Electron-electron interaction in doped GaAs at high magnetic field

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We observe an inversion of the low-temperature dependence for the conductivity of doped GaAs by application of a magnetic field. This inversion happens when $\omega_c\tau_{\rm tr}\simeq 1$, as predicted by Houghton *et al.* [Phys. Rev. B **25**, 2196 (1982)] for the correction to conductivity due to screened Coulomb repulsion in the diffusive regime. This correction follows the oscillating behavior of the transport elastic time entering the Shubnikov–de Haas regime. For $\omega_c\tau\geqslant 1$, we observe that the Hartree part of the interaction correction is suppressed. Moreover, the total correction seems strongly reduced, although its dependence stays logarithmic. [S0163-1829(98)06604-1]

Electron-electron interaction (EEI) and weak-localization corrections determine the low-temperature dependence for the conductivity of disordered metals and highly doped semiconductors. In the two-dimensional case, following Altshuler Aronov, and Lee (AAL),[1] the EEI correction to the conductivity is given in zero magnetic field and in absence of any spin relaxation by

$$\delta\sigma(T)=\frac{e^2}{2\pi^2\hbar}\left(1+\frac{3\lambda^{(j=1)}}{4}\right)\ln\left(\frac{k_bT\tau}{\hbar}\right), \qquad (1)$$

where τ is the elastic relaxation time. The first universal term describes interaction between an electron and a hole with total spin $j=0$ and is due to the exchange (Fock) term while $\lambda^{(j=1)}$ is related to the direct (Hartree) term in the Hartree-Fock approximation of the Coulomb repulsion. In the absence of any attractive virtual potential between electrons, $\lambda^{(j=1)}$ depends only on the Fermi surface and on the screening length. The exchange term dominates the Hartree term, if the interaction potential is sufficiently smooth, i.e., its extension is larger than λ_F.[2]

For magnetic fields higher than $H_c=k_bT/g^*\mu_B$, the spin degenerescence is broken by Zeeman splitting, and the correction due to interaction becomes

$$\delta\sigma(T)=\frac{e^2}{2\pi^2\hbar}\left(1+\frac{\lambda^{(j=1)}}{4}\right)\ln\left(\frac{k_bT\tau}{\hbar}\right). \qquad (2)$$

Expressions (1) and (2) are valid for a diffusive motion,[3] and are modified when the cyclotron frequency $\omega_c=eH/m^*$ is comparable to the elastic relaxation time τ.

In this classically high-magnetic-field case, it is known that the tensor of conductivities is anisotropic:

$$\sigma_{xx}(\omega_c)=\frac{1}{1+(\omega_c\tau)^2}\sigma(\omega_c=0), \qquad (3)$$

$$\sigma_{xy}(\omega_c)=\frac{-\omega_c\tau}{1+(\omega_c\tau)^2}\sigma(\omega_c=0).$$

Houghton[4] has shown that

$$\delta\sigma_{xx}(\omega_c)=\delta\sigma(\omega_c=0), \qquad (4)$$

$$\delta\sigma_{xy}(\omega_c)=0.$$

Equation (4) is a general result which is valid for any dimensionality and any kind of interaction between electrons.[1] A a consequence the correction to the conductivity, *measured in a standart Hall bar geometry*, is $\delta\sigma(\omega_c)=\delta[(\sigma_{xx}^2+\sigma_{xy}^2)/\sigma_{xx}]=\delta\sigma(\omega_c=0)[1-(\omega_c\tau)^2]$, despite the fact that $\sigma(\omega_c)=\sigma(\omega_c=0)$. With Eq. (2), one finally obtains

$$\delta G=\delta\sigma=\frac{e^2}{2\pi^2\hbar}\left(1+\frac{\lambda^{(j=1)}}{4}\right)[1-(\omega_c\tau)^2]\ln\left(\frac{k_bT\tau}{\hbar}\right) \qquad (5)$$

This remarkable result is valid for $\omega_c\tau\lessgtr 1$, as demonstrated in Refs. 4 and 5. It means that the logarithmic correction to the conductance due to interaction increases steadily as a function of magnetic field, changing its sign at $\omega_c\tau=1$. Equation (5) also shows that oscillations of $\tau(H)$ with magnetic field (in the Shubnikov–de Haas regime) may give oscillations of δG.

The aim of this work is twofold: first we will show a direct experimental observation of the inversion of the correction. We will confirm the temperature [ln(T)] and magnetic field $[1-(\omega_c\tau)^2]$ dependences according to Eq. (5). Then we will use this fact to extract unambiguously the $[1+(\lambda^{(j=1)}/4)]$ term in both low and high classical magnetic fields, and find that the Hartree contribution is suppressed ($\lambda^{(j=1)}\simeq 0$), once $\omega_c\tau\geqslant 1$. To our knowledge such a result have been never reported up to now. Moreover, the substraction of the Shubnikov–de Haas oscillations permits

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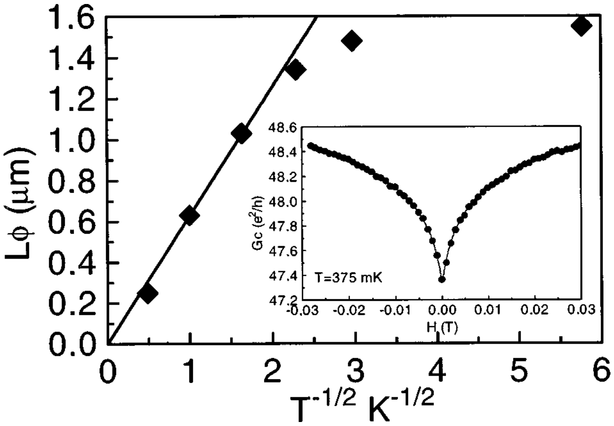

FIG. 1. L_ϕ vs temperature. The solid line is the prediction of Eq. (7). Inset: the low-field magnetoconductance at $T=375$ mK, with a weak localization fit.

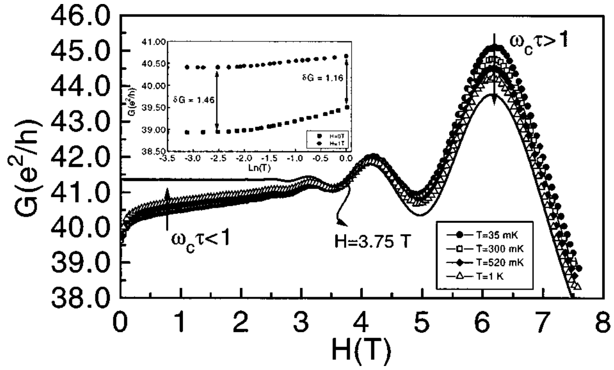

FIG. 2. $G(H)$ at various temperatures. The solid line is the fit of the Shubnikov–de Haas oscillations. Inset: the correction of G vs T at two magnetic fields (squares: $H=0$ T; circles: $H=1$ T). The absolute variation between these two fields is given by the weak localization and the Zeeman splitting terms.

us to show that the logarithmic term in Eq. (4) is effectively given by $\ln(k_b T\tau_{tr}/\hbar)$ at low magnetic field, where τ_{tr} is the transport relaxation time. However,the amplitude of this factor seems strongly reduced at high field.

Equation (5) means that $\delta\sigma$ increases in amplitude as the square of the magnetic field, leading eventually to an Hall insulating state characterized by $\sigma_{xx}\simeq 0$ and σ_{xy} constant. This prediction was studied by Murzin and Jansen[6] in three-dimensional (3D) doped semiconductors at high-magnetic-field. But the crossover at $\omega_c\tau\simeq 1$ has not been studied. A merit of our samples is to conjugate a relatively high disorder, giving a large EEI correction to the conductivity even at small fields, and a classical high magnetic field regime above 3 T. Electron interactions have been also studied in 2D high mobility GaAs heterostructures in Ref. 7. They observed that the correction to conductivity due to interaction varies like $(\omega_c\tau)^2$ [for $\omega_c\tau\geqslant 1$ and $T\geqslant 1$ K), and use that fact to study extensively the amplitude of the correction for various geometries. Our experiment differs from Ref. 7 because the samples are in the diffusive regime, where AAL theory is applicable. In addition, the Shubnikov–de Haas oscillations do not depend on temperature in our sample, because the elastic mean free path is much smaller (large Dingle temperature) than in Ref. 7 (and the experiment is performed at lower temperature). This permits us to extract the temperature dependence of the correction and not only the associated magnetoresistance. For this limitation the sign inversion predicted in Eq. (4) was not seen in Ref. 7.

We have used molecular-beam-epitaxy-(MBE)-grown GaAs doped at 2.2 10^{23} Si m^{-3}. Because our samples are based on a 300-nm-thick layer, in the low-temperature regime considered, samples are effectively two dimensional: both the phase breaking length and the thermal length $L_T=\sqrt{\hbar D/k_B T}$ (D is the diffusion constant) are larger than the thickness below 1 K. A 250×200-$\mu\mathrm{m}^2$ sample with Ohmic AuGeNi contacts is defined by etching. The system is characterized by the following parameters: $D=3.2\ 10^{-3}\ \mathrm{m}^2\mathrm{s}^{-1}$, $k_f l=6.5$, $\tau_{tr}=1.01\times 10^{-13}$ s, $E_f=240$ K, $a_b=95$ Å, and $R_c=492\ \Omega$ is the resistance per square.

To separate the EEI correction, we first analyze the weak-field magnetoconductance which is entirely due to the weak localization correction (see the inset of Fig. 1):[1]

$$\delta\sigma(H)=\frac{e^2}{2\pi^2\hbar}f_2\left[2\left(\frac{L_\phi}{L_H}\right)^2\right], \tag{6}$$

with $f_2(x)=\ln x+\psi(x+\frac{1}{2})$. $L_H=\sqrt{\hbar/2eH}$ is the magnetic length and $\psi(x)$ the digamma function.

Figure 1 shows the temperature dependence of the phase-breaking length $L_\phi=\sqrt{D\tau_\phi}$ (τ_ϕ is the phase-breaking time). For electron-electron interaction in two dimensions, Altshuler *et al.*[1] obtained

$$L_\phi=\sqrt{\frac{2\pi D\hbar^2}{k_b e^2 R_c \ln\left(\frac{\pi\hbar}{e^2 R_c}\right)}}T^{-1/2}. \tag{7}$$

With the measured sample parameters, we find $L_\phi(\mu\mathrm{m})=0.63\ T^{-1/2}$, in excellent agreement with the weak-localization measurement between 150 mK and 4 K (see Fig. 1). At very low temperature ($T\leqslant 150$ mK) a saturation is nevertheless observed, attributed either to high-frequency heating, to dephasing due to magnetic impurities, or to general electromagnetic environment considerations.[8]

In the intermediate magnetic field regime ($0.02\leqslant H\leqslant 0.5$) both Zeeman effect and weak localization make non-negligible and opposite contributions to the magnetoconductance. In addition, a crossover in the effective dimensionality occurs when the magnetic length is comparable to the sample thickness. For these reasons we do not fit the magnetoconductance in this intermediate regime. From the temperature dependence of the conductance correction both in zero magnetic field and for $H=1$ T, we determine self-consistently $\lambda^{(j=1)}$ (see inset of Fig. 2). In fact, above $H=1$ T the weak-localization contribution is negligible and the Zeeman level degeneracy breaking is effective for our lowest electron temperature. The conductance correction should obey to Eq. (5), i.e., the slope of $\delta G(e^2/h)$ versus $\delta\ln T$ (divided by $1-[\omega_c\tau(H)]^2$, that is, 0.929 at $H=1$ T) is given by $(1/\pi)(1/1.3)[1+\lambda^{(j=1)}/4)]$ (where the factor $d=1.3$ corresponds to the the length divided by the width of the sample). At zero magnetic field the same slope is given by $(1/\pi)(1/1.3)[1+(3\lambda^{(j=1)}/4)+1]$, where the last factor 1 is due to the weak localization term $(1/\pi)(1/1.3)\times(e^2/h)\ln(\tau_\phi/\tau)$ (with $\tau_\phi\propto T^{-1}$). The first evaluation gives

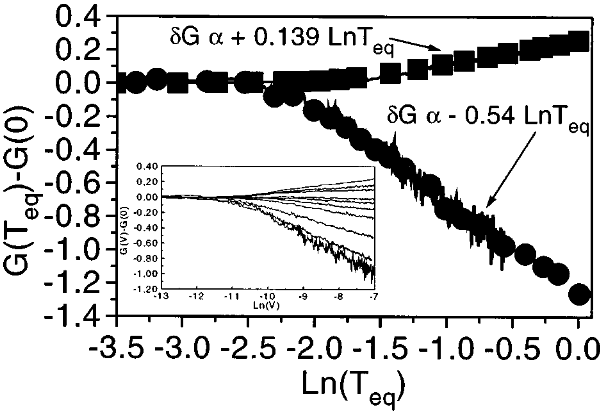

FIG. 3. Conductance vs temperature for $H=0.5$ T (squares) and $H=6.6$ T (circles). T_{eq} is the measured temperature for measurements in the linear regime or the temperature deduced from the bias-temperature relation discussed in the text for nonlinear conductance measurements (solid lines). Inset: $G(V)$ at various intermediate magnetic fields 1, 1.5, 3, 3.75, 4, 4.5, 5, 5.5, 6, and 6.6 T (top to bottom).

$\lambda^{(j=1)}\simeq -1.55+-0.1$, while the second estimation is compatible with $\lambda^{(j=1)}\simeq -1.2$. This corresponds to a strong screening case in $d=2$. The small discrepency may be related to a small spin splitting a zero magnetic field,[9] or to additionnal terms, for instance the Maki-Thomson term.

Moreover, the absolute magnetoconductance between $H=0$ and 1 T is well accounted by balancing the the weak-localization suppression and the Zeeman splitting effects (see the inset of Fig. 2):

$$G(1T)-G(0)\simeq \frac{1}{\pi}\frac{1}{1.3}\frac{e^2}{h}[\ln(\tau_\phi/\tau_{\mathrm{tr}})] + \frac{\lambda^{(j=1)}}{4}\ln[k_B T\tau_{\mathrm{tr}}/\hbar)]. \quad (8)$$

For instance at $T=1$ K, we find $G(1T)-G(0)\simeq 1.16(e^2/h)$, and we estimate $G(1T)-G(0)\simeq 1.34(e^2/h)$ ($T=1$ K, $\tau_\phi=1.32\times 10^{-10}$ s and $\lambda^{(j=1)}=-1.55$). Our value $\lambda^{(j=1)}$ corresponds to a screening larger than the estimation based on the Thomas-Fermi approximation,[1] but is not surprising considering the relatively high carriers concentration.

After elimination of the weak localization and Zeeman splitting effect, it is possible to investigate precisely the correction due to interaction above $H=1$ T. First one has to consider the $[1-(\omega_c\tau)^2]$ term in Eq. (4), which give two main effects: a change of sign for $\delta G(T)$ as $\omega_c\tau\simeq 1$ and oscillations of $\delta G(T)$ resulting from oscillations of $\tau(H)$.

Figure 2 shows the absolute magnetoconductance at various temperatures. One can see the change in the temperature dependence of the conductance at a magnetic field of about 3.75 T. This is confirmed in the Fig. 3, which details the correction to the conductivity versus temperature for two magnetic fields: $H=1$ and 6.6 T. The correction varies like the logarithm of the temperature as predicted by Eqs. (1) and (5). Note that the cancellation of the correction at $H=3.75$ T permits to determine precisely the Drude conductance: $G_{\mathrm{Drude}}=41.36(e^2/h)$.

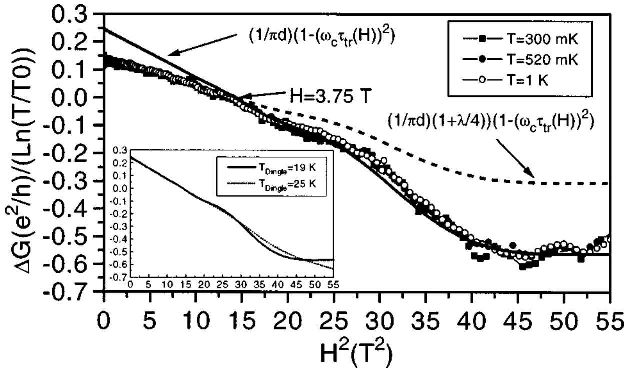

FIG. 4. The correction to conductivity due to electron interaction vs H^2. The base temperature T_0 is 150 mK; the symbols refer to different temperatures. The dashed and solid lines are, respectively, Eq. (5) with $\lambda^{j=1}\simeq -1.55$ (excellent at low field) and with $\lambda^{j=1}=0$ (excellent at high field). Inset: Two fits with $\lambda^{j=1}=0$ for two different Dingle temperatures, showing the sensitivity to this parameter.

Figure 3 also shows the conductance versus bias at $T\simeq 100$ mK. Due to finite electron-phonon coupling, the effective electronic temperature T' is increased above the T_0 phonon temperature at finite bias, according to the expression[10]

$$T'\simeq\left(T_0^5+\frac{V^2}{\Sigma\varrho L^2}\right)^{1/2}\propto V^{2/5}, \quad (9)$$

with ϱ the resistivity, L is the sample length (larger than the electron-phonon scattering length), and $\Sigma=0.524\alpha\gamma$, with $\tau_{\mathrm{ep}}^{-1}=\alpha T'^3$ and $\gamma=\pi^2\nu k_B^2/3$. ν is the density of states, and α is a numerical model-dependent constant. α and Σ characterize the electron-phonon coupling.

We use this formula to rescale the voltage as an effective temperature in Fig. 3: $T'=\beta V^{2/5}$ [$\beta(H=1$ T$)=15$, $\beta(H=6.6$ T$)=10$; this change is not explained]. These values of β correspond to $\Sigma\simeq 4\ 10^{-4}$ nW μm^{-3} K^{-5}, an estimation 10^{-4} lower than in metals.[10] This small electron-phonon coupling is essentially due to the low density of electrons.

We observe that the change in the temperature or bias dependence of the conductivity happens precisely when $\omega_c\simeq\tau_{\mathrm{tr}}^{-1}$ where τ_{tr} is the transport relaxation time (see Figs. 2 and 4). In that range of fields, the sample exhibits pronounced Shubnikov–de Haas oscillations periodic in $1/H$ (see Fig. 2), which permits us to determine the thermodynamic relaxation time τ_{thermo} to be 6.410^{-14} s.[11] We find that $\tau_{\mathrm{tr}}\simeq 1.5\tau_{\mathrm{thermo}}$. The diffusion by impurities is quite isotropic. Moreover, this time corresponds to a large value of the Dingle temperature: $T_D\simeq 19$ K. This value much larger than our experimental range [0.1 and 1 K] implies that the temperature changes of the conductance are strictly related to the EEI effects. This permits us to investigate the absolute values for exchange and direct terms (see Fig. 4). Indeed, by substracting conductance vs magnetic field at different temperatures T and T', we can estimate the term $\delta G=G(T)-G(T')=(1/\pi 1.3)[1+(\lambda^{(j=1)}/4)][1-(\omega_c\tau)^2]\ln(T/T')$. From low-field analysis we have obtained $\lambda^{(j=1)}\simeq -1.55+-0.1$, that corresponds to a relatively strong screening case, which makes the direct term comparable to the exchange term. Figure 4 shows that this estimation is valid up to $\omega_c\tau_{tr}\simeq 1$. But, as $\omega_c\tau_{tr}\gg 1$, the fit deviates strongly from the data. In this high-magnetic-field regime, we find that

$$1.3\delta G=\delta\sigma=\frac{e^2}{2\pi^2\hbar}[1-(\omega_c\tau_{tr})^2]\ln(T/T'), \qquad (10)$$

without any adjustable parameter: the direct term is destroyed ($\lambda^{(j=1)}\simeq 0$), and the correction is just given by the exchange part, qualitatively as if the screening becomes much less efficient.

The reason for the cancelation of the Hartree term needs to be clarified, taking into account that it happens as $\omega_c\tau \gg 1$. That suggests an orbital effect, perhaps due to the reinforcement of the forward scattering as compared to the backward ones: as $\omega_c\tau\gg 1$, the backward scattering $\Delta k\simeq 2k_F$ is diminished as compared to the forward scattering $\Delta k\simeq 0$. The direct (exchange) correction is proportionnal to $\Delta k \simeq 2k_F$ ($\Delta k\simeq 0$), that could explain our experimental observation.

To complete our analysis, we have substracted the Shubnikov–de Haas fit in order to extract the total correction to the conductance, i.e., to evaluate the absolute value of the $\ln(k_bT\tau/\hbar)$ term in Eq. (5). This is possible because of the excellent evaluation obtained for the other terms. At weak field, we verify that the absolute value of the correction agrees perfectly with the prediction of Eq. (5) with $\tau=\tau_{tr}$. But at higher magnetic field ($\omega_c\tau_{tr}\gg 1$), Eq. (5) predicts a larger correction than the one measured. A quantitative agreement is obtained if the term $(k_bT\tau_{tr}/\hbar)$ is multiplied by a factor 50. Note that this factor does not enter in the relative $\delta G=G(T)-G(T')$ measurement. This result suggests that departure from the diffusive regime—strictly valid only at low magnetic field— is accompanied by a strong absolute reduction in amplitude for the correction due to electron-electron interaction.

In conclusion, our diffusive GaAs sample exhibits large corrections due to disorder and interaction in zero magnetic field. Above $H\gg 1$ T, only interaction corrections due to exchange and Hartree terms of the screened Coulomb repulsion persist. These corrections leads to $\delta\sigma/\delta T<0$ at low temperature. When a high magnetic field is applied such that $\omega_c\tau_{tr}\simeq 1$, the temperature dependence changes its sign leading to $\delta\sigma/\delta T>0$, as predicted by Houghton, Senna, and Ying.[4,5] The whole functional dependence of the correction in $1-\omega_c\tau_{tr}(H)^2$ is obtained, including the the Shubnikov–de Haas oscillations of $\tau(H)$. We have been able to normalize the magnetoconductance curves at various temperatures, and we show that the Hartree term is canceled when $\omega_c\tau_{tr}(H)\gg 1$. Moreover, we have measured the absolute value for the interaction correction. Its predicted dependence is verified at low magnetic field, but when $\omega_c\tau_{tr}>1$ it is strongly reduced.

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