Determination of the electronic phase coherence time in one-dimensional channels

J. A. Katine,* M. J. Berry, and R. M. Westervelt
Division of Applied Sciences and Department of Physics, Harvard University, Cambridge, Massachusetts 02138

A. C. Gossard
Materials Department, University of California-Santa Barbara, Santa Barbara, California 93106

(Received 17 January 1997; revised manuscript received 29 September 1997)

We utilize the weak-localization effect in an array of one-dimensional channels to perform a quantitative measurement of the electronic phase coherence time at temperatures between 100 mK and 4.2 K in the two-dimensional electron gas formed in a GaAs/$Al_xGa_{1-x}As$ heterostructure. The measured times agree well with theories of dephasing based on electron-electron scattering. [S0163-1829(98)07103-3]

I. INTRODUCTION

The electronic phase coherence length l_ϕ determines the range over which quantum interference effects may be observed. As such, knowledge of this length is crucial for designing and analyzing experiments involving such interference phenomena. In GaAs/$Al_xGa_{1-x}As$ two-dimensional electron gas (2DEG) systems, previous experiments have used the weak-localization correction to the classical resistance of a quantum wire to determine the value of this length.[1–3] Although the results obtained from these experiments have been useful, these earlier measurements all possessed some limitation that precluded a precise determination of l_ϕ; for example, if the wires are too narrow, the phase coherence time is comparable to the elastic scattering time, in which case the theory used to extract l_ϕ from the localization correction is no longer valid.[1] Magnetoconductance fluctuations due to quantum interference can also limit the accuracy of the measurement.[3]

In this paper, we utilize the weak-localization effect to make a quantitative measurement of l_ϕ in narrow wires between 4.2 K and 100 mK, and find excellent agreement with the dephasing rate calculated from models based on electron-electron interactions.[4] Section II provides a brief introduction to the weak-localization effect. Section III summarizes the important experimental details of the work. Finally, in Sec. IV we discuss our results.

II. WEAK LOCALIZATION

The weak-localization effect is a quantum correction to the classical Drude conductance. This correction arises from the constructive interference of time-reversed pairs of backscattered electron trajectories, which increases the probability for backscattering, resulting in a decreased conductance.[5] The amplitude of this correction is proportional to the number of phase-coherent backscattered trajectory pairs. In a channel whose width W is shorter than l_ϕ, electron diffusion is constrained to one dimension and the amplitude of this correction is[6]

$$\delta G_{WL} = -\frac{2e^2}{h}\frac{D^{1/2}}{L}\left[\tau_\phi^{1/2} - \left(\frac{1}{\tau_\phi}+\frac{1}{\tau}\right)^{-1/2}\right], \qquad (1)$$

where L is the channel length, τ_ϕ is the phase coherence time, τ is the elastic scattering time, and $D=v_f^2\tau/2$ is the diffusion coefficient. Since motion along the channel is diffusive, the phase coherence length is related to the phase coherence time by the diffusion coefficient, $l_\phi=(D\tau_\phi)^{1/2}$.

Experimentally, the easiest way to determine δG_{WL} is to apply a magnetic field perpendicular to the sample. At sufficiently large fields, time-reversal symmetry is broken and the magnetoconductance saturates at its classical value:

$$\delta G_{WL} = G(0) - G(B > B_{\text{sat}}). \qquad (2)$$

This saturation field B_{sat} is reached when the magnetic flux enclosed by all of the time-reversed pairs is sufficient to dephase them. The half-width of the localization feature is extremely narrow ($\sim 1\ \mu T$) for our bulk 2DEG, but by constraining diffusion in the lateral dimension, we may broaden the width of the localization feature. In addition, when the channel width is narrower than the mean free path l, electron motion between the walls is ballistic, leading to a flux cancellation effect that further increases the characteristic field of the localization effect (~ 1 mT). Beenakker and van Houten (BvH) have developed a theory for the line shape of the magnetoconductance in this quasiballistic one-dimensional limit:[7]

$$\delta G(B) = -\frac{2e^2}{h}\frac{D^{1/2}}{L}\left[\left(\frac{1}{\tau_\phi}+\frac{1}{\tau_B}\right)^{-1/2} - \left(\frac{1}{\tau_\phi}+\frac{1}{\tau_B}+\frac{1}{\tau}\right)^{-1/2}\right]. \qquad (3)$$

In Eq. (3),

$$\tau_B = \frac{9.5 l_m^4}{W^3 v_f} + \frac{24 l_m^2 \tau}{5W^2}$$

for specular boundary scattering, where $l_m^2 = \hbar/eB$.

III. SAMPLE DESIGN AND MEASUREMENT

Our channels are formed by applying a negative bias voltage to a split gate on the heterostructure surface that depletes the electron gas located 420 Å below these gates. The 2DEG has $n_s = 3.9\times 10^{15}\ \text{m}^{-2}$, $\mu_e = 4.0\times 10^5\ \text{cm}^2/\text{V s}$, and $\tau = 15.2$ ps. The gates are patterned using electron-beam li-

0163-1829/98/57(3)/1698(5)/$15.00

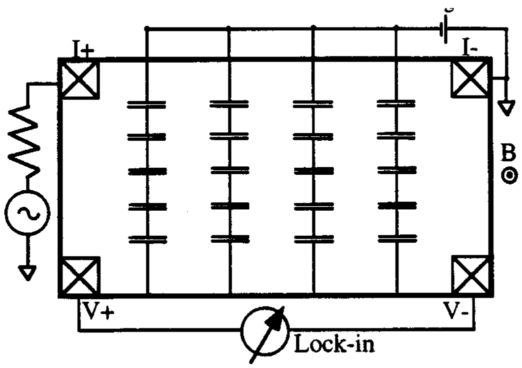

FIG. 1. Schematic of sample and four-probe measurement circuit. The dimensions of the 20 channels have been exaggerated for clarity.

thography followed by Cr/Au metallization. Figure 1 is a schematic overview of the sample design and measurement circuit. As indicated, annealed $Au_xNi_{1-x}Ge$ pads in the corners of the sample are used to make Ohmic contact to the electron gas. Greatly enlarged in the schematic, the individual channels are 105 μm long and 800 nm wide. The sample contains four identical columns in series, each of which consists of five parallel channels separated vertically by 500 μm. The metal wires connecting these parallel channels are negatively biased, allowing current to flow from source to drain only through the channels. All of the channels have the same negative bias applied to them, allowing simultaneous measurement of the resistance of the entire array. In performing four-probe low-frequency, current-biased lock-in measurements of the sample resistance, a small enough bias current (2 nA) is used to prevent hot carrier effects that can attenuate the phase coherence length.[8]

As alluded to in the Introduction, the array of channels is necessary for a quantitative measurement because of magnertoconductance fluctuations.[6] At sub-Kelvin temperatures, these manifestations of universal conductance fluctuntions are comparable in magnitude to the localization correction in channels of length less than l_ϕ. Unlike the localization feature, these fluctuations can be positive or negative at zero field, which allows them to be suppressed by making the channels longer than l_ϕ and by averaging over multiple channels.

The sample is oriented perpendicular to the magnetic field in a helium dilution refrigerator. To isolate the sample from radiation, all of the leads to the sample pass through cold metal film rf blocking resistors and two stages of Cu-powder filters.[9] Such precautions are necessary, because any radiation that reaches the electron gas can shorten the phase coherence time.

In quasiballistic channels, the nature of boundary scattering strongly influences the diffusion coefficient, and with it, the amplitude of the localization correction [Eq. (1)]. For specular boundary scattering, the diffusion coefficient is unchanged from its value in the bulk, because momentum along the channel is conserved during boundary collisions. Conversely, for perfectly diffuse boundary scattering, D is reduced. In the limit that $l \gg W$,[10]

$$D = \frac{v_f W}{\pi} \ln\left(\frac{l}{W}\right), \tag{4}$$

which implies that perfectly diffuse boundary scattering in our channels would reduce D to about 20% of its value in the bulk 2DEG.

Previous experiments have demonstrated reflections from electrostatically defined boundaries to be $>95\%$ specular.[11] From the present data, we can use the classical resistance of the channels to place a lower bound on the specularity. Consider a simple model[12] in which boundary collisions are fully specular with probability p, and fully diffuse with probability $1-p$. In this model, the resistance of a single channel is

$$R = \frac{L}{W}\,[p\rho_s + (1-p)\rho_d], \tag{5}$$

where ρ_s and ρ_d are the resistivities assuming perfectly specular and diffuse scattering, respectively: $\rho_s = m^*/e^2 n_s \tau$ and $\rho_d = \pi l \rho_s / 2W \ln(l/W)$. The sheet density in the channels can be directly measured from Shubnikov–de Haas oscillations ($n_s = 3.74\times 10^{15}\ \mathrm{m}^{-2}$). Although the small decrease in electron density from its value in the bulk likely results in a small reduction of the elastic scattering time in the channels, we may place a lower bound on ρ_s by using the bulk value for τ.

Unless noted otherwise, our measurements are made with a negative bias of 400 mV applied to the gates defining the channels, slightly larger than the depletion voltage (330 mV). Typically, the depleted region of the electron gas extends a short distance beyond the edge of the surface gate, so we expect that the channels are actually slightly narrower than their lithographic width of 800 nm. The most conservative estimate of boundary specularity, however, is made with Eq. (5) assuming W equal to the lithographic width, and establishes a lower bound of $p > 92\%$. If we assume this edge depletion width is equal to the depth of the 2DEG (giving $W = 700$ nm), our model finds $p = 96\%$. Finally, if we assume $p = 100\%$, the channel resistance places a lower bound on the width at 610 nm.

IV. RESULTS AND DISCUSSION

Figure 2(a) shows the magnetoconductance of the array at 4.2 K and 385 mK. In addition to the weak-localization correction centered at zero field, we observe a gradual negative magnetoconductance. Other investigators studying similar sample geometries have observed the same behavior as the result of classical size effects;[11] a small perpendicular field may deflect electrons whose velocity is parallel to channel walls into a boundary, and such collisions will lead to the observed negative magnetoconductance for even the small amount of diffuse boundary scattering present in our experiment.[13] At higher fields (~ 80 mT) we observe a crossover to a positive magnetoconductance, which is likely due to a magnetic suppression of backscattering.[6] In order to use Eq. (2) to determine τ_ϕ, these classical size effects must be separated from the localization correction. This is easy to accomplish because, as shown in Fig. 2(a), the localization correction decays rapidly with increasing temperature, while the classical size effects remain unaffected. From our higher-temperature measurements, we see that the magnetoconductance of the classical size effects over the magnetic field range of interest is well described by the parabolic fit shown

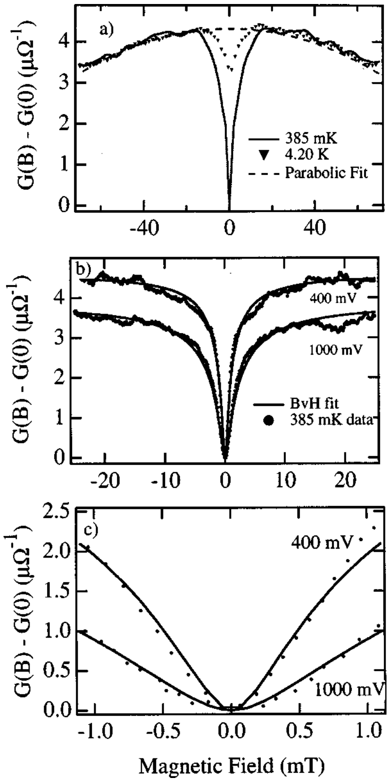

FIG. 2. (a) $G(B)-G(0)$ vs perpendicular field (400 mV bias) at 385 mK (solid) and 4.2 K (triangles). Dashed line is a fit to the classical background. (b) BvH theory [Eq. (3)] fit to 385-mK data following subtraction of the classical background for the array biased at 400 and 1000 mV. (c) Expanded view of low-field region.

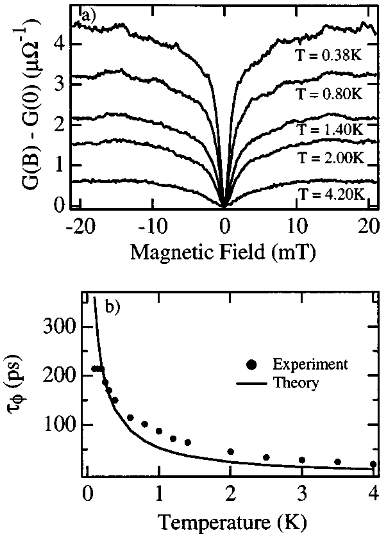

FIG. 3. (a) Representative traces demonstrating growth of the localization correction with decreasing temperature. (b) Dephasing times extracted from amplitude of localization effect plotted vs temperature. Solid line is calculated from Eq. (6).

in Fig. 2(a) (dashed line). Following subtraction of this classical background, we may use Eq. (2) to determine the amplitude of the localization correction.

Figure 2(b) is an enlargement of the 385-mK low-field magnetoconductance following subtraction of the parabolic background. After using Eq. (2) to find δG_{WL}, from the amplitude of the weak-localization correction Eq. (1) is used to extract the phase coherence time of $\tau_\phi = 150 \pm 5$ ps ($l_\phi = 9\ \mu$m). The quoted uncertainty comes from the magnetoconductance fluctuations present (though suppressed) despite averaging by the array. In using Eq. (1), perfectly specular boundary scattering is assumed in order to calculate D. As discussed earlier, nonspecular boundary scattering could impede diffusion along the channel, so the assumption of perfect specularity could result in a systematic undervaluing of τ_ϕ. Our lower bound on the specularity, however, assures us that any such systematic error is fairly small.

A second possible source of systematic uncertainty involves neglecting the magnetoconductance of electron-electron interactions.[14,15] At low temperatures, Coulomb interactions introduce quantum-mechanical correlations to the electron motion that affect the conductance. These e-e interactions are dominated by a magnetic-field-independent (but temperature-dependent) channel. There is, however, a small field-dependent component whose effects in our channel measurements are indistinguishable from that of weak localization, so neglecting these interactions in our analysis could inflate the inferred value of τ_ϕ. Fortunately, even at dilution refrigerator temperatures, the measured value of these field-dependent e-e interactions is less than 5% the size of the localization correction in the current measurement.[16]

Included in Fig. 2(b) and Fig. 2(c) is a fit to the BvH theory [Eq. (3)] for the line shape of the localization feature. The value of τ_ϕ obtained from the amplitude of the localization correction is held fixed, leaving the channel width as the only free parameter. The fit is restricted to fields less than ± 1 mT [the range shown in Fig. 2(c)], as the BvH theory is only strictly valid when $l_m > W$. This restriction exists because the lateral confinement may become irrelevant at fields large enough for dephasing to occur before the electron reaches the channel walls, effectively making the channel a two-dimensional system. As seen in Fig. 2(b), the fit is quite good over a much larger field scale. This is because the long mean free path and specular boundaries in our channels insure that, on average, several boundary collisions must occur before an electron may be backscattered and contribute to the localization effect. It is only when the width is larger than the mean free path that high fields induce a crossover to two-dimensional behavior when $l_m \sim W$. The fit for the 400-mV gate biased array shown in Fig. 2 yields a channel width of 725 nm, in excellent agreement with the range for the value of W inferred earlier from the classical resistance of the channels. Although previous work was able to qualitatively demonstrate that flux cancellation plays an important role in the localization line shape[1,2] the large uncertainty in the width of the etched channels used in those experiments prevented the quantitative check of the BvH theory possible in the current measurement.

As shown in Fig. 3(a), lower temperatures produce a larger amplitude localization correction as the result of the increased phase coherence length. Using the procedure described earlier, the amplitude of the weak-localization correction can be used to determine the phase coherence time at each temperature, and these values are plotted in Fig. 3(b) between 100 mK and 4.2 K. It is possible to study the temperature dependence of τ_ϕ in this manner, because the other quantities in Eq. (1), D and τ, are temperature independent below 4.2 K.[6] At these low temperatures, loss of phase coherence is believed to be the result of e-e scattering events between the injected electrons and those in the Fermi sea. In the limit of a one-dimensional channel, the dephasing rate calculated from such scattering is given by[4]

$$\frac{1}{\tau_\phi}=\frac{\pi}{2}\frac{(kT)^2}{\hbar E_f}\ln\left(\frac{E_f}{kT}\right)+\left(\frac{\pi kT}{D^{1/2}Wm^*}\right)^{2/3}. \tag{6}$$

The first term comes from momentum conserving processes involving large energy transfers. The second (Nyquist) term describes dephasing through many small energy transfer events that occur as electrons are scattered quasielastically by the fluctuating electric fields generated by the diffusive motion of other electrons. This Nyquist dephasing dominates below about 1.5 K for our sample parameters. Included in Fig. 3(b) is a plot of the calculated phase coherence times when the known parameters of our system are inserted into Eq. (6). Although our measured values for τ_ϕ are about 25% larger than those calculated, the agreement between the two is quite good considering that there are no adjustable parameters in the theory. This agreement strongly suggests that e-e interactions are indeed the dominant source of dephasing over this temperature range for 2DEG systems.

In order to examine the reproducibility of our results, we studied three samples fabricated on adjacent chips taken from the same heterostructure wafer. For all thermal cycles of all samples, virtually identical temperature dependencies are obtained for τ_ϕ, which is not surprising as the array is designed to average over the disorder that exists in the 2DEG. We also performed our measurements at different channel widths. This was accomplished by varying the bias voltage applied to the gated array. Changing the bias voltage by a small amount (25 mV) is sufficient to completely rearrange the appearance of the conductance fluctuations, but only changes the width of the channel by about 10 nm.[17] We find that such a small change in the channel width has a negligible impact on our measurements of τ_ϕ.

We have also examined the impact of more dramatic changes in the gate bias voltage. Included in Fig. 2(b) is data taken with the array gates biased at 1000 mV. Shubnikov–de Haas measurements show that the average sheet density in the channels is reduced to $n_s=3.38\times10^{15}\ \text{m}^{-2}$ at the larger bias voltage. Previous studies have found the elastic scattering time is roughly proportional to $n_s^{3/2}$.[18] Using this relationship to estimate the effect of the reduced sheet density on the elastic scattering time in our channels, the measured 76% increase in the classical resistance of the channels indicates that the width decreases from $W\approx700$ to ≈510 nm as the gate bias is increased from 400 to 1000 mV. Repeating the procedure outlined earlier, we find $\tau_\phi=135$ ps from the amplitude of the localization correction; a small decrease in τ_ϕ is expected for the narrower channel according to Eq. (6). We also fit the BvH theory, Eq. (3), to the data for the channels gate biased at 1000 mV. The fit remains excellent for the narrower channels, and the width $W=495$ nm extracted from the fit agrees well with the value inferred from our classical resistance measurements.

Finally, note that the measured dephasing rate shown in Fig. 3(b) saturates below 200 mK. As the temperature decreases from 200 to 100 mK, the measured value of the conductance continues to decrease due to the magnetic-field-independent e-e interactions alluded to earlier, but the amplitude of the weak-localization correction and the universal conductance fluctuntions remain unchanged. This is a clear indication that although the electron gas temperature continues to decrease, the phase coherence time has saturated.

Saturation of the weak-localization correction would be expected once the phase coherence length exceeds either the channel length or the Anderson localization length α^{-1}.[19,20] At saturation, however, $l_\phi=11\ \mu$m, which is significantly shorter than either the channel length or our 100 μm estimate for α^{-1}.[21] Therefore, we think it is likely that the saturation is due to high-frequency radiation coupling to electron gas despite the filtering present in the system. Calculations of the effect of such radiation on weak localization in 2DEG systems demonstrate that attenuation of the localization correction occurs in the presence of such radiation.[22] In the future, it would be interesting to experimentally test this hypothesis, especially in light of recent reports on phase coherence saturation at low temperatures.[23]

V. CONCLUSIONS

Using the weak localization effect in electrostatically defined quasiballistic one-dimensional channels, we have performed a quantitative measurement of the electronic phase coherence time. Over the temperature range examined, the measured dephasing rate agrees well with that predicted by theories based on e-e interactions. At very low temperatures, there is a saturation of the weak-localization correction, which may be due to phase breaking caused by external radiation.

ACKNOWLEDGMENTS

This work was funded at Harvard by ONR Grants Nos. N00014-95-1-0104 and N00014-95-0866, and by NSF Grants Nos. DMR 95-0-1483 and DMR 94-0-0396, and at UCSB by AFOSR Grant No. F49620-94-1-0158.

*Present address: School of Applied and Engineering Physics, Cornell University, Ithaca, NY 14853.

[1] H. van Houten, C. W. J. Beenakker, M. E. I. Broekaart, M. G. H. J. Heijman, B. J. van Wees, H. E. Mooij, and J. P. Andre, Acta Electron. **28**, 27 (1988).

[2] C. Kurdak, A. M. Chang, A. Chin, and T. Y. Chang, Phys. Rev. B **46**, 6846 (1992).

[3] J. A. Katine, M. J. Berry, R. M. Westervelt, and A. C. Gossard, Superlattices Microstruct. **16**, 211 (1994).

[4] B. L. Altshuler, A. G. Aronov, and D. E. Khmelnitsky, J. Phys. C **15**, 7367 (1982).

[5] S. Chakravarty and A. Schmid, Phys. Rep. **140**, 193 (1986).

[6] C. W. J. Beenakker and H. van Houten, Solid State Phys. **44**, 1 (1991).

[7] C. W. J. Beenakker and H. van Houten, Phys. Rev. B **38**, 3232 (1988).

[8] A. Yacoby, U. Sivan, C. P. Umbach, and J. M. Hong, Phys. Rev. Lett. **66**, 1938 (1991).

[9] J. Martinis, M. Devoret, and J. Clarke, Phys. Rev. B **35**, 4682 (1987).
[10] H. van Houten, C. W. J. Beenakker, B. J. van Wees, and J. E. Mooij, Surf. Sci. **196**, 144 (1988).
[11] T. J. Thornton, M. L. Roukes, A. Scherer, and B. P. van der Gaag, Phys. Rev. Lett. **63**, 2128 (1989).
[12] K. Fuchs, Proc. Cambridge Philos. Soc. **34**, 100 (1938).
[13] E. Ditlefsen and J. Lothe, Philos. Mag. **14**, 759 (1966).
[14] B. L. Altshuler and A. G. Aronov, in *Electron-Electron Interactions in Disordered Systems*, edited by A. L. Efros and M. Pollak (North-Holland, Amsterdam, 1985), p. 1.
[15] H. Fukuyama, in *Electron-Electron Interactions in Disordered Systems*, edited by A. L. Efros and M. Pollak (North-Holland, Amsterdam, 1985), p. 155.
[16] K. K. Choi, D. C. Tsui, and S. C. Palmateer, Phys. Rev. B **33**, 8216 (1986).
[17] This is known from measurements of electrostatically tunable electron interferometers, J. A. Katine *et al.*, Phys. Rev. Lett. (to be published).
[18] H. Z. Zheng, H. P. Wei, D. C. Tsui, and G. Weimann, Phys. Rev. B **34**, 5635 (1986).
[19] D. J. Thouless, Phys. Rev. Lett. **39**, 1167 (1977).
[20] R. G. Mani, K. von Klitzing, and K. Ploog, Phys. Rev. B **48**, 4571 (1993).
[21] M. Suhrke and S. Wilke, Phys. Rev. B **46**, 2400 (1992).
[22] B. L. Altshuler, A. G. Aronov, and D. E. Khmelnitsky, Solid State Commun. **39**, 619 (1981).
[23] P. Mohanty and R. A. Webb, Phys. Rev. B **55**, 13 452 (1997); P. Mohanty, E. M. Q. Jariwala, and R. A. Webb, Phys. Rev. Lett. **78**, 3366 (1997).