Spin-dependent resistivity at transitions between integer quantum Hall states

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The longitudinal resistivity at transitions between integer quantum Hall states in two-dimensional electrons confined to AlAs quantum wells is found to depend on the spin orientation of the partially-filled Landau level in which the Fermi energy resides. The resistivity can be enhanced by an order of magnitude as the spin orientation of this energy level is aligned with the majority spin. We discuss possible causes and suggest a new explanation for spike-like features observed at the edges of quantum Hall minima near Landau level crossings.

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The integer quantum Hall (QH) effect is a relatively well-understood phenomenon in condensed matter physics that occurs when a two-dimensional electron system (2DES) is subjected to a large perpendicular magnetic field $\mathbf{B}$. Associated with the effect are oscillations in the longitudinal magnetoresistivity (MR), $\rho_{xx}$, including wide regions of zero resistivity, that occur as the Fermi level passes through successive disorder-broadened Landau levels (LLs) containing bands of localized and extended states. At transitions between integer QH states, there are peaks in $\rho_{xx}$ that have been extensively studied in the context of universal scaling phenomena. The magnetic field resolves the LLs into up and down spin branches split by the Zeeman energy, causing the $\rho_{xx}$ maxima to split as well, but the spin degree of freedom is generally thought to play a minor role in determining the value of the $\rho_{xx}$ maxima in the spin-resolved QH regime.

We observe unexpectedly large changes of $\rho_{xx}$ at transitions between integer QH states as the 2DES is tilted with respect to an applied magnetic field such that the LL spins are polarized (Fig. 1). Specifically, the value of $\rho_{xx}$ at a transition is higher when the spin of the energy level in which the Fermi energy ($E_F$) resides is aligned with the majority, lowest LL spin, and lower for those levels with the opposite spin orientation. The resistivity difference can be as large as an order of magnitude.

We performed measurements on electrons confined to a narrow (45Å wide) AlAs quantum well (QW) with Al$_0$Ga$_0$As barriers, grown on a GaAs (100) substrate. The 2DES in this system occupies a single, out-of-plane valley with an isotropic Fermi contour, a band effective mass of 0.21$m_0$, and a band g-factor of 2. We also made measurements on 2D electrons in a strained, wide (110Å wide) AlAs QW, where a single, in-plane valley with an anisotropic Fermi contour with transverse and longitudinal band effective masses of 1.1$m_0$ and 0.21$m_0$ is occupied. We studied a total of four samples from three different wafers, all of which were lithographically patterned in a Hall bar configuration with Ohmic contacts made by depositing AuGeNi and alloying in a reducing environment. Metallic front and back gates were deposited for in situ control of the charge density, $n$. The peak mobilities are 5.0 m$^2$/Vs in the narrow QW samples and 17 m$^2$/Vs for the low-mass direction of the wide QW. We performed measurements down to $T = 20$ mK using low-frequency lock-in techniques. The samples were mounted on a single-axis tilting stage to allow the angle, $\Theta$, between the normal to the sample and the magnetic field to be varied from 0° to 90° in situ. We define $B_{tot}$ as the total magnetic field, and $B_{\parallel}$ and $B_{\perp}$ as the components parallel or perpendicular to the 2DES plane.

Because the Zeeman energy, $E_Z$, depends on $B_{tot}$ while the cyclotron energy, $E_c$, depends on $B_{\perp}$, the ratio $E_Z/E_c$ can be adjusted by changing $\theta$. At the so-called coincidence angles, $E_Z/E_c$ takes integer values as different spin branches of the various LLs are brought into energetic coincidence. These level crossings can be labelled by a crossing index given by the difference between the LL indices for the crossing up and down spin levels, $i = \mathcal{N}_i^+ - \mathcal{N}_i^-$. By tilting the sample through coincidence angles, the spin for a LL associated with a given $\rho_{xx}$ maximum can be reversed [see Fig. 2(a)]. This argument is based on a single-particle picture and, as discussed later, many-body effects can alter this simple behavior near the coincidence angles.

AlAs QWs are very well suited to tilted-field measurements. Their relatively large value of $g^*m^* \quad (g^* \quad \text{and} \quad m^* \quad \text{are the renormalized Landé g-factor and electron mass respectively})$ allows LL coincidences to be reached at modest tilt angles as compared, for example, to GaAs 2DESs which have an order of magnitude smaller $g^*m^*$. In narrow AlAs QWs, $g^*m^*$ is small enough to allow the $i = 1$ coincidence to be observed, while wide AlAs QWs are past this coincidence at $\theta = 0^\circ$ for most accessible densities. Because of their relatively narrow well widths, AlAs QWs have the further advantage of being largely free of finite-thickness effects associated with the coupling of $B_{\parallel}$ to the orbital degree of freedom.

In Figs. 1(a) and (b), we show $\rho_{xx}$ in a narrow QW at several different $\Theta$. The value of $g^*m^*$ is about 1.5 at $n = 2.1 \times 10^{11}$ cm$^{-2}$, and previous measurements have indicated no strong dependence of $g^*m^*$ on magnetic field or filling factor ($\nu$) in this system. This means that $E_Z < E_c$ at $\theta = 0^\circ$, so that the $\rho_{xx}$ maximum near $B_{\perp}$
FIG. 1: (color online) $\rho_{xx}$ vs. $B_\perp$ at $T = 300$ mK in a narrow AlAs QW as the sample is tilted through the first ($i = 1$) coincidence angle ($\theta \approx 45^\circ$). (a) A resistance spike is visible near the $\nu = 2$ minimum, and (b) subsequently merges with the $\rho_{xx}$ maximum between $\nu = 2$ and 1, leading to a dramatic enhancement of $\rho_{xx}$ at that maximum. The two lowest energy levels, labelled by their LL index and spin, and the location of $E_F$ are shown schematically for the indicated $\rho_{xx}$ maxima.

$E = 6$ T corresponds to $E_F$ passing through the $0 \uparrow$ energy level (transition between $\nu = 2$ and 1 QH states) [see Fig. 1(b)]. As the sample is tilted to angles near the $i = 1$ coincidence, a spike develops in the $\nu = 2$ minimum [Fig. 1(a)]. Similar spikes were observed previously in wide AlAs QWs [9] and were associated with scattering at boundaries between different spin domains concomitant with a first-order, many-body spin reversal [3]. The properties of the spikes observed in narrow QWs are similar, including hysteric behavior and the movement of the spike to higher $B_\perp$ as $\theta$ is increased. In contrast to previous studies, however, when the spike eventually merges with the $\rho_{xx}$ maximum between $\nu = 2$ and 1 [Fig. 1(b)], the spike’s height and width increase continuously until it is no longer distinct from the maximum itself, and we are left with a single, enhanced $\rho_{xx}$ maximum. Even after the spike is unresolvable, the $\rho_{xx}$ maximum continues to increase in height as the sample is tilted until a saturation angle is reached. The apparent position of the QH minimum also depends on the spin orientation of the adjacent LLs. It is evident in Fig. 1 that $\rho_{xx}$ at $\nu \gtrsim 2$ attains finite values at lower $B_\perp$ when the second lowest LL is $1 \downarrow$ than when it is $0 \uparrow$. This causes an apparent shift in the position of the $\nu = 2$ minimum to lower $B_\perp$ [9]. The Hall resistance plateaus (not shown) are truncated in step with the $\rho_{xx}$ minima.

To further understand the enhancement, we plot $\rho_{xx}$ at half-integer $\nu$ in Fig. 2 as a function of $1/\cos(\theta)$ [(b) and (c)] and $B_{tot}$ [(d) and (e)] for two different $n$ in a narrow QW. We also show schematically the evolution of the energy levels with increasing $\theta$ for the narrow QW in Fig. 2(a), although we have omitted peculiarities of the level configurations associated with many-body effects near the coincidence angles. Prior to the $i = 1$ coincidence, $\rho_{xx}$ at $\nu = 3/2$ and $7/2$ is locally minimal as a function of $1/\cos(\theta)$ while at $\nu = 5/2$ it is maximal. From the energy level diagram in Fig. 2(a), we see that $3/2$ and $7/2$ correspond to $E_F$ passing through spin-up
branches while 5/2 corresponds to a down branch. As the \( i = 1 \) coincidence is traversed, the spin for the energy levels corresponding to these \( \nu \) flip, and accordingly \( \rho_{xx} \) increases or decreases as the spin flips to down or up respectively. With further increase of \( \theta \), this correspondence between \( \rho_{xx} \) and spin continues, and eventually \( \rho_{xx} \) at a given \( \nu \) saturates at the angle beyond which there are no further spin reversals. Beyond this saturation, there sometimes appears a slow decrease in \( \rho_{xx} \) (e.g. \( \nu = 3/2 \) at \( n = 1.3 \times 10^{11} \text{ cm}^{-2} \)), however this is small compared to the initial increase and does not appear to depend on the spins of the nearby (higher) LLs (i.e., it is unaffected by passage through subsequent coincidence angles).

Plotting \( \rho_{xx} \) at the various half-integer \( \nu \) as a function of \( B_{\text{tot}} \) [Figs. 2(d,e)] reveals that the enhancement saturates at a roughly \( \nu \)-independent \( B_{\text{tot}} \) which is close to the saturation field, \( B_F \), for the parallel-field MR where \( \theta = 90^\circ \) and \( B_{\text{tot}} = B_0 \). \( B_F \) has been shown experimentally [10, 11] and theoretically [12] to correspond to the total spin polarization of the charges. Thus, the correspondence between the saturation fields, which is consistent with previous observations of isotropy and field-independence of \( g^* m^* \) in narrow AlAs QWs [5], corroborates the assertion that the enhancement of the \( \rho_{xx} \) maxima are associated with spin polarization. The analogy with parallel-field MR implies a possible explanation for the occurrence of spin-dependent \( \rho_{xx} \) in the QH regime.

In the absence of orbital effects, the parallel-field MR can be attributed largely to the spin-polarization dependence of screening in a 2DES [12]. This is essentially a consequence of the Pauli exclusion principle, whereby like-spin charges screen disorder from each other less effectively than opposite spin charges. This might seem irrelevant to the QH case where the density of states is quantized into discrete (albeit disorder-broadened), spin-resolved, and macroscopically degenerate LLs. However, screening associated with inter-LL excitations can indeed play a significant role in this regime [13]. In particular, considering that the relatively large band mass in AlAs causes the electron-electron interaction energy, \( E_x \sim e^2/4\pi\epsilon l_B \) (\( l_B \) is the magnetic length), to be greater than \( E_c \) for all accessible fields, it becomes evident that mixing between different LLs must be taken into account. Disorder may then be screened most effectively when the spin of the LL in which \( E_F \) resides is different than the majority spin of the low-lying LLs.

Treating different spin levels as parallel conduction channels could also point to a cause of spin-dependent \( \rho_{xx} \). The spin-channel Hall resistivities \( (\rho_{xy}^s, s \in \{\uparrow, \downarrow\}) \) are different by virtue of the different filling factors for the two spin channels. Thus, if the spin-channel longitudinal resistivities \( (\rho_{xx}^s) \) are close, then the spin-channel longitudinal conductivities, \( \sigma_{xx}^s = \rho_{xx}^s/(\rho_{xx}^s + \rho_{xy}^s) \), are different and, when added and inverted, give a spin dependent total \( \rho_{xx} \). Of course, \( \rho_{xx}^\uparrow \) and \( \rho_{xx}^\downarrow \) are not necessarily close, and if the spin-channel conductivities are calculated directly instead of the resistivities, as is usually the case, then some additional mechanism must be postulated to account for the spin-dependence. For \( \rho_{xx}^s = \rho_{xx}^\uparrow \ll \rho_{xx}^\downarrow \), this model gives a factor of 9 change in \( \rho_{xx} \) at \( \nu = 3/2 \) when tilting past the \( i = 1 \) coincidence, which is in reasonable agreement with the narrow QW results. As we show below, however, measurements in wide AlAs QWs give results that are not accounted for by this simple model for spin-dependent resistivity.

Our measurements of strained, wide AlAs QWs with a single anisotropic valley occupied yield give results that are similar to those in the narrow QWs but with some notable differences. Figure 3 displays MR data for current (I) oriented (a) along and (b) transverse to the long-axis of the occupied valley, as indicated schematically in each panel. The inset of (a) shows a closeup of the \( \theta = 0^\circ \) trace, and the LL indices and spins for the lowest three energy levels are shown in (b) for the indicated traces.
spikes form near the coincidence angles (the $i = 2$ coincidence at $n = 4.5 \times 10^{11}$ cm$^{-2}$ in this system occurs at $\theta \approx 46^\circ$) and large $\rho_{xx}$ changes occur when these angles are passed [Figs. 3(a,b)]. In contrast to the narrow QWs, however, the enhancement at $\nu = 5/2$ in the wide QW is anisotropic, with a remarkable factor of $\sim 19$ enhancement in (a) but only $\sim 5$ in (b). By straining the sample to depopulate one valley and populate the other and repeating the measurement, we have confirmed that the enhancement anisotropy arises from the valley (mass) anisotropy and not from the orientation of $B_{||}$ with respect to $I$ or the valley. The narrow wells, which have an isotropic in-plane Fermi contour, show no enhancement anisotropy. Such an anisotropic enhancement is not explained by the aforementioned two-channel model for spin-dependent $\rho_{xx}$. We also note that the strong spin-dependence is suppressed when the sample is strained to occupy a second valley. This might be expected in the context of an inter-LL screening picture, since the additional valley degree of freedom would allow like-spin charges to screen disorder from each other as effectively as opposite-spin. That spin-dependent $\rho_{xx}$ is strongest in single-valley systems with large electron effective mass may explain why this effect has not been reported for Si-MOSFETs (metal-oxide-semiconductor field-effect transistors) or GaAs heterostructures.

Whatever the cause, it is clear that $\rho_{xx}$ at QH transitions is spin-dependent. This observation prompts a simple explanation for $\rho_{xx}$ spikes that exist on the flanks of QH states that is independent of domain-wall scattering. It has been shown theoretically that, when opposite spin LLs are brought near energetic coincidence, $B_{\perp}$ can induce a discontinuous transition where the nearly coincident LLs suddenly trade places in energy. For example, at low $B_{\perp}$ the energy levels may have the configuration of the middle panel of Fig. 2(a), but as $B_{\perp}$ is increased enough to begin depopulating the low lying LLs a rapid transition to the configuration of the left panel of Fig. 2(a) can occur. Spikes in $\rho_{xx}$ have been shown to accompany such transitions [10], and have been explained as arising from scattering at boundaries between domains of the two competing energy level configurations [5]. In the case of a spin-dependent $\rho_{xx}$, however, such spikes would accompany a collective spin-reversal transition at the edges of QH states even in the absence of domain-wall scattering. As an example, for an appropriate $\theta$, the system may have the configuration of the middle panel of Fig. 2(a) as $E_F$ enters the extended states of the $1_\perp$ energy level, and because the spin of that LL is down, $\rho_{xx}$ would be large. As $B_{\perp}$ is increased and a transition to the configuration in the left panel of Fig. 2(a) is made, the partially occupied LL would change to $0_\parallel$ and $\rho_{xx}$ would rapidly drop to the un-enhanced values. This would register in $\rho_{xx}$ as a spike-like feature.

The spikes occurring at the edges of QH states are observed to follow the same rising curve defined by the enhanced (spin-down) $\rho_{xx}$ maxima (Figs. 1 and 3), which is consistent with them originating from spin-dependent $\rho_{xx}$. Therefore, the relative sizes of the spikes as they pass across the landscape of $\rho_{xx}$ are neatly accounted for, as is the fact that the largest spikes are observed when the spin-dependent enhancement is strongest [e.g. Fig. 3(a)]. It is important to point out, however, that resistivity spikes are also observed deep within the QH minima (e.g. $\theta = 44.8^\circ$ trace in Fig. 1(a)), i.e. in the localized states, that do not fall within the envelopes defined by the enhanced $\rho_{xx}$ maxima. Domain-wall scattering may account for these, but we reiterate that large spike-like features at the edges of QH states are an inevitable consequence of a spin-dependent $\rho_{xx}$ with a sudden spin-reversal transition. Spikes will occur even for second-order (continuous) transitions as long as the transition is sufficiently rapid in magnetic field on the scale of the QH features. In fact, the sharpness of the transitions appears to diminish as they move deeper into the extended states (i.e. deeper into the flanks of the QH minima) [Fig. 1(b)], possibly signaling a crossover from a first to a second-order transition. Finally, for spikes arising from spin-dependent $\rho_{xx}$, the spike maximum may not correspond to the center of the phase transition as was previously thought to be the case, nor does the full-width-at-half-maximum give the width of the transition. These facts place restrictions on deriving information about the exchange energy or Curie temperature associated with the ferromagnetic transition for spikes occurring at the edges of QH states.

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A spin-dependent $\rho_{xx}$ is evident, though not highlighted, in [11]. Due to the large 2DES width, there is likely also an orbital component to the tilted-field MR enhancement in that measurement.

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