Magnetic-Field-Driven Destruction of Quantum Hall States in a Double Quantum Well

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The magnetotransport of double-quantum-well structures reveals a high-magnetic-field regime in which quantum Hall states at odd-integer ν are missing, where ν is the number of filled energy levels. In these double wells, the ν=odd quantum Hall states originate from the symmetric-antisymmetric energy gap. With increasing barrier thickness between the two wells, initially the magnetic field destroys only the $\nu=1$ state, then both the $\nu=1$ and $\nu=3$ states disappear. We attribute their destruction to the onset of interwell electron-density fluctuations.

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The integral[1,2] and fractional[2-5] quantum Hall effects were both discovered in high-mobility two-dimensional electron systems. Fabrication of multiple 2D electron layers in close proximity allows the controlled introduction of additional degrees of freedom associated with the third dimension. The double quantum well (DQW) is the simplest of these structures and preserves both high electron mobility and external gating of the electron density in each layer. Interlayer Coulomb interactions in a multilayer structure are predicted to lead to even denominator fractional quantum Hall states[6-8] and increased stability of a new collective state, perhaps the Wigner crystal.[9,10] Experimentally, both the integral[11] and fractional[12] quantum Hall effects have been observed in multilayer systems with substantial interlayer tunneling. In this Letter, we instead report the *absence* of selected integral quantum Hall states in these structures in a high-magnetic-field regime.

We present magnetotransport results from three DQW samples consisting of nominally identical quantum wells separated by barriers of different widths. The samples are grown by molecular-beam epitaxy and consist of two GaAs wells of width $d_W=139$ Å, separated by a $Al_{0.3}Ga_{0.7}As$ barrier of thickness $d_B=28$, 40, or 51 Å (see inset of Fig. 1). The electrons are provided by remote delta-doped donor layers ($N_{Si}=6\times10^{11}$ cm^{-2}) set back from each side of the DQW by ~600-Å-thick $Al_{0.3}Ga_{0.7}As$ spacer layers. The as-grown electron density in each well is determined from the two distinct peaks in the Fourier transform of the longitudinal resistivity (ρ_{xx}) oscillations versus inverse magnetic field. Biasing of a back-side gate (~0.5 mm from the double well) can vary the ratio of the two densities, indicating that there are two spatially separated layers of highly mobile electrons in our samples. When the wells are greatly out of balance, the density difference between the layers varies linearly with applied gate bias, as expected for two decoupled single quantum wells. However, as gate bias brings the single-well electron states into resonance, tunneling begins to couple the single-well states. Finally, when the electron densities between the two wells are balanced, the single-well electron states form symmetric and antisymmetric DQW states (inset of Fig. 1) separated by an energy gap, $\Delta_{SAS}=E_{AS}-E_S$, where E_S and E_{AS} are the symmetric state and antisymmetric state energies. In two of our balanced samples, there remains a resolvable pair of oscillation frequencies in ρ_{xx} vs (tesla)$^{-1}$, resulting from the higher electron density in the symmetric state, n_S, than the antisymmetric state, n_{AS}. This finite, minimum density difference between two electron states is clear experimental evidence of a symmetric-antisymmetric (SAS) gap in our DQW samples. All data presented in this paper are taken from samples whose electron densities are balanced between the quantum wells. (We note that a fixed gate potential can lead to a magnetic-field-driven density imbalance resulting from different gate capacitance to each single well. This effect, however, is vanishingly small for our 0.5-mm-distant back-side gate.) The total electron densities, n_{2D} ($n_{2D}/2$ per well), and mobilities, μ, at $T=0.3$ K are given in Table I, along with the calculated Δ_{SAS}^{calc}s, resulting from a self-consistent solution of the DQW Schrödinger equation and Poisson's equation.

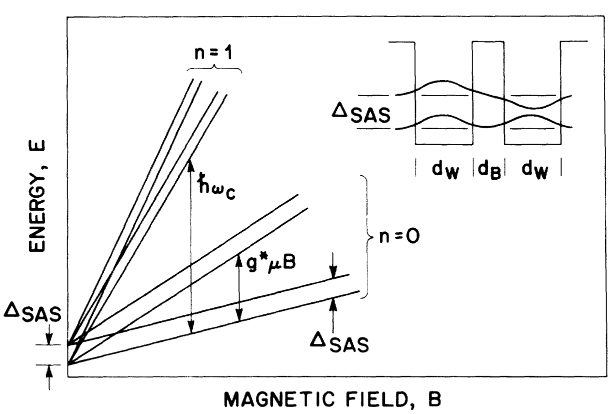

FIG. 1. Energy diagram for a double quantum well. The cyclotron ($\hbar\omega_c$), Zeeman ($g^*\mu_B B_{tot}$), and symmetric-antisymmetric (Δ_{SAS}) energies are indicated. Inset: The symmetric and antisymmetric states of the double quantum well.

TABLE I. Sample parameters.

d_B (Å)	$10^{-11}n_{2D}$ (cm^{-2})	μ ($cm^2/V\,s$)	Δ_{SAS}^{calc} (K)	Missing states
28	4.2	740000	17.3	...
40	3.8	630000	8.1	$\nu=1$
51	3.9	350000	3.9	$\nu=1,3$

With the application of a magnetic field perpendicular to the 2D planes, the symmetric and antisymmetric states each form a fan of spin-split Landau levels, as shown in Fig. 1 for the $n=0$ and $n=1$ Landau levels. The energy-level filling factor is defined as the number of occupied energy levels, $\nu=n_{2D}h/eB_\perp$ (where $B_\perp$ is the component of magnetic field perpendicular to the 2D layers). As in a single quantum well, each minimum in ρ_{xx} and the accompanying plateau in the Hall resistivity, ρ_{xy}, correspond to a magnetic field at which the Fermi energy, E_F, lies in an energy gap. The filling factor, ν, is determined experimentally from the quantization of the Hall plateaus, $\rho_{xy}=h/\nu e^2 \sim 25.8\ k\Omega/\nu$.

Experimental results at low magnetic fields demonstrate that the samples behave in accordance with the single-electron energy diagram of Fig. 1. (1) From the measured n_S and n_{AS} in the balanced double wells, an experimental Δ_{SAS} is determined: $n_S/n_{AS}=(E_F-E_S)/(E_F-E_{AS})$. In the $d_B=28$ Å sample, Δ_{SAS} is measured to be 17.5 K (± 0.5 K) in excellent agreement with the calculated value in Table I. In the $d_B=40$ Å sample, the measured Δ_{SAS} is 5.5 K ($^{+2.0}_{-1.0}$ K). (2) At $B\sim 0.5$ T, there is a doubling of the ρ_{xx} oscillation frequencies, followed by another doubling at $B\sim 0.9$ T. This implies the existence of three distinct energy gaps in the DQW samples which become resolved at different magnetic fields. The three relevant energies are labeled in Fig. 1: the cyclotron energy, $\hbar\omega_c=\hbar eB_\perp/m^*$; the Zeeman energy, $g^*\mu_B B_{tot}$; and the symmetric-antisymmetric energy Δ_{SAS}. (3) Finally, as the data of Fig. 2 demonstrate, the assignment of each quantum Hall state to one of the energy gaps of Fig. 1 is made unambiguous by tilting the samples in a magnetic field. When an in-plane magnetic field is applied, the quantum Hall states at $\nu=4n+2$ ($n=0,1,2,\ldots$) are found to *increase* in strength (deeper ρ_{xx} minima and broader ρ_{xy} plateaus). Since an in-plane magnetic field increases the Zeeman energy (proportional to B_{tot}) at a given filling factor (fixed $B_\perp$), we ascribe the states at $\nu=4n+2$ to the Zeeman energy gap. By contrast, each quantum Hall state at $\nu=$odd is *destroyed* by a $\sim$1.5-T in-plane magnetic field (see Fig. 2). This is compelling experimental evidence that these states result from the Δ_{SAS} energy gap, which would collapse from decreased tunneling between quantum wells in the presence of an in-plane magnetic field.[13] The body of data summarized above demonstrates that each of the three samples realizes a double quantum well with interwell tunneling in accordance with the single-electron picture of Fig. 1. These data also indicate that the observed $\nu=$odd quantum Hall states correspond to the Fermi energy lying in the SAS energy gap.

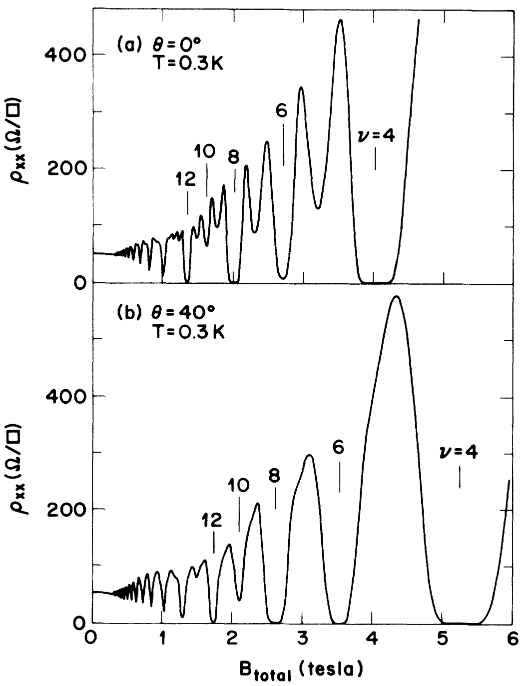

FIG. 2. Longitudinal resistivity, ρ_{xx}, as magnetic field is tilted from $\theta=0°$ (normal to the 2D planes). The minima at $\nu=6$, 10, and 14 become deeper, while all minima at $\nu=$odd are destroyed.

Figure 3 contains the most striking result of our experiments: the low-temperature ρ_{xx} and ρ_{xy} data from the $d_B=51$ Å sample in a perpendicular magnetic field to $B=23$ T. Note that there are well-formed plateaus at all integers, $4<\nu<16$ (see the 3× enlarged data in the inset). Strong fractional quantum Hall states at $\nu=\frac{2}{3}$, $\frac{4}{3}$, and $\frac{8}{3}$ also exist. Despite the fact that this high-mobility sample resolves all three single-electron energies ($\hbar\omega_c$, $g^*\mu B$, and Δ_{SAS}) and, additionally, the collective excitation gap of the fractional quantum Hall effect, *there is no evidence of a quantum Hall state at either $\nu=3$ or $\nu=1$.* Particularly striking is the well-developed $\nu=\frac{8}{3}$ fractionally quantized state and the missing $\nu=3$ quantized state at nearly identical magnetic fields.

The other two samples show different behavior at $\nu=1$ and $\nu=3$. In the $d_B=40$ Å sample, the $\nu=3$ state is visible; however, there is no evidence of the $\nu=1$ state. Finally, in the $d_B=28$ Å sample, all the quantum Hall states at integer $\nu<16$ are well developed. The missing quantum Hall states are summarized in Table I. Our data demonstrate that high magnetic fields perpendicular to the quantum wells can destroy the Δ_{SAS} energy gap and the associated $\nu=$odd quantum Hall states. Furthermore, the mechanism for the destruction of the quantum Hall states depends strongly on the thickness of the interwell barrier.

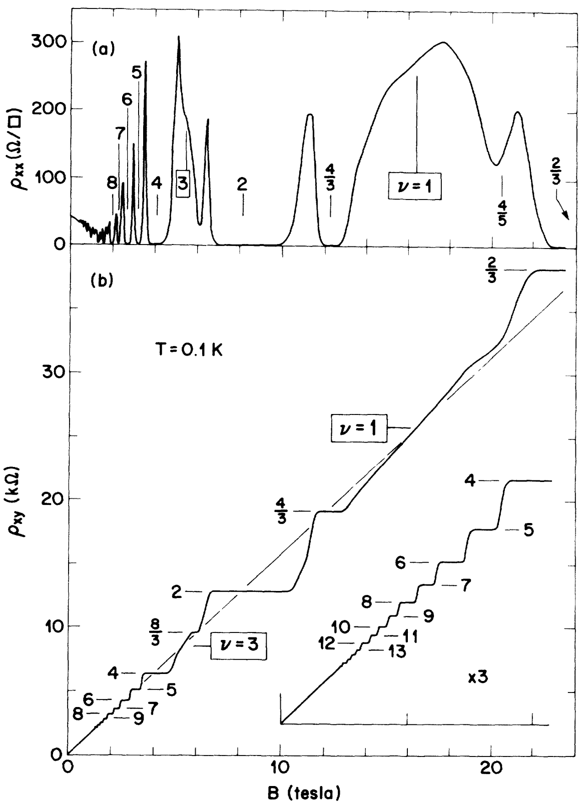

FIG. 3. Longitudinal resistivity, ρ_{xx}, and Hall resistivity, ρ_{xy}, for the $d_B = 51$ Å sample in a perpendicular magnetic field. There is no evidence of quantum Hall states at $\nu = 1$ and $\nu = 3$. Inset: An expansion of the low-field ρ_{xy} data. For clarity, the ρ_{xx} data at $4 < B < 8$ T have been reduced by a factor of 2; at $B > 8$ T by a factor of 10.

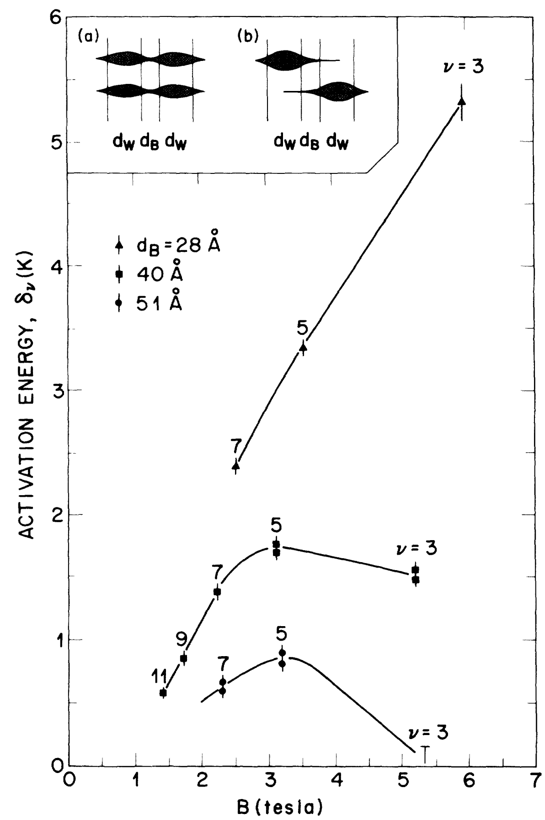

FIG. 4. Experimental activation energies at ν=odd for three different barrier thicknesses. The lines are guides to the eye. Insets: The charge density of nearest-neighbor electrons: (a) occupying the symmetric double-well state and (b) confined to opposite quantum wells (see discussion in text).

The fractional quantum Hall states in Fig. 3 offer further evidence that the Δ_{SAS} gap has collapsed in the high-magnetic-field regime. Normally, if the SAS energy gap were intact, the observation of the $\nu = \frac{2}{3}$, $\frac{4}{3}$, and $\frac{8}{3}$ fractional states would imply that the $\nu = \frac{5}{3}$ and $\frac{7}{3}$ states should also exist. Instead, there is no evidence of fractional quantum Hall states at $\nu = \frac{5}{3}$ and $\frac{7}{3}$ in Fig. 3. If the Δ_{SAS} gap is destroyed in high magnetic fields, the DQW becomes energetically identical (in a single-electron picture) to a system of two independent single quantum wells with no interwell tunneling. In this case, the filling factor ν, as defined in this Letter, would represent filling factor $\nu/2$ in each single quantum well. It follows that the well-developed quantized states at $\nu = \frac{2}{3}$, $\frac{4}{3}$, and $\frac{8}{3}$ in Fig. 3 would correspond to the well-known $\frac{1}{3}$, $\frac{2}{3}$, and $\frac{4}{3}$ states existing independently in each well. Also, the lack of fractional quantum Hall states at $\nu = \frac{5}{3}$ and $\frac{7}{3}$ in the DQW results from the absence of fractional states at filling factors $\frac{5}{6}$ and $\frac{7}{6}$ in the individual wells.

Finally, the temperature dependence of ρ_{xx} at the ν=odd minima indicates that ρ_{xx} is thermally activated over more than an order-of-magnitude dynamic range: $\rho_{xx} = \rho_0 \exp(-\delta_\nu/T)$, where δ_ν is the activation energy. Figure 4 contains activation energies for various ν=odd minima from all three samples. The data suggest that δ_ν depends roughly linearly on magnetic field until this dependence weakens and collapses to zero as the ν=odd quantum Hall states are destroyed in high magnetic fields. Note that the calculated activation energies, $\Delta_{SAS}^{calc}/2$, are more than twice the largest measured activation energies, δ_ν, for each of the three samples. Also, in the single-electron model of Fig. 1, Δ_{SAS} is expected to be magnetic-field *independent*. The reduced magnitudes and striking magnetic-field dependence of the activation energies in Fig. 4 suggest that band mixing due to many-electron interactions and/or disorder becomes important in the DQW system in the quantum limit.

Before speculating further on the physics underlying the above observations on the DQW system, we present two similar systems which exhibit missing quantum Hall states for different physical reasons: (1) In a heterojunction with two occupied subbands, missing quantum Hall states have been observed when the Fermi energy lies at a degeneracy of two energy levels.[14] This effect, however, cannot explain the states destroyed in the DQW samples: (i) For example, in the $d_B = 51$ Å sample, for the $\Delta_{SAS} = g^* \mu B$ degeneracy to occur at $B = 16.5$ T (and thus destroy the $\nu = 1$ state), the enhanced g^* factor at $\nu = 1$ must be $\lesssim 0.34$ *smaller* than the bulk GaAs value. Similarly, setting $\Delta_{SAS} = \hbar\omega_c$ at $B = 5.5$ T (to destroy the $\nu = 3$ state) implies m^* more than 13 times its GaAs value. (ii) Also, the in-plane magnetic-field data of Fig. 2 identify the $\nu = 4n+2$ states in our samples with the Zeeman gap and the $\nu =$odd states with the SAS energy gap. As Fig. 1 makes clear, the observed $n =$odd states correspond to the SAS energy gap only if the $\Delta_{SAS} = g^* \mu B$ degeneracy occurs at some magnetic field smaller than our observed $\nu = 15$ state at $B \sim 1$ T.

(2) In an idealized two-layer model (layers of zero thickness, no disorder), quantum Hall states at $\nu =$odd have been predicted to exist in the absence of interlayer tunneling.[6-8] In this model, interwell Coulomb interactions open an excitation gap with a $B^{1/2}$ magnetic-field dependence and a magnitude similar to the activation energies of Fig. 4. Also, high magnetic fields are predicted to destroy these $\nu =$odd states. However, this idealized model may not adequately describe the experimental situation, since there is demonstrated to be substantial tunneling in our DQW samples: (i) The model $\nu =$odd states are rapidly destroyed in the presence of tunneling[8,15] and also (ii) an in-plane magnetic field would leave the $\nu =$odd model states essentially unchanged,[15] in contrast to the data of Fig. 2.

We therefore suggest an alternate mechanism for the existence of a high-magnetic-field regime in which the Δ_{SAS} energy gap has collapsed. In the single-electron model of Fig. 1, each electron is shared equally between the two quantum wells. This situation is represented schematically in inset (a) of Fig. 4, where the charge densities of nearest-neighbor electrons are depicted. At $\nu =$odd, the lowest-energy excitation consists of exciting an electron from a symmetric state to an antisymmetric state. Through these excitations, local charge-density fluctuations can be established between the two quantum wells. An extreme case of these fluctuations is represented in inset (b) of Fig. 4, in which each electron is confined entirely within a single quantum well and nearest-neighbor electrons are forced into opposite wells. The Δ_{SAS} energy cost of the excitations will be at least partially offset by the gain in electron correlation energy, which becomes much more important at the lower filling factors in high magnetic fields. As such, the Δ_{SAS} energy gap at $\nu =$odd will soften and could eventually collapse in sufficiently strong magnetic fields. This would account for the observations of the missing integral and fractional quantum Hall states in a double quantum well.

In conclusion, we reveal a high-magnetic-field regime in double-quantum-well structures which is evidenced by missing quantum Hall states at energy-level filling factors $\nu =$odd. For samples of increasing barrier thickness between the two wells, first the $\nu = 1$ state is destroyed, then both the $\nu = 1$ and $\nu = 3$ states are destroyed. This magnetic-field-driven regime depends critically on the additional degree of freedom in the third dimension provided by the double-quantum-well structure.

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