New Nonlocal Magnetoresistance Effect at the Crossover between the Classical and Quantum Transport Regimes

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A new type of nonlocal oscillatory magnetoresistance effect is observed in heavily doped n^+-GaAs wires over a limited range of temperature ($10 \leq T \leq 50$ K) and at magnetic fields sufficiently large to give rise to Landau quantization. The effect is observed when quantum ballistic transport along the edges coexist with diffusive and dissipative conduction in the bulk.

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The observation of nonlocal effects in the resistance and magnetoresistance of small, multiterminal metallic and semiconducting structures has provided new insights into the physics of electrical conduction [1-7]. Nonlocal resistance is the term used to describe the appearance of a potential difference in regions of the structure which are well separated from the classical current paths. To date, observations of this effect fall into three distinct categories: (1) nonlocal universal conductance fluctuations (UCF) in "dirty" semiconductors [2] and metals [3] in which the electrical conduction is diffusive; (2) ballistic transfer resistance in two-dimensional electron gases (2DEGs) [1,4]; and (3) nonlocal Shubnikov–de Haas (SdH) oscillations in the quantum Hall effect regime of 2DEGs [5-7]. In all three categories the nonlocal effect is best seen at low temperatures. For nonlocal UCF to be observed the electron phase-coherence length l_ϕ, which decreases with increasing temperature, must be of the same order as the distance between the current and voltage probes. In the other two cases, high electron mobilities are required.

In this Letter we report the existence of a new type of nonlocal oscillatory magnetoresistance under conditions in which diffusive transport would be expected to occur but which is not reliant on the phase coherence of the electrons. The effect has two striking features. First, it is completely *quenched* at low temperatures ($T < 4$ K) and, second, it corresponds to a 100% modulation of the nonlocal magnetoresistance, even though the local SdH oscillations are relatively weak. We attribute this effect to the coexistence of classical bulk conduction and quantum conduction in edge states [8]. The edge states are able to carry the current to classically inaccessible regions, but the voltage appears only when there is coupling between the edge states and bulk with subsequent dissipation. At low temperatures the dissipation is reduced and this causes the quenching of the effect in sharp contrast to high-mobility systems where the dissipation occurs in the contact reservoirs.

The structure used in our experiment is shown schematically in Fig. 1(a) (inset). It was fabricated using electron-beam lithography and dry etching from an n^+ GaAs epilayer grown by molecular-beam epitaxy on a semi-insulating substrate. The electron concentration n is approximately 1.1×10^{24} m^{-3}. We estimate an effective conducting channel thickness of ~30 nm with four electrically quantized, two-dimensional subbands occupied. Channel widths w of 150, 250, 350, and 450 nm, all allowing for sidewall depletion, have been investigated. Adjacent pairs of probes are separated by $L = 1$ μm, although there are additional pairs (not shown in Fig. 1) beyond e and f with separation 2 and 50 μm. Each voltage probe has the same thickness and width as the main wire and is at least 7 μm long before opening out to a large contact pad. We refer to this pad as the *contact* and the length of wire between it and the main wire as the *probe*. The electron mobility μ is 0.16 m^{2}V^{-1}s^{-1}, giving a bulk elastic mean free path λ_B of 50 nm and $\omega_c\tau \sim 1$ at 6 T. The bulk resistivity of the material $\rho \approx 800$ $\Omega\square^{-1}$. Weak localization and UCF measurements give $l_\phi \sim 0.3$ μm at 4.2 K, decreasing as $T^{-1/2}$ at higher temperatures.

We have measured the nonlocal resistance $R_{ij,kl}$, where V_{kl} is the voltage difference between contacts k and l due to a current I_{ij} between contacts i and j. Standard low-frequency (12 Hz) lock-in techniques were employed with currents down to 50 nA to avoid electron heating [9]. At all temperatures, the measured resistance was checked to be independent of frequency and driving current. For voltage probes a distance L from the current probes, a and b in our devices, the expected classical voltage drop is $V_{kl} = \rho I_{ab}\exp(-\pi L/w)$. We present data only for $w = 150$ and 250 nm where this voltage is vanishingly small ($V/I < 0.01$ Ω). Magnetoresistance $R_{ab,cd}$ traces are shown in Fig. 1(a) for three temperatures and $w = 150$ nm. In this configuration $L = 1$ μm. For $T < 10$ K nonlocal UCF are clearly visible; their rms amplitude, ΔR, is well described by $\Delta R \sim \rho^2(e^2/h)\exp(-aL/l_\phi)$, with $a = 1.1 \pm 0.2$ and a temperature-dependent l_ϕ [2]. In low magnetic fields ($B < 3$ T), with $\omega_c\tau < 1$, the Lee-

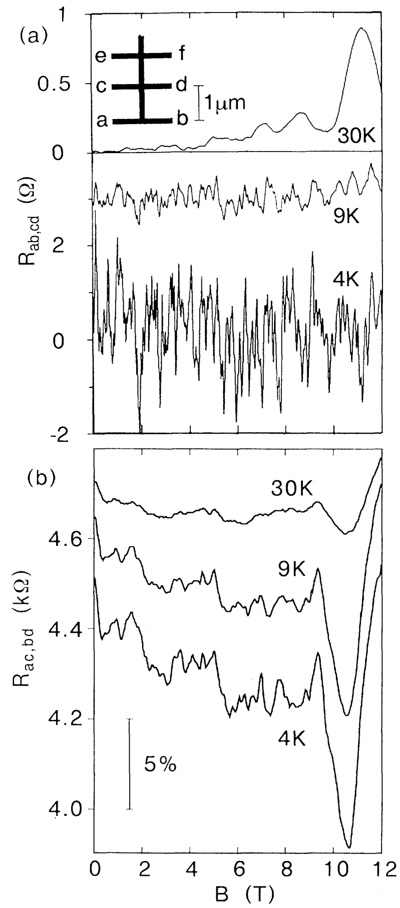

FIG. 1. (a) The nonlocal magnetoresistance traces for $L=1$ μm, $w=150$ nm at different temperatures. The 9-K trace is offset upwards by 3 Ω for clarity. The 30-K trace has threefold magnification. Inset: The sample geometry. (b) Local magnetoresistance $R_{ac,bd}$ at three temperatures. The 9- and 30-K traces are offset upwards by 0.2 and 0.4 kΩ, respectively, for clarity.

Stone correlation field [10], ΔB, is

$$\Delta B = \beta(h/e)/l_\phi[\min(w,l_\phi)] ,$$

with $\beta \sim 0.4$. For $\omega_c\tau > 1$, ΔB increases, as can be seen clearly in the raw data of Fig. 1(b). This increase is not predicted by theory [10] and will be the subject of a future paper.

For $T > 10$ K the UCF are exponentially damped and we observe a new type of oscillation as shown in Fig. 1(a) with threefold magnification, for $T=30$ K. In zero magnetic field the nonlocal resistance is zero within experimental uncertainty, the new peaks occurring only with finite magnetic field. Local SdH oscillations measured in the configuration $R_{ac,bd}$ are shown in Fig. 1(b). Their amplitude decreases monotonically with increasing temperature. Note also that at $T=30$ K the amplitude of the local oscillation around 10 T is only $\sim 1\%$ of the average resistance. Furthermore, peaks in the local magnetoresistance do not, in general, coincide with peaks in the nonlocal configuration. This means that the measured nonlocal resistance cannot be due to an artifact involving the "leakage" of the local resistance. It also distinguishes our effect from the nonlocal SdH oscillations investigated in high-mobility 2DEGs [5-7]. Nonlocal resistance oscillations are seen in all our samples and for many different probe configurations. The amplitude of the peaks decreases with increasing L and no effect is observable when the probes are separated by 50 μm.

Figure 1(a) shows clearly that the nonlocal oscillations do not increase in size in cooling from 30 to 9 K, although the effect is masked by the presence of the UCF at 9 K. For comparison, the local SdH oscillations increase by a factor of 4 in the same temperature interval [Fig. 1(b)]. The behavior with temperature is shown more clearly in Fig. 2 where $R_{ab,ef}$ (i.e., $L=2$ μm) at high magnetic field is shown for five different temperatures. In this configuration the UCF are sufficiently small, even at helium temperatures, to allow a quantitative study of the new effect. Figure 3 shows the amplitudes of the peaks at 8.7 and 11 T plotted against temperature. The amplitude is defined as the actual value of the resistance at the relevant field. Figures 2 and 3 show clearly that the effect disappears at both low and high temperatures. The inset in Fig. 3 shows the temperature dependence of the relative amplitude of the *local* SdH oscillations for a section of length 1 μm. The solid line in

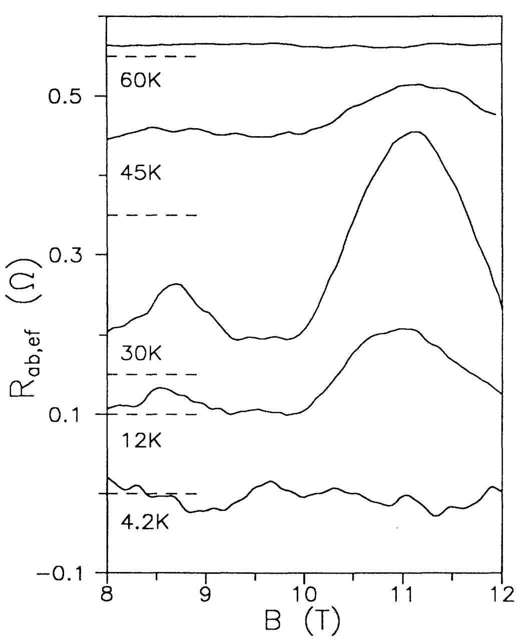

FIG. 2. The nonlocal high-field magnetoresistance $R_{ab,ef}$ measured at voltage probes with $L=2$ μm and $w=150$ nm. The dashed lines represent the respective zero-field resistances at the five temperatures.

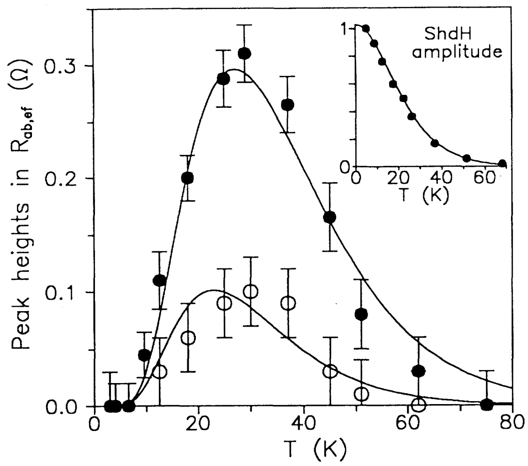

FIG. 3. The temperature dependence of $R_{ab,ef}$ for the two peaks shown in Fig. 2. Solid lines are the best fits to the temperature dependence using Eq. (1); circles are the experimental data. Inset: The temperature dependence of the amplitude of local SdH oscillations (the solid line represents theory).

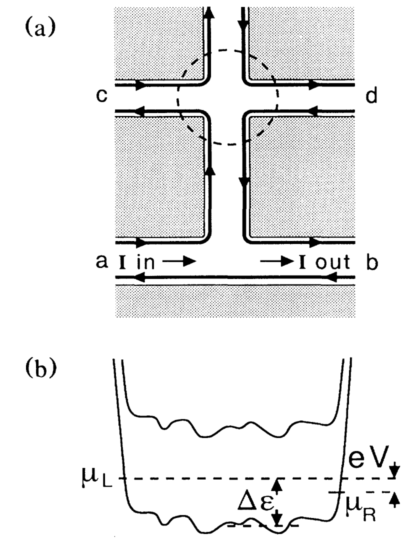

FIG. 4. (a) Schematic diagram showing the edge currents and the region of dissipation (dashed line) to which the nonlocal resistance $R_{ab,cd}$ is sensitive. (b) Landau levels showing the origin of the activation term in Eq. (1).

the inset is the expected temperature dependence [11] $[(k_BT/\hbar\omega_c)\text{csch}(2\pi^2k_BT/\hbar\omega_c)]$.

The considerations above are generally true for all nonlocal configurations. However, individual nonlocal magnetoresistance traces may be different from each other. Peak amplitudes are not generally the same for equal probe spacing, although generally the amplitude decreases as L increases. Also, the position in magnetic field of the nonlocal oscillations depends on the direction of magnetic field, particularly for an asymmetric probe configuration such as $R_{ab,ed}$. However, the Onsager condition, $R_{ij,kl}(B)=R_{kl,ij}(-B)$, is always satisfied for all configurations. In contrast, the local SdH oscillations have no change in phase and only very small changes in amplitude for different sections of the wire. Despite the differences in phase, the periodicity of the oscillations is the same in both the local and nonlocal configurations.

It is well known that in high-mobility 2DEG systems the edge states are decoupled from the bulk. We may consider the electrons in the edge states to have a mean free path, λ_E, which in those systems may be of the order of 1 mm [5]. In our low-mobility structure we may still expect $\lambda_E \sim 1$ μm if we scale λ_E with the bulk mean free path λ_B. This is reasonable if we assume that both λ_E and λ_B scale with the density of scattering centers. The large difference between λ_E and λ_B induces a step, $\Delta\mu_{EB}$, between chemical potentials at the edge and the bulk states which is proportional to the applied current [12]. We attribute the observed effects to this absence of local equilibrium between edge and bulk states.

Referring to the inset in Fig. 1(a), to measure the voltage difference between contacts c and d, when the current is applied between contacts a and b, involves dissipation occurring somewhere between c and d. In the high-mobility case the dissipation of the edge currents occurs in the *contact* regions. This cannot happen in our low-mobility structures because the edge electrons which are injected sideways from the main current path cannot reach the voltage contacts. Instead they are either transferred to the opposite edge without energy losses or thermalized within the bulk somewhere in the wire and probe regions. The voltage in our experiment, therefore, measures the dissipation which occurs in and around the *probe* regions. The region where the voltage V_{cd} is sensitive to the thermalization process is shown schematically in Fig. 4(a) as the dashed circle. We emphasize that this is totally different from the effect described, for example, in Refs. [5,6]. In particular, at low temperatures, our effect disappears when the electrons entering the bulk region can traverse the width of the wire without suffering an inelastic collision and dissipating energy.

We expect the voltage drop V_{cd} to have the form

$$V_{cd} \propto I_{ab}P(T)\exp(-L/\lambda_E)f(T)/\lambda_E\,. \tag{1}$$

The factor $f(T)=\text{csch}(2\pi^2k_BT/\hbar\omega_c)$ is the same factor which limits the bulk SdH oscillations at high temperatures. We assume that the potential drop V_{cd} is proportional to the number of electrons $[\sim\lambda_E^{-1}\exp(-L/\lambda_E)]$ leaving the edge at a distance L from the current path ab. Those that reach the opposite edge without losing energy will not give rise to this potential drop and only the fraction $P(T)$ that thermalize will contribute. Thermalization involves a transition from edge to bulk followed by a relaxation to local equilibrium. Using the energy balance equation, we can write $P(T)\propto[1+\exp(\Delta\epsilon/k_BT)]^{-1}$, where $\Delta\epsilon$ is the difference between the Fermi energy and the energy of the maximum of the density of states in the bulk [see Fig. 4(b)]. We expect that λ_E will be a maximum when the Fermi energy is between bulk Landau levels, but at this magnetic field the dissipation term $P(T)$ will be small. On the other hand, there will be strong coupling between edge states and bulk when the Fermi energy lies within a bulk Landau level but λ_E will be correspondingly small. The maximum values of V_{cd}, given by Eq. (1), have the same periodicity but not necessarily the same phase as the Landau levels crossing the Fermi energy, in agreement with experiment. We can determine $\Delta\epsilon$ directly from the measured phase shift between the local SdH maxima and the corresponding nonlocal magnetoresistance peaks. For an asymmetric probe configuration, the fraction of the edge current which reaches a voltage *probe* depends on the length of the current path to this probe along the edges. Therefore this fraction depends on the direction of current circulation [13]. This dependence is responsible for the observed asymmetry with respect to the inversion of the magnetic field.

The solid lines in Fig. 3 are the best fit to the temperature dependence of Eq. (1) using $\Delta\epsilon = 4$ meV for the upper curve and $\Delta\epsilon = 3.2$ meV for the lower curve, calculated as discussed above. The only adjustable parameter is the overall amplitude. The fitted values for the amplitudes of both peaks shown in Fig. 3 scale with the amplitude of the corresponding SdH maxima. Given the approximate derivation of Eq. (1) the agreement with experiment is very good.

To summarize, we have observed a new nonlocal magnetoresistance effect in the intermediate regime between quantum edge conduction and classical bulk transport. Unlike previous examples of nonlocal resistance, the observed effect vanishes as $T \rightarrow 0$ K and is due to dissipation in the bulk rather than in the electrical contacts. However, quantum conduction in the edge states plays an important role in carrying the current to classically inaccessible regions. Using a simple model we are able to explain the observed temperature dependence of the effect and account for the phase of the nonlocal oscillations relative to the local SdH oscillations.

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