Magnetoresistance of two-dimensional p-GaAs/$Al_{10.3}Ga_{10.7}As$ structures in the vicinity of metal-insulator transition: Effect of superconducting leads

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Experimental and theoretical studies on transport in semiconductor samples with superconducting electrodes are reported. We focus on the samples close to metal-insulator transition. In metallic samples, a peak of negative magnetoresistance at fields lower than critical magnetic field of the leads was observed. This peak is attributed to restoration of a single-particle tunneling emerging with suppression of superconductivity. The experimental results allow us to estimate tunneling transparency of the boundary between superconductor and metal. In contrast, for the insulating samples no such a peak was observed. We explain this behavior as related to properties of transport through the contact between superconductor and hopping conductor. This effect can be used to discriminate between weak localization and strong localization regimes.

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I. INTRODUCTION

Recently we reported an observation of crossover from strong to weak localization in 2D p-GaAs/AlGaAs structures.[1] The magnetoresistance of our samples in weak localization regime has demonstrated, in addition to the standard antilocalization behavior, a small peak of negative magnetoresistance (NMR) at weak fields $<0,02$ T. The nature of this peak was not clear. It could not be attributed to weak localization since at higher magnetic fields the samples clearly demonstrated positive magnetoresistance related to antilocalization. Since the peak disappeared at the critical temperature of In contacts, 3.4 K, it was natural to assume that it originates from the superconducting contacts. However, simple considerations would predict PMR at magnetic field destroying superconductivity of the contacts. Another point was that the effect was not observed for the hopping regime even when the resistance of the samples was not much larger than the resistance of the metal samples. Thus, further analysis of this effect was necessary.

Usage of superconducting contacts is a common practice in studies of the samples close to the metal-insulator transition (MIT) or deep in the hopping regime. However, though a significant attention was paid to the properties of a superconductor–normal-metal interface, we are not aware of studies of the interface between superconductor and the sample close to MIT transition.

Here we present results of experimental and theoretical studies of a role of superconducting contacts to structures in the vicinity of MIT from both sides of the transition. We will prove that the peak of the magnetoresistance mentioned above is related to a presence of an insulating barrier between the superconductor and semiconductor. This barrier suppresses the Andreev reflections, the single-particle channel being affected by the superconducting gap. Thus the mechanism of the observed negative magnetoresistance is suppression of the gap by the magnetic field. We will show that this model allows one to explain the experimental results in detail.

The samples on the dielectric side of MIT do not demonstrate any traces of the effect—even with an account of the sample resistance increase. We will show that this fact is related to specific properties of the hopping transport including a larger value of the effective energy band than for metals and the topology of the percolation cluster. We believe that the unusual magnetic field dependence of electron transport in systems with superconducting electrodes can be used as a tool to discriminate between the regimes of weak and strong localization.

II. EXPERIMENT

We have chosen $GaAs/Al_{0.3}Ga_{0.7}As$ multiwell structures with the well widths (d) of 10, 15, and some larger barrier width of 25 nm, doped by an acceptor dopant Be. The binding energy of Be dopant is $E_A \sim 28$ meV which yields the localization length $a_b = 2$ nm being much less than (d). The method of growing multilayer structures by molecular beam epitaxy was described in our work.[2] In sample 2, by selective doping of the central regions of the wells with relative widths 1/3 we prepared the system where the lower Hubbard (LH) impurity band was formed. In the sample 1,3 by selective doping of the central regions of both wells and barriers (with equal doping concentrations) we prepared a system where the upper Hubbard (UH) impurity band was partly occupied in the equilibrium. The dopant concentration $N_a \sim (0,6-2) \times 10^{12}$ cm^{-2} was near the critical concentration for the metal-insulator transition in 2D structures $N_c^{1/2} a_b \sim 0,3$. The contacts with a form of drops (with a diameter 0.5 mm) were produced by firing indium with a low zinc concentration during 2 min at temperature 450 °C. Magnetoresistance was measured with a help of four probes method; magnetic field was applied normal to the plane of the structure.

Close to the room temperature, the temperature dependences of hole concentration show activation behavior caused by the transition of holes from the impurity band to valence band. From these parts of the curves we estimated the Fermi energies in samples 2 and 1—15–20 and 6 meV, respectively. In sample 2 Fermi level is located in LH band,

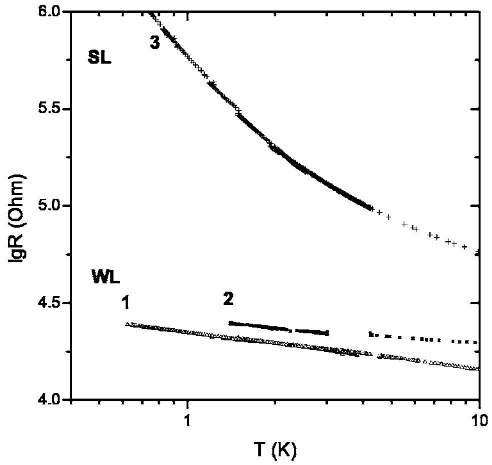

FIG. 1. Temperature dependences of the conductivity for three samples: samples 1,2 are in the weak localization regime, sample 3 is in the strong localization regime. Sample 1: wells (10 nm) and barriers are doped (the bulk Be concentration is 6×10^{17} cm^{-3}), sample 2: only wells (15 nm) are doped (bulk Be concentration is 10^{18} cm^{-3}), 3: wells (15 nm) and barriers doped (bulk Be concentration is 4×10^{17} cm^{-3}).

while in sample 1 it is in the UH band. At low temperatures, conductivity of these two samples depends on temperature very weakly (Fig. 1). As we have shown, it can be described by the weak localization theory.[1] In other words, these samples with large enough impurity concentrations exhibit properties of a "dirty metal." In contrast, sample 3 with lowest impurity concentration exhibited typical hopping conductivity which correspond to "strong localization" limit. It was despite of the fact that in this sample the upper Hubbard band was partly occupied and the corresponding localization length was estimated from quadratic positive magnetoresistance to be about 10 nm.[2]

A standard antilocalization behavior (positive magnetoresistance) was observed in 1,2 samples for magnetic fields >0.02 T, Fig. 2. For lower fields these two samples demonstrated also a small region of negative magnetoresistance (NMR), Fig. 3(a). In the contrast at the same temperatures and fields this NMR region was absent for sample 3, Fig. 3(b). Such a behavior was observed in 1,2 samples only for temperatures lower than ~3.4 K which corresponds to the

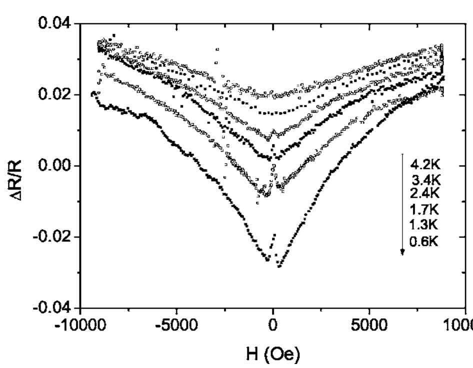

FIG. 2. High-field magnetoresistance at different temperatures for sample 1 (weak localization regime).

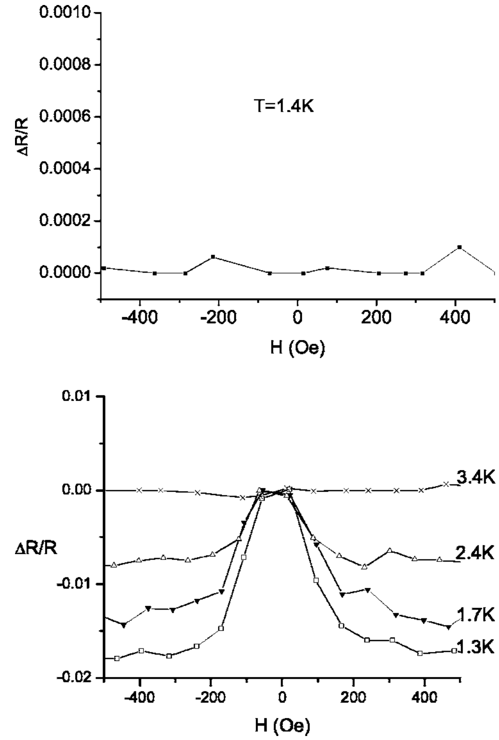

FIG. 3. Low-field magnetoresistance at different temperatures: (a) for sample 1 (weak localization regime), (b) for sample 3 (strong localization regime).

temperature of the superconducting transition in In. The NMR magnitude (1–10 %), depends on a concrete realization of the contact. Shown in Fig. 4 is the temperature dependence of this low field NMR. In the temperature region 3–1.2 K one can see the increase of this NMR magnitude with a temperature decrease following by saturation at temperatures 1.2–0.6 K.

III. THEORY

Let us consider a tunnel contact between a superconductor and a semiconductor sample. First we will discuss the situa-

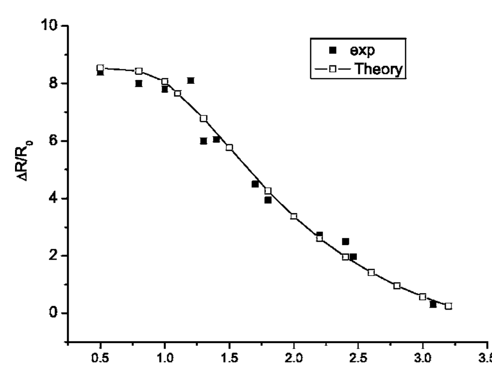

FIG. 4. Temperature dependence of negative magnetoresistance peak magnitude for sample 1 (experiment and theory). Theoretical fitting equation $\Delta R/R_0=(1/\{R_N/R_T+[\Delta(T)/\Delta(0)]^2R_N/R_A\}-1)R_N/R_0$. Where R_T, R_A, R_N are the single-particle tunneling, Andreev and normal state contact resistances, respectively. The fitting parameters are $R_N/R_0\approx4.2$, $R_N/R_A\approx0.33$.

tion on the metal side of the MIT, that is of the tunneling between a superconductor and a normal metal. As it is known, see, e.g., Ref. 3, the single-electron tunneling current can be written as

$$I = \frac{4\pi e}{\hbar}|T_0|^2 \int_0^\infty d\varepsilon\, \nu_{n1}(\varepsilon + eV)\nu_s(\varepsilon)[n_2(\varepsilon) - n_1(\varepsilon + eV)], \tag{1}$$

where T_0 is the tunneling matrix element and ν_{n1} is the density of states in the 2D metal. The density of quasiparticle states in a superconductor is

$$\nu_s = \nu_{n2}\frac{|\varepsilon|}{\sqrt{\varepsilon^2 - \Delta^2}} \quad \text{at } |\varepsilon| > \Delta;$$

$$\nu_s = 0 \quad \text{at } |\varepsilon| < \Delta, \tag{2}$$

where ν_{n2} is a density of state of a superconducting material in normal state. n_1 and n_2 are the quasiparticle occupation numbers in the 2D normal metal and superconductor, correspondingly. Thus one obtains

$$I_{N-S}(V) = \frac{4\pi}{\hbar}e|T_0|^2\nu_{n1}\nu_{n2}\int_{|\varepsilon|>\Delta}[n_1(\varepsilon) - n_2(\varepsilon + eV)]\frac{|\varepsilon|d\varepsilon}{\sqrt{\varepsilon^2 - \Delta^2}}. \tag{3}$$

At low temperature, $T \ll \Delta$, the integral (3) in the linear approximation in V can be estimated as

$$I_{N-S}(V) = VG\sqrt{\frac{2\pi\Delta}{T}}e^{-\Delta/T}, \tag{4}$$

where

$$G = \frac{4\pi}{\hbar}e^2|T_0|^2\nu_{n1}\nu_{n2}$$

is the contact conductance when the superconductor is in the normal state.

Now let us consider the behavior near the critical temperature $T \to T_c$. In the linear regime $eV < T$ the direct estimate of Eq. (3) yields

$$I_{N-S}(V) = VG\left[1 - \left(\frac{\Delta}{2T}\right)^2\right]. \tag{5}$$

One notes that according to Eq. (4) the current vanishes at $T \to 0$. In addition, there is also a contribution of the Andreev reflections which does not vanish (see, e.g., Ref. 3). Let us estimate the temperature at which the single-particle contribution crosses over to the Andreev contribution. The latter can be estimated as

$$G_A = \frac{e^2}{\hbar}A\Gamma^2, \tag{6}$$

where $\Gamma = \nu_{n1}n_{n2}|T_0|^2$ is the tunneling transparency and A is a constant. The quasiparticle contribution can be rewritten as

$$G_T = \frac{e^2}{\hbar}\sqrt{\frac{2\pi\Delta}{T}}e^{-\Delta/T}A\Gamma. \tag{7}$$

Correspondingly, these contributions are equal at some crossover temperature T^*. Considering T^* as given, one can estimate the tunneling transparency Γ in a simple way:

$$\Gamma = \sqrt{\frac{2\pi\Delta}{T'}}e^{-\Delta/T'}. \tag{8}$$

Now let us consider the insulator side of the MIT transition when one has tunneling between the superconductor and semiconductor in hopping regime. In this case the single-particle tunneling between semiconductor and superconductor banks can be controlled either by the direct resonant electron tunneling or by the phonon-assisted tunneling. The latter process dominates if the effective hopping energy band ε_0 is less than the superconducting gap. In this case the temperature behavior of the conductance is given as

$$G \propto e^{-(\Delta-\epsilon_0)/T}. \tag{9}$$

To the contrary, if $\varepsilon_0 > \Delta$ the contribution of resonant tunneling is expected to dominate.

The character of transport is also expected to be sensitive to the strength of the tunneling barrier. Indeed, transport in the semiconductor is controlled by the percolation cluster which allows self-averaging of the conductivity. Thus the contact resistance is the resistance between the superconductor and the percolation cluster. It consists of a sum of the resistance related to the last hop between some localized state in the semiconductor and the superconductor, and the resistance of the branch connecting this localized state to the percolation cluster. If the tunnel barrier transparency Γ is much less that the critical hopping exponent $\exp(-2r_h/a)$, where r_h is the typical hopping length, than the contact resistance is dominated by the "last" hop. In its turn, the contact conductance is a sum of the conductances corresponding to these hops. This fact allows us to average over these conductances (see Ref. 5). For each of the localized state i the corresponding conductance is

$$G_{T,i} \propto \frac{e^2}{\hbar}\nu_{n2}|T_0|^2\frac{1}{aT}\frac{\varepsilon_i}{\sqrt{\varepsilon_i^2 - \Delta^2}}\exp\left(-\frac{2x_i}{a}\right), \tag{10}$$

where x is the coordinate normal to the contact. Correspondingly, for the average one has

$$G_{SS} \propto \frac{e^2}{\hbar}|T_0|^2\frac{\nu_{n2}}{aT}\int_\Delta^{\epsilon_0} d\varepsilon\,\nu_{n1}(\varepsilon)\frac{\varepsilon}{\sqrt{\varepsilon^2 - \Delta^2}}\int d^2r dx \exp\left(-\frac{2x}{a}\right), \tag{11}$$

where ν_{n1} now plays the role of the density of localized states in semiconductor. Evaluating the expression (11) one obtains

$$G_{SS} \propto \frac{e^2}{\hbar}\Gamma\frac{\varepsilon_0}{T}\sqrt{1 - \left(\frac{\Delta}{\varepsilon_0}\right)^2}. \tag{12}$$

Now let us consider the situation when $\Gamma > \exp(-2r_h/a)$. In this case the contact between the superconductor and the percolation cluster is supported by some branches of hopping resistors. Since these branches are in parallel, the conductance of this structure is given as

$$G_{SS} = \sum_i \frac{1}{R_{T,i} + R_{c,i}}, \tag{13}$$

where $R_{T,i} = G_{T,i}^{-1}$ is the resistance of the hop form the last localized state to superconductor, while $R_{c,i}$ is the resistance of the branch connecting this last hop with the percolation cluster. One can separate the contact contribution to resistance as

$$\left(\sum_i \frac{1}{(R_{T,i} + R_{c,i})} \right)^{-1} - \left(\sum_i \frac{1}{R_{c,i}} \right)^{-1} = \frac{\sum_i \frac{R_{T,i}}{R_{c,i}(R_{T,i} + R_{c,i})}}{\sum_i \frac{1}{R_{c,i}} \sum_i \frac{1}{R_{c,i} + R_{T,i}}}. \tag{14}$$

Since the branches support the current flow through the system, according to ideas of the percolation theory they should also be considered as a part of the percolation cluster. Thus one expects that all of $R_{c,i}$ are of the order of that corresponding to the percolation threshold R_c. So the only average should be taken with respect to the localized state i corresponding to the last hop, actually—with respect to x_i. One notes that the upper limit for x_i is given by some critical value x_i corresponding to $R_{T,i} = R_c$. The larger resistances do not enter the percolation cluster. Then, one notes that for $x_i < x_c - a$ one has $R_{T,i} \ll R_c$ and the corresponding paths do not contribute effectively to the contact resistance. Thus only a small part of the branches given by a ratio $a/x_c \sim 1/\xi$ is important for the contact magnetoresistance.

Combining this estimate with the one given by Eq. (12) one notes that for the hopping conductivity the contact magnetoresistance is suppressed with respect to the metal conductor due to the two issues: (i) since the hopping energy band width ε_0 entering Eq. (12) is larger than T entering Eq. (5) and additional factor $(T/\varepsilon_0)^2 = \xi^{-2}$ appears at $\Delta < \varepsilon_0$; (ii) in contrast to a metal, in the hopping case only small part of tunneling events contributes to the contact magnetoresistance which gives a factor $1/\xi$. As a result, the contact magnetoresistance in hopping regime is suppressed with respect to metal by a factor ξ^{-3}.

Actually the measurable quantity is a sum of the the contact resistance and the sample resistance. So an increase of the sample resistance in the hopping regime $\propto \exp \xi$ means a decrease of *relative* magnetoresistance if the latter originates from the contact contribution. As a result, a suppression of the *relative* magnetoresistance in the hopping regime is given by the factor

$$\xi^{-3} \exp - \xi \tag{15}$$

rather than by a factor $\exp(-\xi)$ as one may conclude from naive considerations.

Note that as concerns the hopping regime we restricted ourselves by single particle tunneling. According to Eq. (9), this channel is exponentially frozen out at $T \to 0$. This conclusion still holds if the tunnel barrier is weak or even absent since it is based on the single-particle character of transport impossible in a superconductor at $T = 0$. As well known for the interface between a superconductor and a normal metal it is the two-particle Andreev reflections that are responsible for the low temperature transport. However the typical processes leading to hopping transport are single-particle ones. Note that the activation exponential factor $\exp[-\Delta/T]$ at small temperatures can be much smaller than variable range hopping factor $\exp(-\xi)$. This is the case, e.g., for studies[4] of hopping transport in CdTe structures with In contacts where the temperatures were as low as ~ 30 mK. Thus a process similar to the Andreev reflection, i.e., involving tunneling of Cooper pairs, should be considered. To the best of our knowledge, no detailed studies of such transport were reported until now. Indeed, the processes of Andreev reflections in *S-I-N* structures (where *I* is Anderson dielectric) were extensively studied by Frydman and Ovadyahu.[7] In particular, the role of coherent electron transport through highly conducting channels formed by the localized states according to Lifshitz-Kirpichenkov scenario[8] (see also Ref. 9) was considered. However the incoherent inelastic hopping transport and its matching to supercurrent was not discussed in Ref. 7. We are going to address this topic in a special paper.

IV. DISCUSSION

As clearly seen, the NMR at low fields can be only related to presence of a superconductor since the effect is absent at $T > T_c$. In principle, the magnetoresistance could be also related to interference contribution to Andreev tunneling, see, e.g., Ref. 6. However, our samples being close to MIT correspond to a limit of extreme dirty metal and thus characteristic magnetic field scales for the interference effects are much larger than critical magnetic field for In (see, e.g., our studies[1]). At the same time, the observed peak occurs at $H \lesssim 0.1$ T, that is of the order of the critical field for In but much less than required for the interference effects.

In Fig. 4 we have plotted the magnitude of the magnetoresistance peak as a function of temperature for the sample on the metal side of MIT. As it is seen, it is proportional to $\propto (T_c - T)$ in the vicinity of T_c. That agrees with Eq. (5), $\Delta R \propto \Delta^2$ since according to the BCS theory $\Delta \propto (T_c - T)^{1/2}$.

For lower temperatures Eq. (4) predicts exponential increase of magnetoresistance which can be seen from Fig. 4. With a further decrease of temperature the magnetoresistance saturates due to a bypassing of the Andreev channel. According to Eq. (8), the crossover temperature allows us to estimate the tunnel barrier transparency. Making use of the experimental data one concludes that $T^* \approx 1.1$ K. It is well known that the superconducting transition temperature for In is $T_c = 3.4$ K and thus $\Delta \approx 6$ K. Substituting Δ and T^* into Eq. (8) we estimate the tunneling probability as $\Gamma \approx 0.05$.

On the Fig. 4 we have also plotted theoretical curve resulting from the summation of the single-particle contribution calculated above and the contribution of the Andreev tunneling. Unfortunately the proper analysis of this latter channel at temperatures close to T_c is difficult due to the fact that our “normal metal” is actually a dirty semiconductor and is characterized by rather small spin-orbital scattering times ($> 10^{-11}$ s) and phase relaxation times ($\sim 10^{-11}$ s).[1] To the best of our knowledge, the detailed theory of Andreev reflections in the vicinity of T_c with an account of the factors mentioned above is absent. We have concluded that the best fitting of the experimental data is obtained if we approximate the Andreev channel contribution at $T \sim T_c$ as $\propto [\Delta(T)]^2/[\Delta(0)]^2$.

Now let us turn to the magnetoresistance in the hopping regime. As follows from Fig. 3(b), in the hopping regime there is no trace of the magnetoresistance peak at $T=1.4$ K with an accuracy at least of the order of 0.01%. Note that the corresponding magnitude of the relative magnetoresistance peak in the metal sample was of the order of $\sim 1\%$, Fig. 3(a). Such a suppression of the relative magnetoresistance in the hopping regime can hardly be explained by a simple increase of resistance. Indeed, the latter is larger than that of the metal sample only by a factor <100. However, this behavior is easily explained by our Eq. (15) predicting much stronger suppression of the magnetoresistance than following from the resistance ratio. In our opinion, this effect can be used to discriminate between the regimes of weak and localization even in the crossover region between the two regimes. The reason is in the principal differences in the physical picture of transport between the regimes which are not clearly seen in the value of the resistance itself. This is why we believe that our model is adequate for the experimental findings.

To conclude, we studied magnetoresistance of 2D p-GaAs/AlGaAs structures with superconducting electrodes (In) close to the metal-insulator transition. We demonstrated that the observed weak field magnetoresistance peak observed for metallic samples is due to restoration of the single-particle tunneling through the superconductor-semiconductor boundary with suppression of superconductivity. The crossover between the two tunneling regimes allowed us to estimate the tunneling transparency. The samples on the dielectric side of MIT did not demonstrate such a peak. the suppression of the magnetoresistance peak in the hopping regime is confirmed by theoretical analysis. We suggest to use this effect for discrimination between the regimes of weak and strong localization.

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