Many-Body Integer Quantum Hall Effect: Evidence for New Phase Transitions

S. Q. Murphy, J. P. Eisenstein, G. S. Boebinger, L. N. Pfeiffer, and K. W. West
AT&T Bell Laboratories, Murray Hill, New Jersey 07974
(Received 25 October 1993)

The $\nu = 1$ quantum Hall effect in bilayer 2D electron systems is shown to continuously evolve from a regime dominated by single-particle tunneling into one where interlayer Coulomb interactions stabilize the state. This many-body integer quantum Hall state exhibits a phase transition to a compressible state at large layer separation. We also find evidence for an intriguing and unexpected second transition to a new incompressible state, driven by an in-plane magnetic field. While the origin of this last result is unclear, we discuss a recent model of a pseudospin textural phase transition.

PACS numbers: 73.40.Hm, 73.20.Dx

In a single layer two-dimensional electron system (2DES), the quantum Hall effect (QHE) arises from gaps in the electronic density of states. For the integer QHE these gaps are generally of single-particle origin, while those in the fractional QHE result from many-body effects. In contrast, in a double layer 2DES both single-particle and many-body regimes can be explored *at a single filling factor* by the tuning of appropriate sample parameters. Associated with a double layer system are two new energy scales, the single-particle tunneling gap (Δ_{SAS}) separating the lowest symmetric and antisymmetric states in the bilayer and the many-body Coulomb interaction between electrons in different layers. Each can lead to a QHE; in concert they can either enhance or destroy a QHE.

For two identical widely separated layers connected in parallel, the QHE spectrum is equivalent to that of each individual layer except that the total filling factor (ν) associated with each Hall plateau is twice the corresponding single layer value. Thus, for example, no $\nu = 1$ QHE would be found since no $\nu = 1/2$ QHE has ever been observed in a single layer 2DES. However, even the smallest amount of interlayer tunneling can lead to *odd integer* QHE states since it opens the symmetric-antisymmetric gap. The ground state at $\nu = 1$ would then be a single fully filled Landau level of (spin polarized) symmetric state electrons and the excitation gap would be Δ_{SAS}.

Even in the absence of tunneling, bilayer QHE states with no single layer analog can occur if the interlayer Coulomb interactions are sufficiently strong [1–3]. The recent observation of a QHE at $\nu = 1/2$ (i.e., 1/4 filling in each layer) in double layer 2D systems [4,5] is a good example. This new state, closely resembling the original Laughlin 1/3 state, is a purely collective phenomenon. Similarly, a bilayer many-body QHE has been predicted to exist at $\nu = 1$ in the absence of tunneling [2,3]. This collective state cannot exist for arbitrarily weak interlayer coupling, unlike the aforementioned tunneling $\nu = 1$ QHE. Instead, it is expected [3,6,7] to suffer a quantum phase transition to a compressible state at some critical layer spacing. Intriguingly, this QHE state is thought to possess a broken symmetry, one not shared by the bilayer $\nu = 1/2$ state, which has been predicted [8,9] to lead to such diverse phenomena as neutral gapless modes, a Kosterlitz-Thouless phase transition, and even Josephson and Meissner effects.

With two distinct physical mechanisms capable of generating a $\nu = 1$ QHE in double layer systems, a rich phase diagram can be anticipated. We present here an experimental determination of that phase diagram showing an apparent continuum of incompressible QHE states between the two limiting regimes outlined above. In the weak tunneling limit, we observe the predicted phase transition from a Laughlin-like incompressible liquid to a compressible phase when the interlayer separation exceeds a critical value. In addition, we find evidence for an unexpected, *additional* phase transition between distinct incompressible $\nu = 1$ states, driven by the application of a magnetic field parallel to the layers.

The modulation-doped double quantum well (DQW) samples used in this study were grown by molecular beam epitaxy. Eleven of the fifteen samples consist of two 180 Å wide GaAs wells separated by a $d_b = 31$ Å $Al_xGa_{1-x}As$ undoped barrier. (The remainder have slightly different well and barrier widths.) By changing the Al concentration ($0.3 < x < 1.0$) in the barrier, the symmetric-antisymmetric gap [10] was varied from a maximum of $\Delta_{SAS} = 8.5$ K down to less than 1 K. Appropriately placed Si doping layers above and below the DQW produced total carrier concentrations ranging from $N_{tot} \sim 0.8 \times 10^{11}$ to 3.2×10^{11} cm^{-2}, equally split between the quantum wells, with low temperature mobilities of 0.5×10^6 to 1.5×10^6 cm^2/V s.

Figure 1 shows the extreme sensitivity of the $\nu = 1$ QHE to sample parameters. These two samples differ only in barrier thickness ($d_b = 31$ vs 40 Å) and (slightly) in carrier concentration ($N_{tot} = 1.26 \times 10^{11}$ vs 1.45×10^{11} cm^{-2}). Although both display clear QHE's at $\nu = 2$, 2/3, and even 4/5, they differ dramatically at $\nu = 1$. The narrower barrier sample exhibits a strong $\nu = 1$ QHE, while the other exhibits no such state. Since the barrier width affects both the tunneling strength and the *interlayer* Coulomb correlations, while the carrier concentration controls the *intralayer* dynamics, it is not obvious *a priori* why the $\nu = 1$ QHE is present in one sample and absent in the other. To determine this, and,

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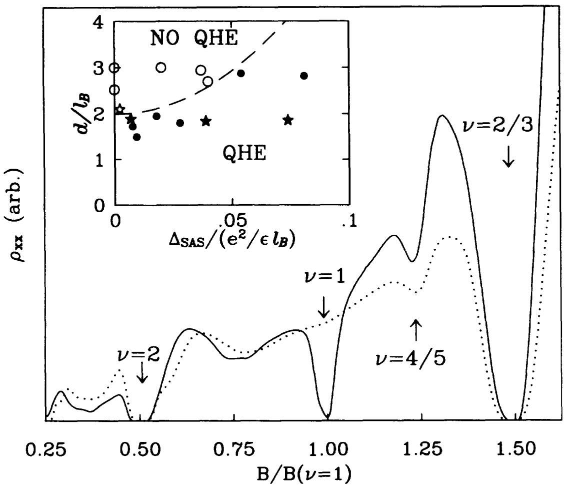

FIG. 1. The resistivity, ρ_{xx}, at 0.3 K vs magnetic field (normalized by the field at $\nu = 1$) for two similar samples. In one ($d_b = 30$ Å, $N_{tot} = 1.26 \times 10^{11}$ cm^{-2}) a strong QHE at $\nu = 1$ is found, while in the other ($d_b = 40$ Å, $N_{tot} = 1.45 \times 10^{11}$ cm^{-2}) the $\nu = 1$ QHE is absent. Inset: The phase diagram at $\nu = 1$. The solid symbols represent samples that show a $\nu = 1$ QHE, open symbols denote those that do not. (The two ρ_{xx} traces are taken from the samples represented as the leftmost open and solid stars.)

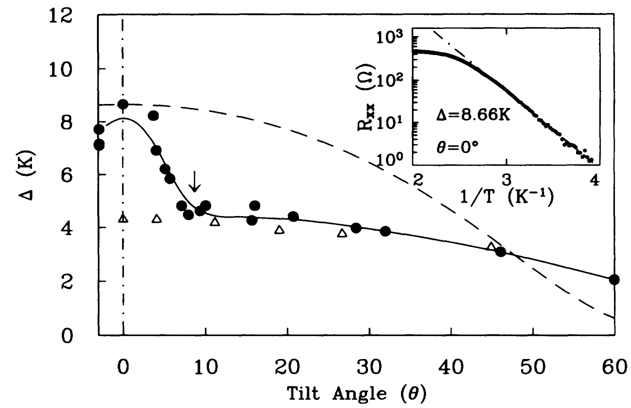

FIG. 2. The energy gap, Δ, as a function of tilt in a weakly tunneling sample ($\Delta_{SAS} = 0.8$ K). The solid circles are for $\nu = 1$, open triangles for $\nu = 2/3$. The arrow indicates θ_c. The dashed line is the predicted behavior of Δ_{SAS} in a tilted magnetic field arbitrarily normalized to the measured Δ at zero angle. Inset: A typical Arrhenius plot at $\nu = 1$.

indeed, to elucidate the distinction between the tunneling and many-body regimes at $\nu = 1$, we consider the phase diagram shown in the inset to Fig. 1. The horizontal axis reflects the tunneling strength ($\propto \Delta_{SAS}$), while the vertical axis is inversely proportional to the interlayer Coulomb energy $e^2/\epsilon d$ (with d the quantum well center-to-center distance). Both axes are normalized by the intralayer Coulomb energy $e^2/\epsilon l_B$ [with the magnetic length $l_B = (\hbar/eB)^{1/2}$ evaluated at $\nu = 1$]. Solid symbols indicate samples exhibiting a clear $\nu = 1$ QHE, while open symbols denote those showing no sign of such a QHE (down to $T = 0.3$ K).

A number of important qualitative conclusions can be drawn from the data in the phase diagram. Most obvious is the existence of a clear boundary separating those samples showing a $\nu = 1$ QHE from those that do not, with the dashed line estimating the position of that boundary. It is apparent that the QHE is destroyed on increasing the separation between the layers even if the tunneling strength is held fixed. In the strong tunneling limit this phase transition, first observed by Boebinger *et al.* [11], reflects the energetic advantage pairs of electrons obtain by abandoning their symmetric DQW wave functions for states localized in opposite wells. While this costs tunneling energy, it saves on Coulomb repulsion. If d/l_B is large enough, this level mixing effectively destroys the tunneling gap and with it the QHE [12]. Such a competition between tunneling and interactions is insufficient, however, to explain the present finding that the phase boundary *intercepts the vertical axis at a nonzero value of* d/l_B (≈ 2). This observation is compelling evidence that a $\nu = 1$ QHE exists in the limit of zero tunneling as predicted theoretically [2,3]. Thus, for the two samples initially discussed (the leftmost stars, open and solid, Fig. 1), both lie sufficiently close to the vertical axis of the phase diagram, where the phase line is essentially horizontal, so that their different tunneling gaps can be ruled out as the source of the destruction of the $\nu = 1$ state. These samples reflect instead the predicted quantum phase transition [3,6,7] from an incompressible Laughlin-like liquid to a compressible state when the critical layer separation is exceeded. Finally, the distribution of samples plotted in the phase diagram strongly suggests that the $\nu = 1$ QHE evolves continuously from a state dominated by single-particle tunneling to one where many-body effects are paramount, with no compressible region apparent in between [13].

We turn now to the behavior of the $\nu = 1$ QHE in tilted magnetic fields where evidence for a novel, additional and unexpected phase transition has been found. These studies were begun in order to exploit the *in situ* suppression of Δ_{SAS} produced by an added in-plane magnetic field $B_\parallel$, to discriminate between tunneling and many-body effects. This suppression, a simple single-particle matrix element effect [14], exhibits a Gaussian dependence: $\Delta_{SAS}(\theta)$ [$\propto \exp(-\alpha^2 \tan^2\theta)$, where $\alpha = d/2l_B$ and $\tan\theta = B_\parallel/B_\perp$]. While in the strongly tunneling regime the $\nu = 1$ QHE energy gap (Δ) should be dominated by Δ_{SAS}, in the many-body limit we expected Δ to be independent of Δ_{SAS} and therefore of angle.

Figure 2 shows the angular dependence of the energy gap at $\nu = 1$ for the weakly tunneling sample ($\Delta_{SAS} \sim 0.8$ K, $N_{tot} = 1.26 \times 10^{11}$ cm^{-2}) indicated by the leftmost solid star in the phase diagram in Fig. 1. Δ was determined from the thermally activated behavior (shown in the figure inset) of the resistivity $\rho_{xx} \sim \exp(-\Delta/2T)$ at $\nu = 1$. Both the large magnitude gap [$\Delta(\theta = 0) \sim 8.7$

K$\gg \Delta_{SAS}$] and the low temperature required ($T < 0.4$ K) to enter the activated regime point to strong many-body effects. It is the angular dependence, however, which is the most striking aspect of the data in Fig. 2. The gap is found to drop sharply for small angles around $\theta = 0$ before crossing over into a roughly angle-independent regime beyond $\theta_c \approx 8 \pm 2°$. Essentially identical angular behavior has been found in two other samples which, while sharing similarly small levels of tunneling, have slightly lower carrier concentrations ($N_{tot} = 0.8 \times 10^{11}$ and 1.06×10^{11} cm^{-2} [15]). This anomalous tilted field behavior is inconsistent with both the angular independence anticipated for a many-body $\nu = 1$ state and the Gaussian suppression (dashed line in the figure) expected if the state is dominated by single-particle tunneling. We emphasize that neither the $\nu = 2/3$ QHE (open triangles) nor the $\nu = 1/2$ state (a distinctly bilayer many-body effect) show any significant angular dependence.

Changes in the energy gap Δ can reflect changes in the spectrum of charged excitations and/or the 2DES ground state itself. Among the possible excitations which might exhibit an angular dependence are those involving spin flips and those in which the tunneling gap $\Delta_{SAS}(\theta)$ is bridged. Both of these can be ruled out by the abrupt angular dependence of Δ. For spin flips the expected angular dependence arises from the Zeeman energy $g\mu_B B_{tot}$, but this has changed by only 15 mK upon tilting from $\theta = 0$ to θ_c, compared to the observed change in Δ of > 4 K. Similarly, the change in Δ_{SAS} (from the Gaussian) is only 20 mK. More importantly, the observed angular dependence is qualitatively inconsistent with a simple level crossing. The excitation dominant at small θ should remain so for $\theta > \theta_c$ as it obviously extrapolates to energy gaps smaller than those observed. We suggest instead that the data reflect a change in the liquid ground state at $\nu = 1$.

To examine this effect further we studied additional samples which have the same quantum well and barrier thicknesses and very similar electron densities, yet have larger tunneling gaps ($\Delta_{SAS} = 4.4$ and 8.5 K) by virtue of different alloy concentrations in the barrier. Nevertheless, these samples, denoted by the center and rightmost solid stars in the phase diagram, exhibit qualitatively similar angular dependences, shown in Fig. 3, to that seen with the weakly tunneling sample. Again, a transition from a regime of strong to weak angular dependence is observed. Importantly, however, $\tan\theta_c$ is found to increase roughly linearly with Δ_{SAS}. This dependence is shown in the inset to the figure where $\tan\theta_c$ is plotted against the (zero angle) $\Delta_{SAS}/(e^2/\epsilon l_B)$. The dependence of the θ_c on Δ_{SAS}, while further reducing the likelihood that the electron spin is involved, does establish a fundamental role for tunneling in the transition. In contrast, the angular suppression of Δ_{SAS}, while perhaps relevant, is apparently not central to the effect since, as already discussed, in the weakly tunneling sample Δ_{SAS} has changed by only 2% when the transition has occurred. The data suggest instead a competition between two ground states, one of which, at $\theta < \theta_c$, takes advantage of tunneling by forming a many-body condensate out of symmetric state electrons. The competing state apparently ignores tunneling, finding, in the presence of the in-plane magnetic field, a more favorable many-body configuration. The correlation advantage of this mixed state must exceed $\Delta_{SAS}/2$ (per electron) for $\theta > \theta_c$. How the in-plane field $B_\parallel$ drives this transition is not evident from the data. Clearly, however, $B_\parallel$ does more than produce a simple scalar reduction of Δ_{SAS} and likely represents an additional dimension in the phase diagram at $\nu = 1$.

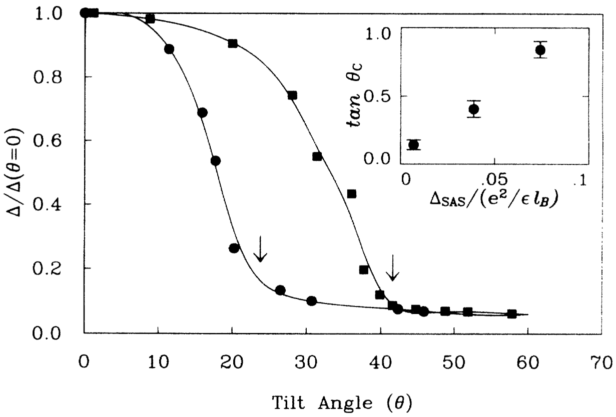

FIG. 3. Angular dependence of the energy gap for two more strongly tunneling samples (solid circles, $\Delta_{SAS} = 4.4$ K; solid boxes, 8.5 K). The arrows indicate the critical angle. Both data sets have been normalized by the measured gap at $\theta = 0$ [$\Delta(0) = 5$ and 14.6 K, respectively]. Inset: Dependence of $\tan\theta_c$ on $\Delta_{SAS}/(e^2/\epsilon l_B)$.

Another clue to the origin of this unusual transition may lie in its absence at filling factors $\nu = 1/2$, 2/3, and 2. This may suggest a connection to the broken symmetry [8,9] which exists at $\nu = 1$, but not at these other fillings. Neglecting Coulomb interactions, the ground state at $\nu = 1$ in a tunneling DQW is a full Landau level of symmetric state electrons. Using a pseudospin quantum number ($s_z = \pm 1/2$) to denote the layer index (top, bottom), this tunneling $\nu = 1$ state corresponds to a *ferromagnetic* alignment of the total pseudospin ($\mathbf{S}$) along $\hat{\mathbf{x}}$, since a symmetric state is just an $s_x = +1/2$ eigenstate. Even in the absence of tunneling, Coulomb interactions (i.e., exchange effects) keep the pseudospins parallel, leaving the ground state at $\nu = 1$ a fully polarized incompressible state, at least in the idealized limit $d/l_B = 0$ [3,16]. While the pseudomagnetization can point in any direction, the smallest amount of tunneling will break the symmetry and orient the polarization along $\hat{\mathbf{x}}$. This representation is not applicable at $\nu = 1/2$, 2/3, or 2; the $\nu = 1/2$ state is believed to not be an eigenstate of $\mathbf{S}$ [3], while $\nu = 2/3$ and $\nu = 2$ are $\mathbf{S} = 0$ singlets [17] in the $d/l_B = \Delta_{\rm SAS} = 0$ limit. Importantly, this ferromagnetic model for $\nu = 1$ can be extended to nonzero (but small) d/l_B. While the interlayer capacitive charging energy forces the pseudospin to lie near the x-y plane, the broken symmetry remains [8,9].

Yang *et al.* [18] have proposed an explanation, based on the ferromagnetic properties of the $\nu = 1$ state, for the tilted field transition reported here. They note that a parallel magnetic field not only uniformly reduces the magnitude, but also periodically shifts the *phase* of the tunneling matrix elements. The wavelength of this shift ($\lambda = \Phi_0/B_{\|}d$) is the distance required to thread one flux quantum (Φ_0) between the layers. For zero applied $B_{\|}$, the pseudospin texture is uniformly polarized along $\hat{\mathbf{x}}$. At nonzero $B_{\|}$, however, the local pseudospin must track the shifting phase in order to satisfy the criteria for tunneling. The resulting twisted pseudospin texture maintains the energetic advantage of tunneling, but since neighboring pseudospins are no longer parallel, this distortion costs exchange energy. As $B_{\|}$ increases λ becomes shorter and this energetic penalty increases. At a critical parallel field ($B_{\|,c}$), Yang *et al.* predict that the system abandons tunneling and makes a transition to a new uniformly polarized state. The calculated $B_{\|,c}$ (or tilt angle, $\tan\theta_c = B_{\|,c}/B_{\perp}$) is proportional to $\sqrt{\Delta_{\rm SAS}}$, at least in the small $\Delta_{\rm SAS}$ limit. While their numerical estimate of θ_c for the weakly tunneling sample of Fig. 2 is only about a factor of 2.5 larger than the measured valued, the predicted $\sqrt{\Delta_{\rm SAS}}$ appears inconsistent with the observations shown in the inset to Fig. 3. This discrepancy may suggest the sample parameters (close to the compressible phase boundary and extending to large $\Delta_{\rm SAS}$) reach outside the predictive range of the theory.

In an alternate approach, Ezawa and Iwazaki [19] have suggested that our results may reflect the predicted $\nu = 1$ Josephson behavior. Although controversial [18], it would be interesting to test for this with direct measurements of the tunneling I-V characteristics.

To summarize, we have presented a number of new results on the interplay of single-particle tunneling and many-body effects in the QHE at $\nu = 1$ in double layer two-dimensional electron systems. In addition to showing that a continuum of incompressible states exists, for small enough layer separation, between the strong and weak tunneling limits, we have found an unexpected and intriguing new phase transition induced by a component of magnetic field parallel to the layers. This may reflect a textural phase transition in the pseudospin field.

The authors acknowledge fruitful discussions with Allan MacDonald, Steve Girvin, Song He, and Francis Ezawa. J.P.E. and G.S.B. acknowledge the hospitality of the Aspen Center for Physics.

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