

Conductivity of a two-dimensional electron gas in a Si/SiGe heterostructure near the metal-insulator transition: Role of the short- and long-range scattering potential

E. B. Olshanetsky,¹ V. Renard,^{2,3} Z. D. Kvon,¹ J. C. Portal,^{2,3,4} N. J. Woods,⁵ J. Zhang,⁶ and J. J. Harris⁶

¹*Institute of Semiconductor Physics, Novosibirsk 630090, Russia*

²*GHMF, MPI-FKF/CNRS, Boîte Postal-166, F-38042, Grenoble, Cedex 9, France*

³*INSA-Toulouse, 31077, Toulouse, Cedex 4, France*

⁴*Institut Universitaire de France, Toulouse, France*

⁵*Imperial College of Science, London WC7 2BW, United Kingdom*

⁶*University College, London WC 15 75, United Kingdom*

(Received 14 March 2003; published 7 August 2003)

We report the observation of a metal-insulator transition (MIT) in a two-dimensional electron gas (2DEG) in a Si/SiGe heterostructure at zero magnetic field. On going through the MIT we observe the corresponding evolution of the magnetic field induced transition between the insulating phase and the quantum Hall (QH) liquid state in the QH regime. Similar to the previous reports for a GaAs sample, we find that the critical magnetic field needed to produce the transition becomes zero at the critical electron density corresponding to the zero field MIT. The temperature dependence of the conductivity in a metalliclike state at zero field is compared with the theory of the interaction corrections at intermediate and ballistic regimes $k_B T \tau / \hbar \geq 1$. The theory yields a good fit for the linear part of the curve. However, the slope of that part of $\sigma_{xx}(T)$ is about two times smaller than that reported in other 2D systems with similar values of r_s . At the same time, the recent theory of magnetoresistance due to electron-electron interaction in the case of arbitrary $k_B T \tau / \hbar$, smooth disorder and classically strong fields does not seem to be quite adequate for the description of the parabolic magnetoresistance observed in our samples. We attribute these results to the fact that neither of these theories deals with the whole scattering potential in a sample but leaves either its long-range or its short-range component out of consideration.

DOI: 10.1103/PhysRevB.68.085304

PACS number(s): 73.20.Fz, 71.30.+h, 72.15.Rn, 73.43.Nq

I. INTRODUCTION

According to the well established one-parameter scaling theory of conductivity¹ any two dimensional electron system can exist only in insulating state and no metal-insulator transition (MIT) should be possible. Thus, it came quite as a surprise when an apparent metallic state and an MIT were observed first in Si metal-oxide semiconductor field-effect transistor (MOSFETs) (Ref. 2) and, later, in several other two-dimensional (2D) systems, such as *p*-Si/SiGe, *n*- and *p*-GaAs/AlGaAs etc.³ The fact that the MIT is generally observed in systems with a large parameter r_s (the ratio of the interaction energy to the kinetic energy) is an indication that electron-electron interaction plays an important role in the MIT. Although the physics behind the MIT observed in strongly interacting systems remains elusive, significant progress has been recently made in the understanding of the nature of the metallic conductivity deeper in the metallic state. Zala *et al.*⁴ have calculated temperature-dependent corrections to conductivity due to electron-electron interactions based on the Fermi liquid approach. For sufficiently strong interactions Ref. 4 predicts a metallic temperature dependence in the entire temperature range. At very low temperatures, in the diffusive regime $k_B T \tau / \hbar \ll 1$ this is Finkelstein's weakly metallic logarithmic conductivity.⁵ At intermediate temperatures, in the ballistic regime $k_B T \tau / \hbar \geq 1$, the predicted temperature dependence is linear.⁶ The results^{4,7,8} have been reported to give a good description of the conductance temperature dependence and magnetoresistance in par-

allel field in several types of 2D systems on the metallic side of the apparent MIT.⁹⁻¹³ It should, however, be noted, that the theory^{4,7,8} relies on an approximation where the scattering potential is treated as pointlike scatterers. While appropriate for Si MOSFET's, this approximation is definitely not true for high mobility heterostructures where the impurities are separated from the 2DEG by an undoped spacer and the scattering potential is predominantly long range. Therefore, a certain caution should be used when applying the results of Refs. 4,7,8 for the description of the experimental data in high mobility heterostructures.

Until recently, in the context of the new approach⁴ no explicit discussion has been made of the interaction contribution to the magnetoresistance in a transverse magnetic field. The point has been recently addressed in Ref. 14 where the interaction contribution to the magnetoresistance in a classically strong transverse field is calculated for a smooth disorder and for arbitrary $k_B T \tau / \hbar$. Soon afterwards the results of Ref. 14 have been tested in an experiment on magnetoresistance in a GaAs/AlGaAs heterostructure with a 2DEG and a good agreement with theory¹⁴ has been found.¹⁵

In the present work we attempt an experimental investigation of these questions in a Si/SiGe heterostructure with a 2DEG. Despite the abundance of information on the MIT in Si MOSFET's and different other types of heterostructures, so far there has been no observation of an MIT in a *n*-Si/SiGe. This may in part be explained by the difficulty in fabricating a gated *n*-Si/SiGe structure with a stable and controllable behavior. At the same time an experimental study of these questions in a *n*-Si/SiGe is expected to be of some

interest. Such a study would provide an opportunity to compare the electron transport properties of a Si MOSFET and a Si/SiGe heterostructure in the vicinity of the MIT and in the “metallic” regime. This comparison may be instructive as these two silicon structures have very similar electron energy spectrum, but differ in the structure of the disorder potential, which is predominantly short range in Si MOSFET’s but has an important long range component in a Si/SiGe heterostructure. In this work we report the observation of the MIT in a Si/SiGe heterostructure with a 2D electron gas and perform an analysis of the magnetoresistance data and the conductance temperature dependence on the metallic side of the MIT using the results of Refs. 4 and 14. The theory⁴ produces a good fit for the linear part of the conductance temperature dependence. However, the value of the interaction constant F_0^σ serving as the only fitting parameter turns out to be about two times smaller than that reported for Si MOSFET’s with similar values of r_s .⁹ Also, in contrast to Ref. 15 we find a rather big divergence between our magnetoresistance data and the behavior predicted in Ref. 14. We attribute all these results to the possibility that in the intermediate and ballistic regimes the case of a mixed mechanism of scattering, i.e., the situation where both the short- and the long-range scattering are equally important for electron transport, cannot be properly described by a model that ignores either one of the other component of disorder.

II. EXPERIMENTAL PROCEDURES

The n -type modulation doped Si/Si_{1-x}Ge_x samples used in this study were grown by MBE.¹⁶ On a high resistivity p -type Si substrate a 1.4 μm graded Si/Si_{1-x}Ge_x layer, in which the Ge content increases linearly from 5 to 35 %, was grown, followed by 0.5 μm of constant composition Si_{0.7}Ge_{0.3} alloy. These layers produced a strain-relaxed but low-dislocation-density surface onto which the tensilely strained Si channel, 110 \AA thick, was deposited, followed by an undoped Si/Si_{0.7}Ge_{0.3} spacer layer, a doped Si/Si_{0.7}Ge_{0.3} supply layer, and finally undoped Si/Si_{0.7}Ge_{0.3} and Si cap layers. The thickness and As doping levels in the supply layer were varied slightly between the samples. To measure the magnetoresistivity of the samples, the layers were photolithographically processed into 50 μm wide Hall bars with the distance between the four voltage probes on each side 100, 250, and 100 μm . The mesa pattern for the Hall bars was etched by reactive ion etching to define the channel geometry. The Ohmic contacts to the Hall bars were prepared by ion doping of phosphorus.

The magnetotransport properties of the structures were measured at 0.4–4.2 K in magnetic fields up to 15 T using a superconducting solenoid and He₃ and VTI cryostats. The data was acquired using standard a.c. lock-in techniques with a frequency around 10 Hz. Driving currents through the Hall bar of 0.1 μA were chosen to avoid the heating effects. All our samples behaved in a more or less similar way. In what follows we discuss the results obtained on a sample with the electron density $n_s = (3.53 - 6.23) \times 10^{11} \text{ cm}^{-2}$ depending on the prior illumination and the highest mobility $\mu \approx 6 \times 10^4 \text{ cm}^2/\text{Vs}$.

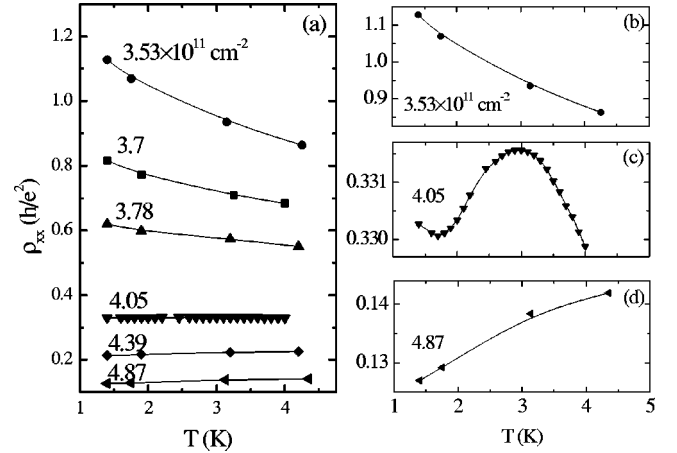


FIG. 1. (a) Resistivity versus temperature for different electron densities (b)–(d). Some of the traces from (a) given in more detail.

III. RESULTS AND ANALYSIS

Despite a large number of articles on the MIT in various types of 2D semiconductor systems, so far there has been no report of a MIT in a Si/SiGe heterostructure with a 2DEG. The MIT is typically observed by measuring the temperature dependence of a sample resistance for different carrier densities which are usually controlled by voltage applied to the electrostatic top gate. Reportedly, it has been found difficult to fabricate a n -Si/SiGe heterostructure with a gate that does not leak. The leakage from the gate to the 2DEG results in a sample instability that makes it unsuitable for experiment. In the present work, after several unsuccessful attempts to fabricate an operational gated n -Si/SiGe structure, we decided to perform our study using structures without gate. Among the several samples that we have tried there were some that, upon cooling down to low temperatures in the dark were found in high resistance low electron density states. The samples were also found to be very sensitive to illumination. Exposing them to a specially selected dose of red LED radiation permitted us to fine-tune the sample resistance with a good precision. The sample states of different resistivity thus obtained were characterized by a high stability yielding reproducible dependences of resistance versus temperature and magnetic field over a considerable, up to several hours, period of time. Shown in Fig. 1 is a series of resistance versus temperature traces for different electron densities obtained using the procedure described above. Similar to the previously reported MITs in other 2D semiconductor systems³ we observe a transition from the insulator to metalliclike behavior as the electron density is varied from 3.53×10^{11} to $4.87 \times 10^{11} \text{ cm}^{-2}$. To our knowledge this is the first observation of this kind in a Si/SiGe heterostructure with a 2DEG. In Figs. 1(b), 1(c), 1(d) some of the curves are given in more detail. With our method of varying the electron density and a rather limited temperature range used it is difficult to determine with a high precision the “separatrix,” i.e., the curve corresponding to the transition from the metallic to insulator regime. Nevertheless we consider the curve shown in Fig. 1(c) to be close to that transition point. There are two things about this curve that are worth mentioning. Firstly, the resis-

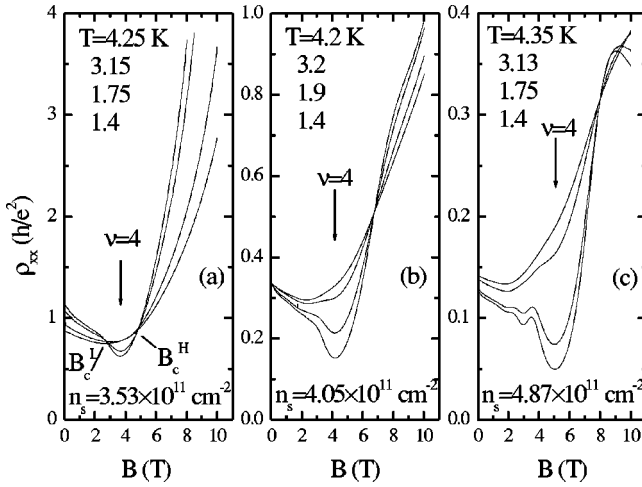


FIG. 2. The magnetoresistance traces taken at different temperatures corresponding to some of the sample states shown in Fig. 1.

tivity of the sample corresponding to the separatrix is about $0.3h/e^2$ which is about ten times lower than that reported for Si MOSFET's.³ Secondly, the peculiar nonmonotonic temperature dependence corresponding to our separatrix is also quite unusual. At present we have no explanation for these observations.

Before the first observation of the MIT in Si MOSFET (Ref. 2) it has been widely accepted that the electrons confined to two dimensions have an insulating ground state at zero magnetic field and zero temperature. On the other hand, the application of a strong perpendicular magnetic field changes the situation, resulting in a transition from the insulating state to a QH liquid state. This magnetic-field-driven transition from the insulator to the QH liquid state has attracted much attention¹⁷ and some interesting hypothesis have been proposed for its explanation.^{18,19} Later, it was suggested²⁰ that there possibly exists a link between the MIT transition observed at zero magnetic field and the finite magnetic field transitions studied previously. In the present work we have made a similar observation in our samples. Figure 2 shows the magnetoresistance traces taken at different temperatures corresponding to some of the zero field sample states presented in Fig. 1. For any given magnetic field it is the sign of the resistance temperature dependence that is the indicator of whether the state of the sample at this particular field is insulating or metallic. Let us start with a sample in an insulating state at zero field. In Fig. 2(a) we see that for $n_s = 3.53 \times 10^{11} \text{ cm}^{-2}$ the zero field insulating state persists up to $B_c^L \approx 2.8 \text{ T}$ where the first insulator-to-QH liquid transition occurs. This QH liquid state corresponding to filling factor $\nu = 4$ is then maintained until the QH liquid-to-insulator transition turns the sample again into insulator at $B_c^H \approx 4.8 \text{ T}$. According to Fig. 1, increasing the electron density from $n_s = 3.53 \times 10^{11} \text{ cm}^{-2}$ brings the sample closer to the zero field MIT. It is interesting to note, that in magnetic field this results in a corresponding gradual decrease of B_c^L until, finally, at the carrier density $n_s = 4.05 \times 10^{11} \text{ cm}^{-2}$ appropriate for the zero field MIT, B_c^L becomes exactly zero, Fig. 2(b). On further increasing the carrier density, the zero field me-

talic state is maintained at finite field and the first magnetic-field-driven transition to occur is thus a QH liquid-to-insulator transition at higher fields, Fig. 2(c). As the authors of Ref. 20, we find the behavior described above to be an indication that both the finite- and zero-field transitions are likely to share a common physical origin.

Up to date there is still no theoretical description available for the behavior of a 2D system in the vicinity of the MIT, i.e., when $\rho_{xx} \approx h/e^2$, and all the discussions of the MIT are, therefore, necessarily qualitative in character. On the contrary there is now a number of comprehensive theoretical models for a situation when $\rho_{xx} \ll h/e^2$. These models offer predictions concerning the corrections to the zero temperature Drude conductivity $\sigma_D = e^2 n \tau / m^*$ under various physical conditions. We now turn to the discussion of these quantum corrections in the limit $\rho_{xx} \ll h/e^2$, i.e., when our samples show metallic behavior.

On the whole, there are two distinct types of corrections to conductivity. First, there is the weak localization correction calculated in the approximation of noninteracting electrons and describing the variation of conductivity due to the electron coherent scattering by the random potential²¹

$$\Delta \sigma_{xx}^{wl} = \alpha p (e^2 / \pi \hbar) \ln(k_B T \tau / \hbar), \quad (1)$$

where it is assumed that the phase breaking time τ_ϕ varies as T^{-p} . The amplitude α is expected to be 1 for normal scattering (-0.5 for pure spin-orbit scattering and 0 for scattering on magnetic impurities). In Si samples only the Coulomb scattering is present. The situation is more complicated when the interaction between electrons is taken into account. The first theory to consider the interaction corrections to conductivity was developed for the diffusive regime $k_B T \tau / \hbar < 1$, i.e., for a situation where the electron experiences multiple scattering by impurities during its life time.⁵ Of the two components of the interaction correction in this regime, the correction in the Cooper channel is negligible when $T_c \ll T$, where T_c is the critical temperature for superconductor transition, and the remaining correction in the Diffusion channel has the form

$$\Delta \sigma_{xx}^{ee} = \frac{e^2}{2\pi^2 \hbar} \ln \left[\frac{k_B T \tau}{\hbar} \right] \left\{ 1 + 3 \left[1 - \frac{\ln(1 + F_0^\sigma)}{F_0^\sigma} \right] \right\}, \quad (2)$$

where F_0^σ is the interaction constant in the triplet channel. This constant is negative if the electron-electron interaction tends to align the electron spins and its absolute value increases with the interaction strength characterized by parameter r_s . Note that the sign of this logarithmic in temperature correction depends on the absolute value of F_0^σ . It is believed that in the limit of very strong interaction the constant F_0^σ approaches -1 leading to the Stoner instability with all electron spins parallel in the absence of magnetic field.

Recently, a theory has been proposed⁴ that along with the diffusive regime covers the intermediate and ballistic regimes $k_B T \tau / \hbar \geq 1$. It is shown that the interaction corrections in all these regimes are a consequence of the same physical process, namely the coherent electron scattering on the modulated density of other electrons (Friedel oscillations).

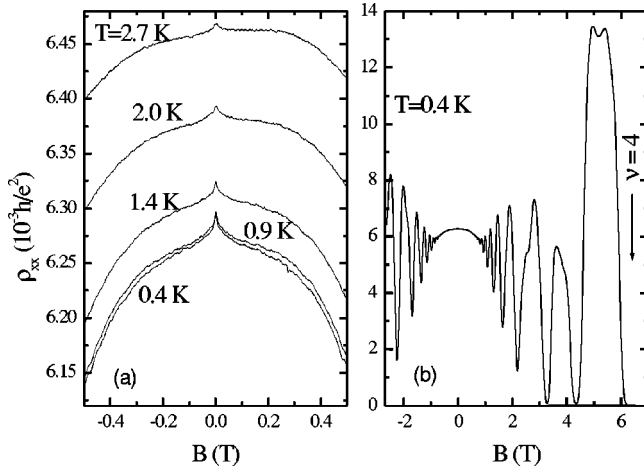


FIG. 3. (a) Magnetoresistance in a metalliclike state of the sample at different temperatures. (b) Magnetoresistance shown in a wider field range at $T=0.4$ K.

tions) caused by an impurity with a short range potential. The quantum interaction correction in this theory can be expressed as

$$\Delta\sigma_{xx}^{ee} = \delta\sigma_C + 15\delta\sigma_T, \quad (3)$$

where

$$\delta\sigma_C = \frac{e^2}{\pi\hbar} \frac{k_B T \tau}{\hbar} \left[1 - \frac{3}{8} f(k_B T \tau / \hbar) \right] - \frac{e^2}{2\pi^2 \hbar} \ln \left[\frac{1}{k_B T \tau / \hbar} \right],$$

is the charge channel correction and

$$\delta\sigma_T = \frac{F_0^\sigma}{[1 + F_0^\sigma]} \frac{e^2}{\pi\hbar} \frac{k_B T \tau}{\hbar} \left[1 - \frac{3}{8} t(k_B T \tau / \hbar; F_0^\sigma) \right] - \left[1 - \frac{1}{F_0^\sigma} \ln(1 + F_0^\sigma) \right] \frac{e^2}{2\pi^2 \hbar} \ln \left[\frac{1}{k_B T \tau / \hbar} \right]$$

is the correction in the triplet channel. We also take into account the two fold valley degeneracy at (100) Si surface that changes the triplet term prefactor from 3 to 15. The explicit expressions for the functions $f(k_B T \tau / \hbar)$ and $t(k_B T \tau / \hbar; F_0^\sigma)$ are given in Ref. 4. In the diffusive limit Eqs. (3) reproduce the well known logarithmic temperature correction (2). In the intermediate and ballistic regimes Eqs. (3) result in a linear with temperature correction to conductivity with the slope and the sign determined by the absolute value of the interaction constant F_0^σ . For small F_0^σ the triplet channel contribution is small and the temperature dependence, governed mostly by the correction in the charge channel, is insulating, whereas for higher F_0^σ values the negative triplet term increases and the behavior becomes metallic. Also, since the triplet term prefactor increases with the valley degeneracy, systems with a higher valley degeneracy show metallic behavior for lower values of F_0^σ and hence for lower interaction strength than systems with lower or no valley degeneracy.

Figure 3 shows the typical magnetoresistance traces ob-

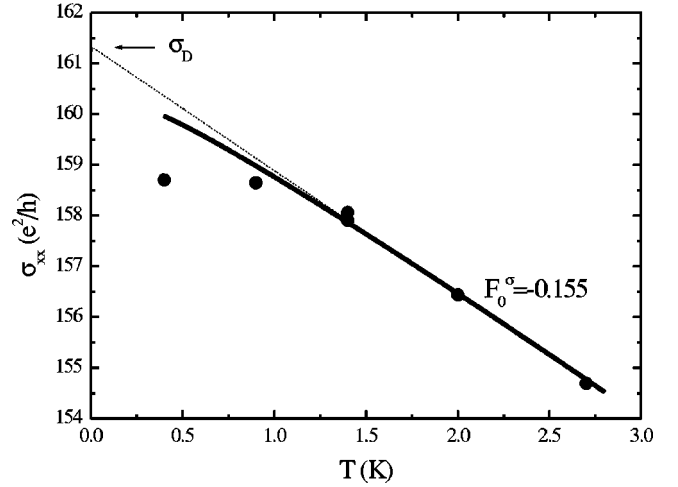


FIG. 4. Conductivity versus temperature at zero field obtained from the magnetoresistance traces in Fig. 3(a). The thick solid line is the sum of the localization term and the interaction correction (Ref. 4) with $F_0^\sigma = -0.155$.

tained at different temperatures in our samples after the electron concentration had been saturated with a maximum LED illumination. The corresponding electron mobility is $\mu = 61\,800$ cm²/Vs and the electron density $n_s = 6.23 \times 10^{11}$ cm⁻². In Fig. 4 we plot the zero field conductivity values from Fig. 3(a) versus temperature. In line with the earlier report on a high mobility *n*-Si/SiGe structure,²² one can see in Fig. 4 that the sample behavior is markedly metallic. The temperature dependence appears to be linear for $T \geq 1.25$ K and saturates at lower temperatures. By extrapolating the linear part of the curve to $T=0$ we find the Drude conductivity and the corresponding momentum relaxation time $\tau = 6.8 \times 10^{-12}$ s. Hence, we get $k_B T \tau / \hbar = 0.89 \times T$ which means that in the experimental temperature range 0.4–2.7 K the sample operates in the intermediate and ballistic regimes. The solid line in Fig. 4 is a theoretical fit given by the sum of the weak localization term [Eq. (1)] and the interaction correction [Eqs. (3)] with the only fitting parameter being the interaction constant $F_0^\sigma = -0.155$. The weak localization term prefactor $\alpha p = 1.5$ used in the fit is obtained from the analysis of the magnetoresistance data in low fields.

It is impossible to compare the calculated value of F_0^σ with the experimental parameter r_s which is ~ 6.7 for our sample since the functional relation between F_0^σ and r_s is unknown for $r_s > 1$. It is, however, interesting to note that systematically higher values of F_0^σ have been reported for various other types of 2D systems and, in particular, for Si MOSFET's with r_s values close to ours. For example, $F_0^\sigma \approx -0.27$ has been found in an *n*-channel Si MOSFET (Ref. 9) and in a *p*-type GaAs/AlGaAs heterostructure¹⁰ with $r_s \approx 7$. The difference in the F_0^σ values obtained in our Si/SiGe heterostructure and in an *n*-channel Si MOSFET (Ref. 9) is, in our opinion, the consequence of a difference in the structure of the scattering potential in these two silicon systems. The linear-in- T term of Ref. 4 arises only due to the scattering on a short range impurity potential (in a purely smooth disorder this term is exponentially suppressed). In a Si MOS-

FET with a predominantly short-range scattering potential the results of Ref. 4 are, thus, quite appropriate. In our samples the single particle scattering time τ_q obtained from the magnitude of the Shubnikov–de Haas oscillations is about 6 times smaller than the momentum relaxation time τ indicating that apparently both types of disorder are present. Therefore the prefactor in front of the linear term would be reduced in our case as compared to the corresponding expression in Ref. 4 and, hence, a lower value of F_0^σ would be obtained from fitting the theory⁴ to experiment which, under these conditions, is unlikely to reflect the actual interaction intensity in the sample. This effect is expected to be even more pronounced in such high mobility structures as *n*- and *p*-GaAs/AlGaAs and, in our opinion, a certain caution should be used when comparing experiment with theory⁴ in these cases.

On the whole the theory⁴ produces a good fit for the experimental curve at temperatures $T \geq 1.25$ K. On the lower temperature side the dependence $\sigma_{xx}(T)$ tends to saturate, in contrast to the theoretical prediction. Similar departures from theory have been observed in all our samples and also by other authors in a Si 2D system different from ours.¹¹ A nonzero valley splitting and a strong sample specific intervalley scattering have been suggested as a possible explanation of the weakening of the $\sigma_{xx}(T)$ dependence. At present there is no theory available that takes intervalley scattering into account.

Let us now analyze the transverse magnetoresistance in our samples, Fig. 3(a). In the absence of any quantum corrections to the conductivity one has $\sigma_{xx}^0(B) = \sigma_D / [1 + (\mu B)^2]$, $\sigma_{xy}^0(B) = -\mu B \sigma_{xx}^0(B)$, $\rho_{xy}^0(B) = -\mu B / \sigma_D$, and $\rho_{xx}^0(B) = 1 / \sigma_D = \text{const}$, i.e., the resistance of the sample is independent of magnetic field. If, on the other hand, there are some small corrections to the conductivity tensor components $\sigma_{xx}(B) = \sigma_{xx}^0(B) + \Delta\sigma_{xx}$ and $\sigma_{xy}(B) = \sigma_{xy}^0(B) + \Delta\sigma_{xy}$, then, converting the conductivity tensor into the resistivity tensor, one gets

$$\rho_{xx}(B) = \rho_D + \rho_D^2 [(\mu B)^2 - 1] \Delta\sigma_{xx} + 2\mu B \Delta\sigma_{xy}. \quad (4)$$

In general there will be contributions to the magnetoresistance both from the weak localization and from the interaction. In the diffusive approximation the WL corrections to $\sigma_{xx}^0(B)$ and $\sigma_{xy}^0(B)$ are restricted to magnetic fields below $B_{tr} \approx \hbar / 2el_e^2$. For higher mobility structures B_{tr} is usually very small (~ 1 mT in our case). At the same time the WL magnetoresistance extends up to fields several times larger than B_{tr} [see Fig. 3(a)].²³

For the interaction contribution to the magnetoresistance it has been established that in the diffusive regime $k_B T \tau / \hbar \ll 1$ the longitudinal conductivity correction $\Delta\sigma_{xx}^{ee}$, given by Eq. (2), remains unchanged at low and classically strong magnetic fields, while, at the same time, the Hall conductivity is not affected by the interaction $\Delta\sigma_{xy}^{ee} \equiv 0$.⁵ As can be seen from Eq. (4) for $\omega_c \tau > 1$ this will lead to a parabolic negative magnetoresistance

$$\rho_{xx}(B) = \rho_D + \rho_D^2 (\mu B)^2 \Delta\sigma_{xx}^{ee}(T). \quad (5)$$

with $\Delta\sigma_{xx}^{ee}(T)$ given by Eq. (2). Finally, at still higher fields $B_Z \geq kT / g\mu_B$ a positive magnetoresistance develops caused by the Zeeman effect on the interaction correction $\Delta\sigma_{xx}^{ee}$.⁵

In contrast to the diffusive regime, the behavior of magnetoresistance is less well understood in the intermediate and ballistic regimes $k_B T \tau / \hbar \geq 1$. The theory of interaction corrections at arbitrary relation between temperature and elastic mean-free time⁴ leaves magnetoresistance in a perpendicular magnetic field out of consideration. It remains unclear whether the corrections to conductivity obtained in this theory for zero field would persist in finite magnetic field and give rise to a magnetoresistance in a manner similar to the diffusive regime (5) or whether they will be modified or suppressed by magnetic field. As has been mentioned above, the results of the theory⁴ are obtained in the approximation where the impurities are treated as pointlike scatterers. This is justified at zero magnetic field where a long-range potential reduces the backscattering and suppresses the interaction corrections in the ballistic regime. The situation, however, changes in a strong magnetic field which increases the probability of an electron to return back and, thus, restores the interaction correction even in the case of a purely smooth disorder. Recently, a theory¹⁴ has been proposed that considers the interaction in the ballistic and intermediate regime in classically strong fields $\omega_c \tau > 1$. The theory¹⁴ takes into account only the long-range component of the scattering potential and shows that in this case the interaction will also lead to a parabolic negative magnetoresistance described by Eq. (5) with a distinct expression for $\Delta\sigma_{xx}^{ee}(T)$. However, as will be shown below, while justified in the case of a predominantly long-range scattering potential, the exclusion of the short-range component of disorder from consideration may be inappropriate in a more general case of a mixed mechanism of scattering. Finally, it might be well to point out that unlike the diffusive regime the conductivity correction $\Delta\sigma_{xx}^{ee}(T)$ extracted from magnetoresistance in the ballistic regime¹⁴ has nothing to do with the correction at zero field.

Up to date there have been several experimental studies of the interaction-related transverse magnetoresistance in the intermediate magnetic field range where the weak localization is already suppressed while the Zeeman effect on the interaction correction has not yet set in. First to mention are the experiments^{24–27} where the observed temperature dependent negative magnetoresistance was compared to the theoretical prediction for the diffusive regime (5). In one of these works²⁷ the samples used were *n*-Si/SiGe heterostructures very similar to ours. However, with the exception of Ref. 25 these experiments have been performed on high mobility structures in the ballistic regime where the validity of Eq. (5) for the description of the magnetoresistance is questionable. Recently an experimental work has been reported¹⁵ whose aim was to study the magnetoresistance caused by electron-electron interactions in the intermediate regime in a *n*-GaAs/AlGaAs heterostructure with long-range fluctuation potential and to compare the results with the theoretical prediction.¹⁴ A good agreement between the theory and experiment has been reported.

Similar to Ref. 15 in our case we have a relatively high mobility 2D electron gas in the intermediate and ballistic

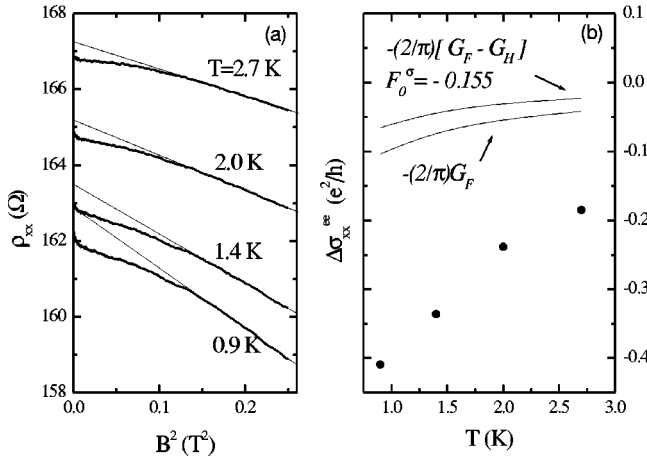


FIG. 5. (a) Magnetoresistance data from Fig. 3(a) plotted versus B^2 . (b) The black circles are $\Delta\sigma_{xx}^{ee}(T)$ obtained using Eq. (5) from the slope of the straight part of the magnetoresistance traces in (a); the solid curves are theory (Ref. 14) (see text).

regimes. The important role of the long-range scattering potential is evidenced by the fact that the transport time is significantly longer than the quantum single-particle time. The sharp change of the magnetoresistance observed in Fig. 3(a) in low magnetic field is caused by the WL effect suppressed at higher fields. Then follows a relatively flat region, which, according to Ref. 14, is supposed to correspond to the suppression of backscattering by the long-range potential in low fields. This effect is more pronounced for higher temperatures. At still higher field the backscattering is restored by the increasing magnetic field and a parabolic magnetoresistance sets in. In our sample the condition $\omega_c\tau=1$ is satisfied at $B=0.16$ T and so the parabolic magnetoresistance in Fig. 3(a) is observed in classically strong fields. On the other hand, except for $T=0.4$ K, the magnetoresistance traces in Fig. 3(a) are measured at fields well below $B_Z[T]=0.74 \times T[K]$ at which the Zeeman effect on the interaction correction should become important. That means that the parabolic magnetoresistance in Fig. 3(a) is observed in the magnetic field region where the results of Ref. 14 are expected to be applicable. In Fig. 5(a) we plot the resistivity from Fig. 3(a) as a function of B^2 . Shown in Fig. 5(b) are $\Delta\sigma_{xx}^{ee}(T)$ obtained using Eq. (5) from the slope of the straight part of the magnetoresistance traces in Fig. 5(a). The theoretical curves plotted in Fig. 5(b) are

$$\Delta\sigma_{xx}^{ee} = -\frac{2}{\pi} G_F(k_B T \tau / \hbar) \quad (6)$$

and

$$\Delta\sigma_{xx}^{ee} = -\frac{2}{\pi} [G_F(k_B T \tau / \hbar) - G_H(k_B T \tau / \hbar; F_0^\sigma)],$$

where $G_F(k_B T \tau / \hbar)$ is the exchange contribution to the interaction correction to conductivity in classically strong magnetic fields and in the case of a long range scattering poten-

tial. $G_H(k_B T \tau / \hbar; F_0^\sigma)$ is the triplet Hartree term. The explicit expressions for these functions are given in Ref. 14. For plotting in Fig. 5(b) we take these functions in the appropriate for our experimental situation limit $\kappa \gg k_F$, where k_F is the Fermi wave vector and $\kappa = 4\pi e^2 \nu$ is the inverse screening length. The Hartree term is plotted for $F_0^\sigma = -0.155$, the interaction constant obtained from the fit of the zero-field resistance temperature dependence to theory.⁴ As can be seen there is a rather big divergence between the experimental points and the theoretical curves in Fig. 5(b). The exchange correction independent of the interaction constant lies closer to the experiment. Taking into account the Hartree contribution with a sign opposite to the exchange term increases the discrepancy, which becomes even greater if it is multiplied by a factor 5 corresponding to the valley degeneracy [compare Eqs. (3)]. In our opinion the explanation why in contrast to Ref. 15 we do not find a good agreement with Ref. 14 in our experiment may lie with the fact that theory¹⁴ was developed with only the long-range component of the disorder potential taken into account. The structures studied in Ref. 15 were of a higher mobility than ours and with a more important role of the long-range scattering potential compared to that of the pointlike scatterers. Probably in that case the approximation made in Ref. 14 is adequate, whereas in our samples the short-range scattering can not be completely neglected and an extension of theory¹⁴ to the case of a mixed disorder is needed to explain our results.

IV. CONCLUSIONS

To conclude, we have observed a MIT in a Si/SiGe heterostructure with a 2DEG. We have also observed the evolution of the magnetic-field-driven transition from an insulator to a QH liquid state as our system is step by step driven through the MIT by gradually increasing the electron density. Similar to Ref. 20 we find that the magnetic field corresponding to the insulator-to-QH liquid state transition becomes zero at the critical electron density appropriate for the MIT in zero field. On the metallic side of the MIT we find that the theory⁴ of the interaction corrections to conductivity for an arbitrary relation between temperature and the momentum relaxation time provides a good description of the linear part of the conductivity temperature dependence at zero field but with the interaction constant F_0^σ about two times smaller than in Si MOSFET's with a similar value of r_s . The recently developed theory¹⁴ of the interaction contribution to magnetoresistance at classically strong magnetic fields and for arbitrary $k_B T \tau / \hbar$ does not seem to be appropriate for the description of the parabolic negative magnetoresistance observed in our system. The theory¹⁴ was built in the assumption that there is only a long-range scattering potential. The importance of a short-range disorder in our samples may be the reason for the observed discrepancy with the theory. Both these theories^{4,14} neglect either the long- or short-range component of disorder and for that reason do not seem to be quite adequate for the description of the case of a mixed disorder.

ACKNOWLEDGMENTS

We are grateful to I.V. Gornyi and A.D. Mirlin for helpful discussion. This work was supported by PICS-RFBR (Grant

No. 1577), RFBR (Grant No. 02-02-16516), NATO Linkage (Grant No. CLG.978991), INTAS (Grant No. 01-0014), programs “Physics and Technology of Nanostructures” of the Russian ministry of Industry and Science and “Low dimensional quantum structures” of RAS.

-
- ¹P.A. Lee and T.V. Ramakrishnan, *Rev. Mod. Phys.* **57**, 287 (1985).
- ²S.V. Kravchenko, G.V. Kravchenko, J.E. Furneaux, V.M. Pudalov, and M. D’Iorio, *Phys. Rev. B* **50**, 8039 (1994).
- ³E. Abrahams, S.V. Kravchenko, and M.P. Sarachik, *Rev. Mod. Phys.* **73**, 251 (2001).
- ⁴G. Zala, B.N. Narozhny, and I.P. Aleiner, *Phys. Rev. B* **64**, 214204 (2001).
- ⁵B. L. Altshuler and A. G. Aronov, *Electron-Electron Interactions in Disordered Systems*, edited by A. L. Efros and M. Pollak (Elsevier Science Publishers, British Vancouver, 1985); A.M. Finkelstein, *JETP Lett.* **37**, 517 (1983).
- ⁶A. Gold and V.T. Dolgoplov, *Phys. Rev. B* **33**, 1076 (1986).
- ⁷G. Zala, B.N. Narozhny, and I.L. Aleiner, *Phys. Rev. B* **65**, 020201 (2002).
- ⁸G. Zala, B.N. Narozhny, and I.L. Aleiner, *Phys. Rev. B* **64**, 201201 (2001).
- ⁹A.A. Shashkin, S.V. Kravchenko, V.T. Dolgoplov, and T.M. Klapwijk, *Phys. Rev. B* **66**, 073303 (2002).
- ¹⁰Y.Y. Proskuryakov, A.K. Savchenko, S.S. Safonov, M. Pepper, M.Y. Simmons, and D.A. Ritchie, *Phys. Rev. Lett.* **89**, 076406 (2002).
- ¹¹V.M. Pudalov, M.E. Gershenson, H. Kojima, G. Brunthaler, A. Printz, and G. Bauer, *cond-mat/0205449* (unpublished).
- ¹²G. Brunthaler, A. Printz, G. Pillwein, P.E. Lindelof, and J. Ahopelto, *cond-mat/0207170* (unpublished).
- ¹³S.A. Vitkalov, K. James, B.N. Narozhny, M.P. Sarachik, and T.M. Klapwijk, *cond-mat/0204566* (unpublished).
- ¹⁴I.V. Gornyi and A.D. Mirlin, *Phys. Rev. Lett.* **90**, 076801 (2003).
- ¹⁵L. Li, Y.Y. Proskuryakov, A.K. Savchenko, E.H. Linfield, and D.A. Ritchie, *Phys. Rev. Lett.* **90**, 076802 (2003).
- ¹⁶J.M. Fernandez, A. Matsumura, X.M. Zhang, M.H. Xie, L. Hart, J. Zhang, and B.A. Joyce, *J. Mater. Sci.: Mater. Electron.* **6**, 330 (1995).
- ¹⁷H.W. Jiang, C.E. Johnson, K.L. Wang, and S.T. Hannah, *Phys. Rev. Lett.* **71**, 1439 (1993).
- ¹⁸S. Kivelson, D.H. Lee, and S.C. Zhang, *Phys. Rev. B* **46**, 2223 (1992).
- ¹⁹D.E. Khmel’nitskii, *Phys. Lett.* **106A**, 182 (1984); *JETP Lett.* **38**, 552 (1983).
- ²⁰Y. Hanein, N. Nenadovic, D. Shahar, H. Shtrikman, J. Yoon, C.C. Li, and D.C. Tsui, *Nature (London)* **400**, 735 (1999).
- ²¹E. Abrahams, P.W. Anderson, D.C. Licciardello, and T.V. Ramakrishnan, *Phys. Rev. Lett.* **42**, 673 (1979).
- ²²T. Okamoto, K. Hosoya, S. Kawaji, A. Yagi, A. Yutani, and Y. Shiraki, *Physica E (Amsterdam)* **6**, 260 (2000).
- ²³H.P. Wittman and A. Schmid, *J. Low Temp. Phys.* **69**, 131 (1987).
- ²⁴M.A. Paalanen, D.C. Tsui, and J.C.M. Hwang, *Phys. Rev. Lett.* **51**, 2226 (1983).
- ²⁵G.M. Minkov, O.E. Rut, A.V. Germanenko, A.A. Sherstobitov, V.I. Shashkin, O.I. Khrykin, and V.M. Daniltsev, *Phys. Rev. B* **64**, 235327 (2001).
- ²⁶K.K. Choi, D.C. Tsui, and S.C. Palmateer, *Phys. Rev. B* **33**, 8216 (1986).
- ²⁷T.J. Thornton, A. Matsumura, and J. Fernandez, *Surf. Sci.* **361/362**, 547 (1996).