-band hole coupled to a negative acceptor in the barrier forming an $\tilde{A}^0$-center. In particular, a hole released by an acceptor in the barrier can be captured by an occupied acceptor in the well forming an $A^+$-center provided

$$U_{A^+} + \frac{e^2}{\kappa \sqrt{r^2 + (d_w + d_b)^2/4}} \geq U_{\tilde{A}^0}. \quad (5)$$

Here $U_i$ is the binding energy of the center of $i$th type, $r$ is a distance between the two acceptors along the plane of the structure, $d_w$ and $d_b$ are the thicknesses of the well and the barrier, respectively. For simplicity we assume that $kT \ll U_i$ and that the dopants form $\delta$-layers in the middle of the well and the barrier. The lhs of equation (5) gives the energy gain due to formation of an $A^+$-center (the second term describes the interaction of the $A^+$-center with the $A^-$-center within the barrier). One concludes that if $U_{\tilde{A}^0} < U_{A^+}$ all the acceptors within the barrier will become the $A^+$-centers (while the acceptors within the barriers all become the $A^-$-centers).

In contrast, if

$$U_{\tilde{A}^0} - U_{A^+} \geq \frac{2e^2}{\kappa (d_w + d_b)} \quad (6)$$

both $A^+$-centers and $\tilde{A}^-$-centers cannot be formed, and only $A^0$- and $\tilde{A}^0$-centers are present. Strictly speaking, in both of these limits the charge transport over the impurity band cannot take place. Such a transport requires partial ‘compensation’ of the acceptors, which is possible provided

$$0 < U_{\tilde{A}^0} - U_{A^+} < \frac{2e^2}{\kappa (d_w + d_b)}. \quad (7)$$

In this case the relation (5) defines the average distance between $A^+$-centers and thus the relative densities of states of the upper Hubbard band (formed by $A^+$- and $A^0$-centers) and of the lower Hubbard band (formed by $\tilde{A}^0$- and $A^-$-centers), $g_a$ and $g_b$, respectively. This distance, $r_c$, corresponds to the equality relation in equation (5).

**Influence of magnetic field.** The external magnetic field tending to align all the spins of the holes decreases the binding energy of the $A^+$-centers, $U_{A^+}$, thus changing the balance between $A^+$- and $\tilde{A}^0$-centers. As is known for $A^+$-centers in GaAs/AlGaAs structures, the total spin of the $A^+$ state is 2 while the spin of a separate hole is 3/2 (see [5]). Thus the application of the magnetic field $H$ leads to an effective increase of the Hubbard energy by the Zeeman energy term $\mu_B g |H|$, $\mu_B$ being the Bohr magneton. As a result, the holes redistribute between the acceptors and the chemical potential

\[ \text{Figure 4. On the impurity complexes in our quantum well structure.} \]

(This figure is in colour only in the electronic version)
is shifted as [6]
\[
\delta \mu \sim \mu g[H] \left( \frac{g_a - g_l}{g_a + g_l} \right).
\]
(9)

**Estimate for magnetoresistance.** To compare the experimental results with our theoretical model let us reconstruct the estimate [1] for magnetoresistance. Consider a hopping site with equilibrium energy \( \varepsilon \) (referred to as the Fermi level) coupled to some two-level system (TLS) with inter-level splitting \( E \) in the equilibrium. The coupling potential \( U(R) \) (where \( R \) is a distance between the hopping site and the TLS) is defined as the difference between the coupling energy for the two TLS states corresponding to the upper and lower levels of the TLS. If one creates an excitation at this hopping site (either an electron or a hole depending on the sign of \( \varepsilon \) the TLS splitting is changed as \( E \to E + U(R) \). If \( U \) is negative and \( |U| > E \), the TLS changes its state with respect to the equilibrium one. Correspondingly, the excitation energy is also changed as \( \varepsilon + U(R) \), and for an electron excitation its energy is lowered. For the hole excitation the same effect will take place if \( U > 0 \). As a result, the presence of a TLS leads to a formation of a polaron gap around the Fermi level having a width of \( 2U \) and depth depending on the concentration of TLSs. The shape of the gap can be found as follows [2]. Specifying the density of states for the TLSs with inter-level spacing \( E \) and relaxation time \( \tau \) as \( P(E, R, \tau) = \bar{P}/\tau \) where \( \bar{P} = \text{const} \) we find the distribution function of polaron shifts, \( U \), from the equation as
\[
\mathcal{F}(U, \tau) dU = 2\pi R dR \int_0^{U(R)} dE \frac{dE}{\tau} P(E, R, \tau) S(t, \tau).
\]
(10)

Here \( S(t, \tau) \) depends on the measurement protocol. \( \mathcal{F}(U, \tau) \) is just the probability to find a TLS providing the polaron shift between \( U \) and \( U + dU \) and located within the layer between \( R \) and \( R + dR \), at time \( t \).

To discuss the results shown in figure 3 we have to specify the function \( S(t, \tau) \) for our measurement protocol. Let us start with the situation without magnetic field, \( H = 0 \), and consider a chessboard aggregate located close to a hopping site. It has some equilibrium configuration, which depends, in particular, on the aggregate’s neighbors. Now we switch on the magnetic field and wait during the time \( \tau_q \). Since the site energies get changed the aggregate deforms to a new equilibrium state according to the law \( (1 - e^{-t/\tau}) \). At \( t = \tau_q \) the probability to reach a new equilibrium state corresponding \( H \neq 0 \) is then \( (1 - e^{-\tau_q/\tau}) \). Since in the new state some aggregates can be destroyed, this factor characterizes a relative decrease of the polaron effect. As a result, application of the magnetic field leads to a decrease in the resistance. After switching off the magnetic field the aggregate returns to its equilibrium state at \( H = 0 \) according to the law \( e^{-t/\tau} \). Therefore, the polaron effect for the corresponding configurations is restored and the resistance is increased. Correspondingly, at time \( t \) the aggregate regains its initial equilibrium state (supporting the initial resistance) with the probability
\[
S(t, \tau) = (1 - e^{-\tau_q/\tau}) e^{-t/\tau}.
\]
(11)

Summing over all relevant configurations one obtains for the slowly relaxing contribution to the distribution of polaron shifts
\[
F(U) \propto - \int_{\tau_{\text{max}}}^{\tau_{\text{min}}} dt P(t)(1 - e^{-t/\tau}) e^{-t/\tau}.
\]
(12)

Here \( \tau_{\text{max}} \) is the maximal relaxation time of the aggregates. In the following we will take into account that \( \tau_{\text{max}} \gg \tau_q, t \) and \( P(t) = P/\tau \propto 1/\tau \). Then
\[
F(U) = -2\pi \mathcal{G}(t, \tau_q) \bar{P} \int_U^{\infty} dU' \frac{dP}{dU'} \mathcal{G}(U, \tau) \approx \int_{\tau_{\text{max}}}^{\tau_{\text{min}}} dt \frac{dP}{dU'}(1 - e^{-t/\tau}) e^{-t/\tau}.
\]
(13)

If both \( \tau_q \) and \( t \) significantly exceed \( \tau_{\text{min}} \), then an excellent approximation for \( \mathcal{G}(t, \tau_q) \)
\[
\mathcal{G}(t, \tau_q) \approx \ln \left( 1 + \frac{\tau_q}{\tau} \right).
\]
(15)

In this approximation the relaxation depends only on the ratio \( \tau_q/\tau \). As was shown in [7], it excellently fits numerous data on ageing in glasses.

Following our previous paper [2] let us consider many-electron aggregates (electronic TLSs) of chessboard type formed from the pairs with one occupied and one empty site. These building blocks can be regarded as pseudo-spins. The two states of the aggregate correspond then to the opposite direction of all of the ‘spins’. The transition between the two states can be due to either a coherent multi-electron hop, or due to motion of a ‘domain wall’ separating parts of the aggregate with different phases of spin orientation. For this model, the size of the optimal aggregate turns out to be relatively large and its coupling to hopping sites is most effective when the hopping site is close to one of the sites belonging to the aggregate. In this case the coupling is just due to Coulomb interaction between the charges of the hopping site and the nearest site of the aggregate. The coupling to the rest sites of the aggregate turns out to be much weaker. The fact that the site belonging to the hopping cluster is close to the aggregate does not imply rapid exchange of particles between the hopping cluster and the aggregate because of the polaron effect. If the occupation of this hopping site leads to formation of the polaron, then in order to change this occupation number one should cover the energy deficiency equal to the polaron gap. Another point is that the aggregates are formed from the sites of the same type as those forming the hopping cluster. However, the typical aggregates do not cause any delocalization of the holes over the whole sample. Otherwise the conductance would be metallic.

Since any site belonging to the aggregate can be coupled to a hopping site, the distribution function \( \bar{P} \) should be multiplied by the number \( N \) of the sites belonging to the aggregate since the hopping site can be coupled to any site within the aggregate. In what follows we will absorb this factor into \( \bar{P} \). For the Coulomb charge–charge interaction, \( U = e^2/2 \), the integrand in equation (13) is \( \propto U^{-2} \), and, correspondingly, the integral is controlled by its lower limit. The apparent divergence at \( U \to 0 \) has a cut-off due to the fact that the small
values of interaction energy $U$ correspond to large distances between the hopping site and the aggregate. At large distance the Coulomb charge–charge coupling is replaced by a much weaker dipole coupling originating from (a random) dipole moment of the aggregate. Having this fact in mind we assume that the cut-off as the energy $\varepsilon_h$ corresponding to the typical coupling energy between the sites separated by the typical hopping length $r_h$. The quantity $\varepsilon_h$ is just the width of the hopping energy band. Following this reasoning we assume that the gaps with $U < \varepsilon_h$ do not effectively influence the hopping transport. As a result, the density of states near the Fermi level as a function of energy can be described as

$$\frac{\delta g_d}{g_d} \sim \frac{2\pi \bar{P}}{\sqrt{\varepsilon^2 + \varepsilon_h^2}} \left( \frac{c^2}{\kappa} \right)^2 G(t, \tau_0). \quad (16)$$

Contrary to the model [2] in the present case the external perturbation is due to variation of the magnetic field. In particular, an instant application of magnetic field leads to a shift of the chemical potential within the polaron gap. Let us estimate the corresponding change in conductance assuming that

$$\delta \mu \lesssim \mu_B R \frac{H}{L} \ll \varepsilon_h.$$ 

From equations (16) and (9) one can expect that an energy shift of $\delta \varepsilon \sim \delta \mu \lesssim \mu_B g|\bar{P}|$ will lead to a relative shift in the density of states, as well as in the conductance, of the order of $(\delta \mu/\varepsilon_h)^2 = (\mu_B g |\bar{P}|/\varepsilon_h)^2$. Therefore, one can expect that the magnetoresistance in an oscillating magnetic field would form a pattern following the magnetic field. In addition to such a pattern, the experiments demonstrate a gradual increase of the conductance with time.

To explain this increase we introduce an additional mechanism, which is due to a slow reconstruction of the polaron gap adjusting its depth to the magnetic-field-dependent number of the ‘active’ sites forming the polaron cloud. The reconstruction can result from a direct influence of the magnetic field on the aggregates forming the polaron clouds. Indeed, the aggregates are formed from pairs of sites allowing electron (hole) transitions within the pairs. These transitions change types of the impurity complexes as $(A^+, A^0) \rightarrow (A^0, A^+)$; $(A^+, A^-) \rightarrow (A^0, A^0)$, and $(A^0, A^0) \rightarrow (A^-, A^+)$. Thus one concludes that the probability for the first configuration to be included into the aggregate is given by the product $N_{A^+} N_{A^0}$ while for the second configuration by the product $N_{A^+}^2 N_{A^0}^2$. In the region specified by equation (8), the magnetic field changes the relation between numbers of the $A^+$ and $A^-$ centers. The effect of magnetic field on the hopping parameter $g_d$ (related to $N_{A^+}$) and $g_l$ (related to $N_{A^0}$) can be estimated from equation (5) for $r = r_e$.

Since the magnetic field changes the binding energy of an $A^+$-center, $\delta U_{A^+} = \mu_B g H$, it changes the characteristic distance, $r_c$. Since $N_{A^+}/N_A \sim r_c^2 N_A$, this change leads to a variation in the number of $A^+$: $\delta N_{A^+}/N_{A^+} \sim r_c^2/2N_A$. Having this fact in mind and using equation (5) one obtains:

$$\frac{\delta N_{A^+}}{N_{A^+}} \simeq 2\delta U_{A^+} \frac{r^2 + (d_w + d_h)^2/4}{r^2 \varepsilon^2} = \frac{\delta U_{A^+}}{U_{A^+} - U_{A^-}}. \quad (17)$$

As was noted above, the concentration of the aggregates is controlled by the product $N_{A^+} N_{A^0} = N_{A^+} (N_A - N_{A^+})$. The variation of this product, in turn, is

$$\delta N_{A^+} / (N_A - 2N_{A^+}). \quad (18)$$

Therefore, one concludes that the effect of magnetic field depends on the sign of $N_{A^+} - 2N_{A^-}$. If this sign is positive, that is the concentration of $A^+$-centers is relatively small, the concentration of the aggregates decreases with an application of magnetic field (decreasing $N_{A^-}$). In turn, it leads to a suppression of the polaron effect and to a decrease of the resistance. If the concentration of $A^+$-centers is large, the magnetic field can lead to an increase of the resistance. Since the experiment demonstrates a decrease of resistance as a result of the application of the magnetic field, we conclude that in our case the density of $A^0$-centers exceeds the density of $A^-$-centers. This assumption is supported by the fact that in our experiments the typical distances between acceptors, $N_A^{-1/2}$, were relatively large, at least $N_A^{-1/2} > (d_w + d_h)$, which is in favor of creating $A^0$-centers rather than $A^+$-centers.

The relative decrease of the density of states of the aggregates, $\Delta P/\bar{P} \simeq \delta N_{A^+}/N_{A^0}$, can be estimated from equation (17) as

$$\frac{\Delta P}{\bar{P}} \simeq \frac{2\mu g |\langle H \rangle|}{U_{A^0} - U_{A^+}}. \quad (19)$$

where $|\langle H \rangle|$ is the time average of the absolute value of the magnetic field. This ratio is equal to the relative (with respect to the total number of sites experiencing the polaron effect) number of the hopping sites where the polaron cloud is destroyed by the magnetic field. It is important that this effect is linear in $|\langle H \rangle|$ and can dominate over the quadratic effect mentioned above.

For further estimates we have to specify the TLS distribution function. A crude estimate for $\bar{P}$ in the case of electronic aggregates is [1]

$$\bar{P} = \frac{N e^{-\lambda N}}{(N^0)^2 N^{1/2} (\varepsilon^2/\kappa \rho)} \quad (20)$$

Here $N$ is the number of pairs of sites forming the aggregate, the exponential is a statistical factor describing the probability to construct the bistable aggregate, $\rho$ is the typical distance between the sites forming the aggregate. Here the first factor in the denominator describes the typical volume of the (2D) aggregate while the second—typical scatter of the TLS energy splitting. The factor $\lambda$ depends on the competition between the Coulomb interactions within the system and scatter of single-particle energies. Indeed, for weak Coulomb interactions the system is in its ground state and occupation of all single-particle states is given. It is the Coulomb correlations that allow to have metastable configurations with close total energies. Unfortunately, the number $\lambda$ is not known; it can depend on realization of the Coulomb glass. However one expects that large relaxation times are available at not too large $N$ and thus the exponential is not too small. Assuming $N \approx 5$, $e^{-\lambda N} \approx 0.1$, $\rho = \xi a \approx 100$ nm (that is of the order of the typical hopping length) one estimates $\bar{P}$ as $\sim 10^{23}$ cm$^{-2}$ erg$^{-1}$. 
Substituting this estimate in equations (16), (19) and assuming that, \( \mu g(H)/(U_A^0 - U_A^+) \sim 3 \times 10^{-2} \) one obtains

\[
\delta G/G \approx \delta g_u/g_u \sim 0.003 \mathcal{G}(t, \tau_g).
\]

(21)

Solid lines in figure 3 show fitting of the experimental results by the expression

\[
\delta U^*/U^* = \pm 0.003 \mathcal{G}(t + 0.6 \tau_{sw}, \tau_g)
\]

(22)

(we have introduced the time shift of 0.6 \( \tau_{sw} \) to allow for non-instantaneous switching off the magnetic field). Although the fitting cannot be considered as a quantitative verification of the suggested mechanism (large error bars, very crude theoretical model, etc) we conclude that the introduced mechanism suggests a reasonable interpretation of the experimental results.

Note that the estimate (21) differs from that of [1] where the effect of magnetic field was attributed to partial suppression of the polaron gap by the magnetic field to an increase of the energy of the upper state of an aggregate. The suppression is due to the influence of the magnetic field on the \( \Lambda^+ \)-centers making some configuration inaccessible even with account of the correlation energy. However, we underestimated a possible effect of an increase of the energy of the lower state involving \( \Lambda^+ \)-centers, which can compensate the effect related to the increase of the energy of the upper level. In this way, we overestimated the slow-relaxing part of the conductance. On the other hand, the experimental result for this quantity was also underestimated because it was not corrected for the temperature fluctuations. We believe that the corrected experimental results and the suggested theoretical model are consistent. Namely, the model explains the experimentally observed gradual increase of conductance with time.

Unfortunately, we do not have sufficient information to fit the results for periodic ramping of the magnetic field shown in figure 2. Indeed, the reliable data are available only for three points where \( H = 0 \). Applying the considerations similar to those discussed above one would expect a crude estimate of the slowly relaxing part similar to equation (21) with replacement \( \mathcal{G}(t, \tau_g) \rightarrow \ln(4t/\tau_g) \) if the logarithm is large. Here \( \tau_g \) is the period of magnetic field ramping. This estimate does not contradict the experimental data.

Thus we have discussed two different mechanisms of slow relaxations (of magnetoresistance) in response to time-dependent external magnetic field, both related to the polaron effect. The first one is to some extent similar to the effect of the gate voltage [8–13]. It is induced by the magnetic-field-driven shift of the chemical potential. As was noted above, this mechanism does not explain the observed gradual increase of conductance in time. We expect that this mechanism will be more pronounced and even dominant in the situation when the induced shift is larger than the width of the hopping energy band.

The second effect is related to a direct dependence of the ‘active’ sites responsible for the polaron effect on the magnetic field. According to our estimates, it is this effect that is responsible for the experimentally observed slowly relaxing response.

It is worth noting that the slow relaxations in response to variations of the gate voltage are usually not observed in doped crystalline semiconductors—the experiments [8–13] were performed using samples with a significant amount of disorder. We believe that there are two reasons for that: (i) the TLS-induced polaron effects are much weaker than the direct influence of the gate voltage on the DOS; (ii) relatively large sweep rates of the gate voltage probably lead to pronounced non-equilibrium behaviors. In the present experiments, the direct influence of the magnetic field sweep is much weaker, and therefore the effects induced by TLS polarons can be observed. At the same time, glassy behavior in two-dimensional Si structures was reported in [14] (see also references therein).

To conclude, we observed a slow relaxation of magnetoresistance in response to applied magnetic field in selectively doped p-GaAs–AlGaAs structures with a partially filled upper Hubbard band. We explain this behavior as related to the properties of the Coulomb glass formed by charged centers with account of spin correlations, which are sensitive to the external magnetic field.

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