Mesoscopic decoherence in Aharonov-Bohm rings

A.E. Hansen, A. Kristensen, S. Pedersen, C.B. Sørensen, and P.E. Lindelof

The Niels Bohr Institute, University of Copenhagen, Universitetsparken 5, DK-2100 Copenhagen, Denmark

(October 24, 2018)

We study electron decoherence by measuring the temperature dependence of Aharonov-Bohm (AB) oscillations in quasi-1D rings, etched in a high-mobility GaAs/GaAlAs heterostructure. The oscillation amplitude is influenced both by phase-breaking and by thermal averaging. Thermal averaging is important when the temperature approaches the energy scale, on which the AB oscillations shift their phase. For the phase-breaking, it is demonstrated that the damping of the oscillation amplitude is proportional to the length of the interfering paths. For temperatures $T$ from 0.3 to 4 K we find the phase coherence length $L_φ \propto T^{-1}$, close to what has been reported for open quantum dots. This might indicate that the $T^{-1}$ decoherence rate is a general property of open and ballistic mesoscopic systems.

PACS numbers: 73.23-b, 73.63.Nm

The understanding of decoherence in quantum mechanical systems gives valuable insight into the cross-over from quantum to classical behavior. Quantum phenomena like weak localization, universal conductance fluctuations and the Aharonov-Bohm effect, that are observed in mesoscopic electronic systems, make these systems well suited for studying decoherence. The loss of electron phase coherence is interesting in its own right because it reveals information about the fundamental physics of the electron scattering mechanisms. Moreover, from the perspective of possible phase-coherent mesoscopic electronic devices, knowledge of phase-breaking length and time scales is crucial.

At low temperatures, electron-electron scattering is usually the dominating source of phase-breaking. In disordered 1D and 2D conductors, the loss of phase coherence at low temperatures has been studied intensively, both theoretically and experimentally. In clean electron systems, the number of investigations are few. In 2D, experiments consistent with the expected electron-electron scattering time $τ_0 \sim (T^2\ln T)^{-1}$ has been carried out. In open quantum dots (a 1D system), an unexpected $T^{-1}$ contribution was found. In general, phase breaking mechanisms in ballistic, mesoscopic systems of dimensionality less than 2, are presently not well understood.

Aharonov-Bohm (AB) rings are obvious systems for probing phase coherence. Here, the interference of two electron paths leads to conductance oscillations of period $h/e$ (or frequency $e/h$) in the magnetic flux enclosed by the paths. The oscillation amplitude is a direct measure of the interference strength, and it has been used to study decoherence in disordered systems. For AB rings with a 2DEG elastic mean free path longer than the circumference of the device (e.g. and references therein), systematic studies of phase-breaking have been scarce.

In this paper, we report measurements of the phase coherence length $L_φ$ via the temperature dependence of AB conductance oscillations in quasi-1D rings, made by shallow etching in GaAs/GaAlAs heterostructures. Two mechanisms are important for the temperature dependence of the oscillation amplitude: phase-breaking, and thermal averaging. At finite temperature, the measured conductance is a weighted average over an energy interval of finite width, proportional to the temperature. We discuss, how thermal averaging influences the AB oscillation amplitude through the phase changes of the oscillations. In the experiment, we detect AB oscillations due to the interference of electrons that encircle the ring up to $n = 6$ times. This is observed as peaks in the Fourier spectra of the magnetoconductance at multiples $ne/h$ of the fundamental AB frequency. Accounting for the effect of thermal averaging, we find that the damping of the amplitude due to phase-breaking depends linearly on $n$, showing directly the relaxation nature of the decoherence. We find the phase coherence length $L_φ \propto T^{-1}$.

The AB rings are fabricated in a two-dimensional elec-
electron gas (2DEG) situated 90 nm below the surface of a modulation doped GaAs/GaAlAs heterostructure. At liquid He temperatures, the unpatterned 2DEG density and mobility is $n = 2.0 \times 10^{15} \text{ m}^{-2}$ and $\mu = 80 \text{ m}^2/\text{Vs}$, corresponding to an elastic mean free path $l_e = 6 \mu\text{m}$. The lateral confinement is obtained by a shallow wet-etch, and the device is covered by a metal gate electrode. Details on the sample fabrication have been presented elsewhere. The sample was cooled in a $^4$He cryostat, and the conductance was measured in a two-terminal configuration. We used a conventional voltage biased lock-in technique, with an excitation voltage of $31.6 \mu\text{V}$ oscillating at a frequency of $116.5 \text{Hz}$.

Two samples with identical designs have been investigated in detail. Here we present measurements on one of them. The main results are reproduced in the second sample. A Scanning Electron Microscope (SEM) image of the ring is shown in Fig. 1a. The ring has a circumference of $3 \mu\text{m}$, $< l_e$. The gate voltage dependence of the conductance is shown in Fig. 1b. At $T = 5.8 \text{K}$, the conductance increases in steps of height $\sim e^2/h$, due to the conductance quantization in the narrow exit and entrance wires. The steps should not be taken as a sign of the population of transverse subbands in the ring itself, where the wires defining the arms are wider. At $T = 0.3 \text{K}$, a UCF-type signal is superimposed on the steps, indicating a large degree of phase coherence in the ring at this temperature.

In Fig. 1b we show an example of the temperature evolution of the conductance in a perpendicular magnetic field $B$. The AB oscillation amplitude decreases with temperature. The period of the oscillations is $5.4 \text{mT}$, in agreement with the period $h/(e\pi r^2) = 5.5 \text{mT}$ calculated from the average radius $r = 490 \text{nm}$ of the ring.

To quantify the degree of coherence in the ring from the AB oscillations, we compute the Fast Fourier Transform (FFT) of the conductance $G(B)$, measured in the magnetic field interval as shown in Fig. 1b. A measure of the AB oscillation amplitude is obtained by integrating the $ne/h$ FFT peaks in the intervals as indicated in Fig. 2a. The amplitudes vary with gate voltage, as exemplified for the $e/h$ peak in Fig. 2a. It is sensitive to the phase difference $\Delta(k_F L)$ at zero magnetic field between electron paths in either arm of the device. The typical phase pickup in an arm is $k_F L \sim 200$, where $L = 1.5 \mu\text{m}$ is half the circumference of the ring and $k_F$ is the Fermi wave number $\sim 1.5 \times 10^6 \text{m}^{-1}$. A shift in the phase of the AB $h/e$ oscillation, between the two values 0 and $\pi$ that are possible in a two-terminal measurement, requires $\Delta(k_F L)$ to change by $\pi/2$. For a real device, $k_F L$ will not increase in exactly the same manner for the two arms, and hence the phase and the amplitude depends on the gate voltage, as we observe. The presence of several transverse subbands in the arms will provide an additional source of amplitude and phase variation.

The temperature dependence of the AB oscillation amplitude is shown in Fig. 2d for 10 different gate voltages. There is some scatter of the data points, around a well-defined average. The temperature dependence of the oscillation amplitude does not depend strongly on gate voltage, in the gate voltage interval used here.

We take advantage of the weak gate voltage dependence, and perform an average of the Fourier spectra obtained at different gate voltages. In Fig. 2b we show Fourier spectra averaged over the 10 gate voltages, for different temperatures. We also show an average Fourier spectrum computed from spectra obtained at 500 different gate voltages at $T = 0.32 \text{K}$. Apart from the peak at the $e/h$ Aharonov-Bohm frequency, clear peaks appear at frequencies $2e/h$, $3e/h$, $4e/h$, and smaller bumps are also visible at $5e/h$, $6e/h$. Electrons can travel around the ring more than once, and a periodicity of $h/ne$ in flux of the conductance means that the interfering electron has enclosed the ring $n$ times. The probability for this type of event to happen will decrease with $n$, and so will also the amplitude of the $ne/h$ oscillation (as seen in Fig. 2b), with a rate that depends on the coupling between the ring and the 2DEG.

In Fig. 3a we show the temperature dependence of the amplitude of the $h/ne$, $n = 1 \ldots 6$ oscillation periods, extracted from the average spectra as shown in Fig. 2b. For high temperatures and frequencies, the FFT amplitudes collapse onto a temperature-independent background spectrum. For temperatures between 0.3 and $4 \text{K}$, the amplitude drops exponentially with temperature for $n = 1 \ldots 4$, but with different rates $b_n$ for different
as has been motivated above. Therefore, the relevant interfering paths changes as function of the Fermi energy, geometrical phase difference, \( \Delta(\phi) \), between the two interfering paths changes as function of the Fermi energy, as has been motivated above. Therefore, the relevant temperature scale on which thermal averaging becomes efficient, is given by the Fermi energy change required to shift the phase of the AB oscillations, rather than for instance the energy level spacing of ring eigenstates.

We use Eq. (3) to estimate the effect of thermal averaging in our experiment. For \( G(E^F, 0, B) \) we use a data set \( G(V_g, T = 0.32K, B) \), where \( V_g \) is changed in steps of 0.6 mV, small enough to resolve all the changes in the AB oscillations. The relation of gate voltage \( V_g \) to Fermi energy \( E^F \) is calibrated by following the spiky features seen on the low-temperature conductance curve in Fig. 1b for finite bias voltages, \( V_{bd} \). They will for small biases move linearly in the \((V_g, V_{bd})\) - plane with a slope \( \delta V_{bd}/\delta V_g \), which is related to the Fermi energy change with gate voltage as \( \delta E^F/\delta V_g = e/(2(\delta V_{bd}/\delta V_g)) \). The extracted \( \delta E^F/\delta V_g \) is close to a simple capacitor estimate. From the simulated data sets \( G(V_g, T > 0.32K, B) \), the average AB oscillation amplitude is extracted in the same manner as for the measured data. The result is shown in Fig. 3b. The calculated oscillation amplitudes do decrease with temperature, but much slower than the measured amplitudes. We conclude that thermal averaging alone can not account for the measured data. Furthermore, there is a clear tendency that the calculated \( h/ne \) amplitudes decay faster for \( n \) odd than even. For \( n \) even, the magnetoconductance oscillations result partly from interference between time-reversed paths, the same paths that ultimately give rise to Altshuler-Aronov-Spivak oscillations in disordered systems. The geometrical phase difference of these paths is zero, which means that the resulting part of the magneto-oscillations does not change phase. Consequently they are insensitive to thermal averaging. This explains why the even frequencies in Fig. 3b have a slower temperature dependence than the odd frequencies. The temperature dependence can be approximated by exponentials, lines in Fig. 3b. For diffusive systems, thermal averaging is not important for temperatures below \( E_c/k_B \approx 15-20 \) K for our system, larger than the measurement temperatures. Here \( E_c = h^2 \pi^2 D/(2L)^2 \) is the standard expression for the correlation energy, and the diffusion coefficient \( D = v_F^2 \tau \) in 1D. But numerical calculations on a ballistic ring have given an exponential temperature dependence for low temperatures.

We can now return to the experimental results in Fig. 3b. As shown, the measured \( h/e \) (\( h/3e \)) decay rate has a larger contribution from thermal averaging than the \( h/2e \) (\( h/4e \)) decay rate. We see that the part of the damping that thermal averaging can not account for, approaches the scaling with \( n \) as foreseen in Eq. (1). To demonstrate this, we show in Fig. 4a the measured \( h/ne \) decay rates vs. \( n \), and in 4b the estimated decay rates due to thermal averaging. The measured rates \( b_n \) are not directly proportional to \( n \). But the \( h/2e \), \( h/4e \) decay rates, which are only little influenced by thermal broadening (since \( b_2/d_2, b_4/d_4 \approx 8 \), do obey the scaling (straight line). Even the \( h/6e \) decay rate extrapolated from the scaling agrees with the data (dashed line in Fig. 4b). From the data in Fig. 4a-b it seems plausible, that with a proper deconvolution of the thermal averaging from the phase-
from these two rings, we have found interference ampli-
0.4. The same analysis performed on data from the
∼
interference amplitude
 exponential decay. The exponents
f
property of phase breaking: the damping of the
amplitude scales with the length of the interfering paths. We
have detected oscillation amplitudes resulting from Eq. (1) that
L
with the slope 0.3 K
ne/h
oscillations shows again an ex-
only slightly smaller than the ones obtained from the
ensemble averaged spectra, showing again that these os-
cillary oscillations are only little influenced by averaging. Finally,
the fits and the Fermi energy calibration. c) and d) as in a)
and b), for another sample.

This has to our knowledge not been done before. As a
further cross-check of the analysis we have also consid-
ered the amplitude of averaged, measured AB oscillations
(b) Exponents𝛼
observed and verified a basic
Ensemble averaged spectra, showing again that these os-
cillations are only slightly smaller than the ones obtained from the
measured data
n
0
from the measured data
n
0
Error bars on the symbols take into account the standard deviations on the fits. Straight line: fit with
bn=α· n, α = 0.3 K
b* (K
f
b
b
b
C.J.B. Ford, A.B. Newbury, M. Pepper, H. Jones, Phys. Rev. Lett.
15356 (2000)
A.G. Huibers, M. Switkes, C.M. Marcus, K. Campman, A.C. Gossard, Phys. Rev. Lett. 81, 200 (1998); A.G.
Huibers, S.R. Patel, C.M. Marcus, P.W. Brouwer, C.I. Duruisz, J.S. Harris, Jr., Phys. Rev. Lett. 81, 1917 (1998)
C. Hodges, H. Smith, J.W. Wilkins, Phys. Rev. B 4, 302 (1971); G.F. Giuliani, J.J. Quinn, Phys. Rev. B 26, 4421
(1982); H. Fukuyama, E. Abrahams, Phys. Rev. B 27, 5976 (1983)
8 See e.g. S. Washburn, R.A. Webb, Rep. Prog. Phys. 55, 1311 (1992); C. Kurdak, A.M. Chang, A. Chin, T.Y.
Chang, Phys. Rev. B 46, 6846 (1992)
9 C.J.B. Ford, T.J. Thornton, R. Newbury, M. Pepper, H. Ahmed, D.C. Peacock, D.A. Ritchie, J.E.F. Frost, G.A.C.
Jones, Appl. Phys. Lett. 54, 21 (1989); C.J.B. Ford, A.B. Fowler, J.M. Hong, C.M. Knoedler, S.E. Laux, J.J. Wainer,
S. Washburn, Surf. Sci. 229, 307 (1990)
10 J. Liu, W.X. Gao, K. Ismail, K.Y. Lee, J.M. Hong, S. Washburn, Phys. Rev. B 50, 17383 (1994)
11 G. Cernicchiaro, T. Martin, K. Hasselbach, D. Mailly, A. Benoit, Phys. Rev. Lett. 79, 273 (1997)
12 S. Pedersen, A.E. Hansen, A. Kristensen, C.B. Sørensen,
The energy spacing $\Delta E_{0,1}$ of the first transverse subband in the exit and entrance wires is $\sim 9$ meV. For the arms of the ring, we estimate $\Delta E_{0,1}$ to 1-3 meV.

It can be argued, that $L$ should be a circumference rather than the length of an arm. In that case, we underestimate $L_\phi$ by a factor of 2.