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Evidence for spontaneous interlayer phase coherence in a bilayer quantum Hall exciton condensate

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Abstract

Recent experimental work on the quantized Hall state at total filling factor $\nu_T = 1$ in bilayer 2D electron systems has revealed a number of striking phenomena, including a giant and sharply resonant enhancement of the interlayer tunneling conductance at zero bias. The tunneling enhancement is a compelling indicator of spontaneous interlayer phase coherence among the electrons in the system. Such phase coherence is perhaps the single most important attribute of the excitonic Bose condensate which describes this remarkable quantum Hall state.

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Keywords: D. Interlayer; D. Tunneling; D. Conductance

1. Introduction

Double layer two-dimensional electron gases in large magnetic fields have been the subject of intense inquiry for more than a decade [1,2]. The reason for this persistent interest stems from the existence of novel correlated electron fluids in the bilayer system which possess no analogs in single layer systems. Early work [3,4] on this problem focussed on the fractional quantized Hall effect (FQHE) state in which the total Landau level filling fraction of the bilayer system is $\nu_T = 1/2$. This state, which occurs when the density of the individual layers places them at $\nu_1 = \nu_2 = 1/4$, was considered especially exciting because of its flagrant violation of the so-called 'odd-denominator rule' which governs all FQHE states of a single layer 2D electron gas in the lowest Landau level.[1] Remarkably, however, a bilayer FQHE state at $\nu_T = 1/2$ had been predicted [5] in advance of the experiments, and a strong candidate ground state wavefunction was available [6]. Indeed, the discovery of the $\nu_T = 1/2$ FQHE may be viewed as a triumph of the theory, pioneered by Laughlin, of electronic correlations in 2D systems at high magnetic fields.

Less heralded at the time was the observation [4] of a QHE state in the same bilayer systems at $\nu_T = 1$. Although this state again has no counterpart in single layers (the individual layer filling fractions being $\nu_1 = \nu_2 = 1/2$), it can exist in the absence of Coulomb interaction effects. If the tunneling between the layers is sufficiently strong, the individual layer eigenstates of the double well potential hybridize into symmetric and antisymmetric double well states. At $\nu_T = 1$ the lower-lying symmetric state is fully filled with electrons and the tunneling-induced splitting Δ_{SAS} between the symmetric and antisymmetric states provides the energy gap required for a quantized Hall effect.

There is, however, a more interesting way to realize a $\nu_T = 1$ bilayer QHE state. It is by now well established, in theory [1] and experiment [2], that a such a state can occur even in the absence of interlayer tunneling. When the layers are close enough together interlayer Coulomb interactions are of comparable importance to intralayer ones and electrons in either layer are thus forced to avoid their

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[1] There are, of course, even denominator quantized Hall states (at $\nu = 5/2$ and $\nu = 7/2$) in the first excited Landau level of single 2D layer systems.

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neighbors in both layers. This fact leads to a QHE at $\nu_T = 1$ with an energy gap to charged excitations which depends solely on Coulomb interactions. In this sense the $\nu_T = 1$ state should be regarded as a fractional QHE state. Fig. 1 displays both the $\nu_T = 1/2$ and $\nu_T = 1$ QHE states in a typical double layer 2D electron gas sample.

There are several equivalent ways to view the strongly correlated bilayer $\nu_T = 1$ QHE phase, including as an easy-plane pseudo-ferromagnet or a Bose condensate of interlayer excitons. Essential to all such descriptions is the concept of spontaneous interlayer phase coherence [1]: Each electron in the system resides in a superposition of individual layer eigenstates. Using a pseudospin language to encode the layer degree of freedom, these superpositions are written as $|\uparrow\rangle + e^{i\varphi}|\downarrow\rangle$, where $|\uparrow\rangle$ denotes an electron in the 'top' layer while $|\downarrow\rangle$ represents an electron in the bottom layer. This superposition state corresponds to the pseudospin lying in the x–y plane, inclined by an angle φ to the x-axis. In the ground state at small layer separation the phase φ is arbitrary, but is the same for all electrons. The system is thus fully pseudospin polarized. The phase φ is a macroscopic quantum phase, exactly analogous to the phase of the order parameter in superfluid helium or conventional superconductors. We stress that this symmetry breaking occurs spontaneously owing to interaction effects; there is no need for any interlayer tunneling. This is quite remarkable, for it implies that in the ground state there is complete quantum uncertainty as to which layer an electron is in, even if there is no tunneling. If tunneling is present, the symmetry breaking is explicit and the phase φ is near zero in the ground state. In this case, each electron is in the usual symmetric linear combination state.

It is possible to re-cast the pseudo-ferromagnet picture of the $\nu_T = 1$ bilayer QHE state into the language of exciton

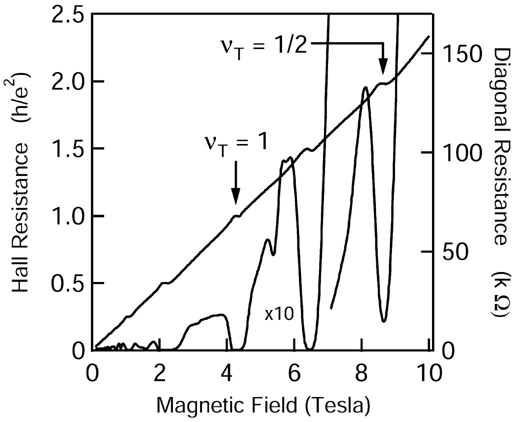

Fig. 1. Hall and diagonal resistances in a high mobility double quantum well showing the $\nu_T = 1$ and $\nu_T = 1/2$ quantized Hall states. Sample consists of two 18 nm GaAs quantum wells separated by a 3 nm pure AlAs barrier layer. After Ref. [4].

condensation [7]. The wavefunction[2] for the ferromagnetic state can be approximated as

$$\Psi = \prod_k \frac{1}{\sqrt{2}}\left[c^\dagger_{k,T} + e^{i\varphi}c^\dagger_{k,B}\right]|0\rangle$$

The operators $c^\dagger_{k,T}$ and $c^\dagger_{k,B}$ create electrons in the kth guiding center state of the lowest Landau level in the top and bottom layers, respectively. In this equation the vacuum state $|0\rangle$ has no electrons in either layer. If, however, one starts with a vacuum $|0'\rangle$ consisting of no electrons in the bottom layer and no holes in the top layer (i.e. the top layer contains a fully filled Landau level), then the above wavefunction can be re-written as

$$\Psi = \prod_k \frac{1}{\sqrt{2}}\left[1 + e^{i\phi}c^\dagger_{k,B}c_{k,T}\right]|0'\rangle$$

The combination $c^\dagger_{k,B}c_{k,T}$ creates an electron–hole pair, the electron being in the bottom layer and the hole in the top layer. Written in this way the $\nu_T = 1$ many-body wavefunction is the same as the BCS wavefunction for a conventional superconductor[3], only the relevant Cooper pairs are interlayer excitons.

The finite separation between the layers in a bilayer 2D electron system lowers the symmetry of the Coulomb interaction. Electrons in the same layer always experience stronger repulsion than electrons in opposite layers. The finite layer separation has two important consequences. First, the pseudospin moment remains close to the x–y plane. Excursions out of the plane put more charge in one layer than the other, and this costs capacitive energy. This was implicitly assumed in the wavefunctions given above. Second, the total pseudospin is in fact not really a good quantum number. As a result, the total pseudospin moment fluctuates, primarily in the x–y plane. At small layer separation these fluctuations are not so severe as to destroy the QHE phase, but at some critical separation there is a quantum phase transition to a compressible non-QHE phase. This makes for an important, and enriching, distinction between the bilayer $\nu_T = 1$ QHE state and the $\nu = 1$ QHE occuring in single layer systems. In the latter case the Coulomb interaction retains complete spin invariance and no such phase transition exists.

The strongly correlated bilayer $\nu_T = 1$ QHE state possesses a number of very remarkable properties, many of which do not occur in any other experimentally studied QHE state [1]. The charged excitations of the system, which determine the QHE energy gap, are vortex-like structures in the pseudospin field known as merons. These excitations can have positive or negative vorticity and carry half-integer charge, $\pm e/2$. In the absence of tunneling, these vortices are logarithmically bound in pairs of charge 0 or $\pm e$ at low

[2] We assume that the true spins are polarized by the Zeeman coupling.

[3] Here the BCS factors u_k and v_k have equal modulus: $|u_k| = |v_k| = 1/\sqrt{2}$. See Ref. [7].

temperatures, but suffer a Kosterlitz-Thouless unbinding transition at some finite temperature T_{KT}, where the QHE is expected to break down. Not surprisingly, long wavelength neutral collective modes are also anticipated, these being essentially pseudospin waves. Owing to the spontaneously broken U(1) symmetry, these modes go soft at $q = 0$ and the dispersion at small q is expected to be linear [8,9].

The most remarkable aspects of the $\nu_T = 1$ bilayer QHE state are Josephson-like interlayer tunneling and counter-flow superfluidity [9–12]. The possibility of these phenomena occurring at $\nu_T = 1$ was first suggested by Wen and Zee [9], but the ideas connect back to earlier work on multi-valley systems by Rasolt and collaborators [13,14]. A major anomaly in the tunneling conductance between the layers would strongly support the notion of spontaneous interlayer phase coherence. If, as expected, electrons are coherently spread through the barrier separating the two layers, then it is plausible to assume there would be very little resistance to tunneling. On the other hand, counter-flow superfluidity results from the invariance of the ground state energy to global changes in the phase variable φ. A uniform spatial gradient in φ corresponds to equal, but opposite, electrical currents in the two layers. This is equivalent to a uniform flow of interlayer excitons in one direction. Such a phase gradient obviously produces a finite kinetic energy. However, just as in superfluid helium and conventional superconductors, the U(1) (or easy-plane) symmetry of the problem prevents the relaxation of supercurrents without generation of phase slips. Of course, being a two-dimensional system, superfluidity in the $\nu_T = 1$ QHE state is expected only in linear response.

2. Separate contacts

It is clear that Josephson-like tunneling and counter-flow superfluidity cannot be directly accessed via conventional resistivity measurements in which the two 2D layers are electrically connected in parallel. To study these exotic phenomena, separate electrical contacts to the individual 2D layers are essential. In 1990 we developed a scheme, based on localized selective depletion of one or the other 2D layer, which allows the establishment of such contacts [15]. This technique has proven enormously effective for studying the physics of the individual layers in double layer systems in which interlayer correlations are negliglible. Important results on the compressibility [16], tunneling density of states [17], and the intra- and interlayer electron scattering rates [18,19], at both zero and high magnetic fields, have been obtained.

Fig. 2 illustrates the application of the separate contact technique to measurements of the interlayer tunneling conductance dI/dV in weakly-coupled bilayer 2D electron systems [17]. The figure shows two tunneling traces, one at zero magnetic field and one at $B = 13$ T, where for the sample in question, the individual 2D layers are near filling factor $\nu = 1/2$. At $B = 0$ a sharp resonance is observed. This resonance reflects the alignment of the ground state energy levels in the individual quantum wells. Unlike the case with 3D electrodes, in 2D the conservation of energy and in-plane momentum produces a sharply resonant tunneling dI/dV vs. V characteristic. The width of the resonance, about 0.2 meV for the data shown, reflects the quantum lifetime of the single electron states in the two 2D gases [19]. For the data shown, the resonance is centered at zero interlayer voltage. This is because the two layers have the same 2D density. If the bilayer is unbalanced, the $B = 0$ tunnel resonance occurs at a finite interlayer voltage set by the density difference.

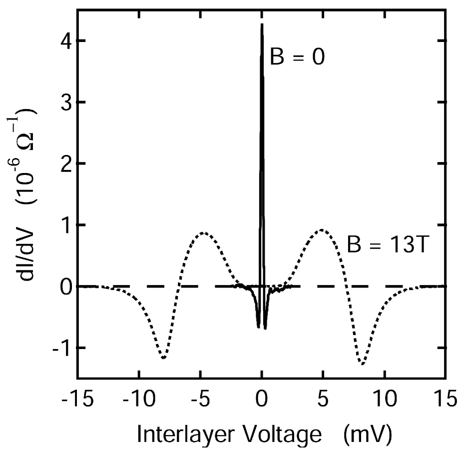

Fig. 2. Tunneling spectra at in a widely-spaced double layer 2D electron system. Narrow resonance at zero magnetic field contrasts sharply with high field data in which a strong suppression of the zero bias tunneling conductance is evident. After Ref. [17].

The tunneling conductance at $B = 13$ T is vastly different than that at $B = 0$. Most importantly, there is a broad region of suppressed conductance around $V = 0$. This is a correlation effect, albeit a single layer one [20,21]. Electrons attempting to tunnel into the 2D gas at high fields effectively have to force their way into an 'interstitial' location in the strongly correlated electron gas created by the high magnetic field. This implies tunneling will only occur at voltages of order the mean Coulomb energy $e^2/\epsilon\ell$ in the system, $\ell = (\hbar/eB)^{1/2}$ being the magnetic length. Even if low energy $N + 1$ particle states of the 2D gas exist (i.e. even if the 2D gas is not in a QHE state, as is the case for the data in the figure), these states cannot be accessed on the very short time-scale associated with tunneling. As a result, the tunneling spectrum shows a Coulomb pseudogap around zero bias. Unlike the resonance seen at $B = 0$, this suppression is pinned to the Fermi level, and thus does not move to finite voltage if the densities in the two layers are made unequal. In different language, when an electron attempts to tunnel between widely-separated 2D electron gases, it is wholly 'unaware' of the correlations present in the layer it is about to enter. As a result, only highly excited

$N+1$ particle states can be accessed. The tunneling conductance near $V=0$ is therefore heavily suppressed. Tunneling appears only at high voltages and is then widely dispersed in energy.

Until recently it had not proven possible to extend the separate contact technique to the regime of strong interlayer correlations. The reason for this is simple: In order to have strong interlayer correlations, the barrier separating the two quantum wells must be narrow. On the other hand, narrow barriers imply stronger interlayer tunneling. If the tunneling resistance is too small, the very notion of separate contacts loses its meaning. Since the tunneling is exponentially sensitive to the barrier thickness, while interlayer correlation effects are much less sensitive, it is not obvious that the desired regime can be reached.

Fortunately, the versatility of molecular beam epitaxy allows one to adjust not only the thickness of the barrier layer, but also its height. A pure AlAs barrier between GaAs quantum wells presents a barrier approximately 1 eV high[4], in contrast to the ~250 meV offered by the customary $Al_{0.3}Ga_{0.7}As$ alloy. Second, the importance of interlayer Coulomb correlations is governed by the ratio of the physical separation d between the 2D layers (which is usually taken to be the center-to-center distance between the GaAs quantum wells) to the mean inter-electron separation in each 2D gas. Thus, for the same physical dimensions, lower density systems will be more strongly correlated than higher density ones. Since the mean inter-electron separation, at fixed filling factor ν, is proportional to the magnetic length ℓ, the ratio d/ℓ provides a convenient dimensionless measure of the effective layer spacing. It is through the combination of nearly pure AlAs barriers and low 2D densities that we have succeeded in reaching the goal of separately contacted bilayer 2D systems possessing both strong interlayer correlations at $\nu_T=1$ and yet very weak interlayer tunneling.

Most of our samples consist of two 18 nm GaAs quantum wells separated by a 10 nm $Al_{0.9}Ga_{0.1}As$ barrier[5]. The cladding layers are $Al_{0.3}Ga_{0.7}As$ and contain Si dopant sheets approximately 220 nm above and below the double well structure. In their as-grown state these samples have a density of about 5.5×10^{10} cm^{-2} per layer and a low temperature mobility of about 1×10^{6} cm^2/V. In order to reach the regime of the $\nu_T=1$ bilayer QHE, the layer densities are symmetrically reduced by about a factor of two using top and back-side gates. We estimate the tunnel splitting to be about $\Delta_{SAS}\approx8$ neV, or about 1 part per million of the mean Coulomb energy ($e^2/\epsilon\ell$) in the system at $\nu_T=1$. With such small tunneling strengths the QHE at $\nu_T=1$ is overwhelmingly dominated by Coulomb interaction effects. Transport measurements (see Fig. 3) of the longitudinal resistivity ρ_{xx} reveal a deep and thermally activated minimum at $\nu_T=1$ from which a quasiparticle charge gap of about $\Delta\approx0.35$ K has been extracted. This is nearly 4000 times the estimated single-particle tunnel splitting, Δ_{SAS}.

3. Tunneling at $\nu_T = 1$

With separate contacts to strongly correlated bilayer samples in hand, the tunneling conductance at $\nu_T=1$ can be examined. Fig. 4 reveals our earliest results on how the tunnel spectrum evolves as the effective layer separation d/ℓ at $\nu_T=1$ in the sample is reduced [22]. At high density (and thus large d/ℓ) the tunnel spectrum exhibits the familiar suppression of dI/dV around zero bias due to intra-layer correlations. As the layer densities are symmetrically reduced, the pseudogap region contracts, but otherwise changes little. At a certain density, however, a small, but sharp, peak appears at $V=0$. For the sample shown, this transition occurs near $d/\ell=1.83$. As the density is reduced

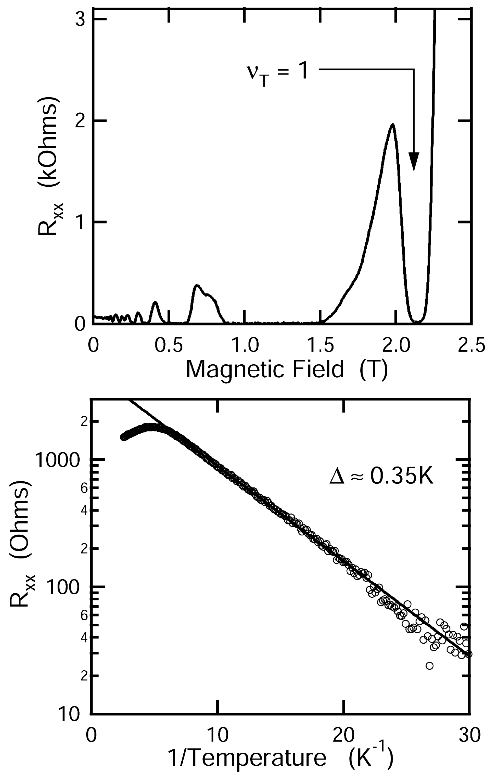

Fig. 3. Top panel: Longitudinal resistance of low density double layer 2D electron system at $T=25$ mK. Strong minimum at $B=2.15$ T reflects $\nu_T=1$ QHE, at $d/\ell=1.6$. Bottom panel: Temperature dependence of resistance at $\nu_T=1$. Slope implies a charge gap of $\Delta\approx0.35$ T.

[4] This is, of course, true only for Γ-point electrons. It turns out, however, that Γ–X mixing is negligible in the situations of interest here, and so this large barrier is in fact the relevant one.

[5] In fact, we use $Al_{0.9}Ga_{0.1}As$ barriers. We find the slight reduction in Al concentration improves the yield of our ohmic contacts without significantly increasing the tunneling strength.

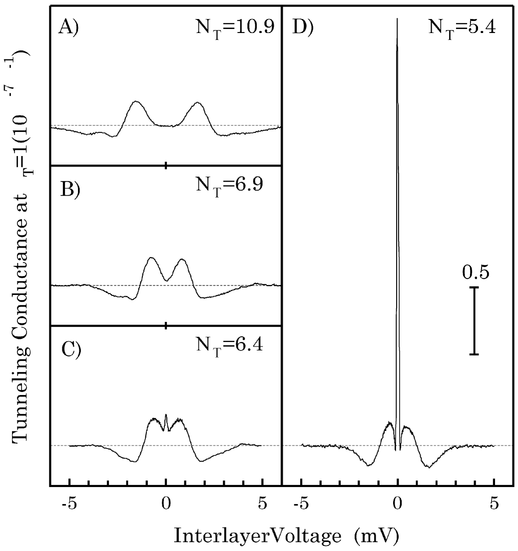

Fig. 4. Discovery of spontaneous interlayer phase coherence in a bilayer quantum Hall system. Sequence of tunneling spectra at $\nu_T = 1$ and $T = 40$ mK for various total densities in the double layer system. High density data show familiar suppression of zero bias tunneling conductance. Low density data reveals a huge and sharply resonant enhancement of the zero bias conductance. After Ref. [22].

further, this peak rapidly grows and eventually dwarfs all other features in the tunnel spectrum.

Improvements to the tunneling measurement have shown that the zero bias peak is much taller and narrower than the data in Fig. 4 suggest. Fig. 5 shows a more recent measurement, taken at $d/\ell = 1.50$. The full-width at half maximum of the conductance peak is only about $\Gamma \approx 2$ μeV at $T = 25$ mK. This is extremely narrow by typical standards for electronic quantum lifetimes in semiconductors. For comparison, the width of the single-particle tunneling resonance at zero magnetic field is some forty times times larger than Γ.

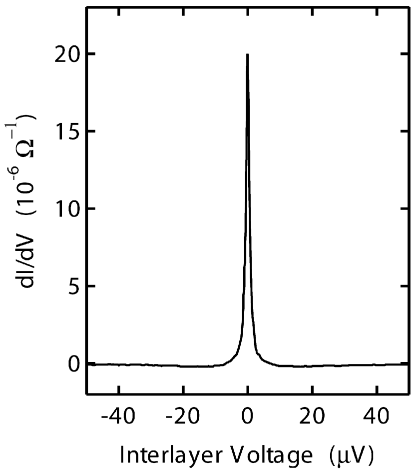

Fig. 5. Tunnel resonance at $\nu_T = 1$ and $T = 25$ mK and $d/\ell = 1.5$. Full width at half maximum of this resonance is less than 2 μV.

The giant enhancement of the tunneling conductance shown in Figs. 3 and 4 occurs only over the narrow magnetic field range in which the $\nu_T = 1$ QHE is observed [22]. Furthermore, as Fig. 6 shows, the peak height falls and the linewidth Γ increases as the temperature rises. We find that the strength of the tunneling anomaly increases smoothly as d/ℓ is decreased below about 1.83. Fig. 7 illustrates this via the measured current–voltage ($I-V$) characteristics of the tunnel junction. The figure displays a family of $I-V$ curves, taken at $\nu_T = 1$ and $T = 25$ mK, for effective layer separations $1.9 \geq d/\ell \geq 1.6$. The curves toward the back, which were taken at large d/ℓ, exhibit the suppression of tunneling near $V = 0$ characteristic of weakly coupled 2D layers. Around $d/\ell = 1.8$ a step-like feature appears at $V = 0$. This feature grows rapidly, but apparently continuously, as d/ℓ is reduced further. Obviously, these near discontinuities in the tunnel current reflect the same phenomenon as the sharp peak in the dI/dV tunneling conductance data. The similarity of the data in Fig. 7 to the dc Josephson effect in superconducting tunnel junctions is striking, and the

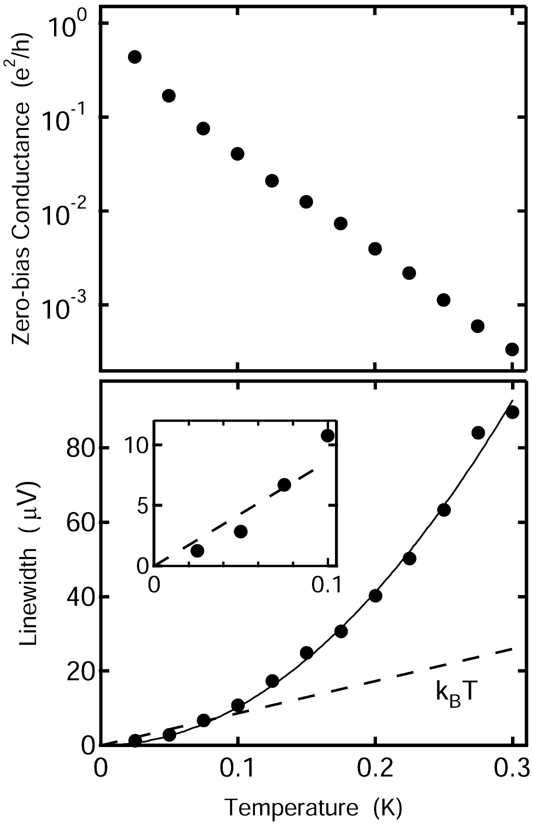

Fig. 6. Temperature dependence of height and width of zero bias tunneling conductance resonance at $\nu_T = 1$ and $d/\ell = 1.5$. Height in units of e^2/h. Note that neither the height nor the width show any evidence of saturation at the lowest temperatures.

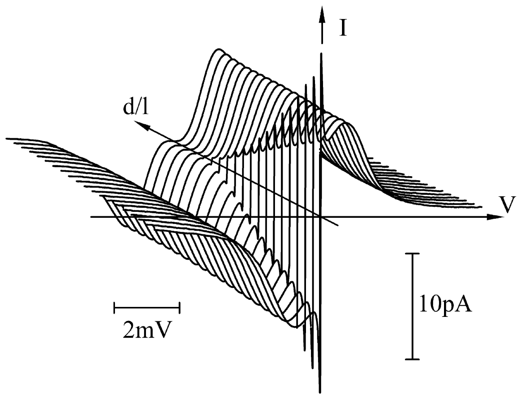

Fig. 7. Sequence of tunneling current–voltage characteristics at $\nu_T = 1$ and $T = 25$ mK, for d/ℓ values ranging from 1.9 (traces at back) to 1.6 (traces in front). Continuous development of near-discontinuity at zero bias is evident.

connection between the two has attracted much attention [23–25].

The tunneling anomalies described above offer dramatic evidence that a strongly collective phenomenon is at work. The experiments demonstrate that these anomalies occur when the $\nu_T = 1$ bilayer QHE appears in the system. The huge enhancement of the zero-bias tunneling conductance suggests that instead of being 'unaware' of the correlations present in the system, as is the case for widely-separated layers, electrons tunneling between the layers of the strongly correlated bilayer $\nu_T = 1$ QHE state are full participants in all relevant correlations and confront little or no energetic penalty in hopping between layers. This is precisely what is expected to result from the spontaneous interlayer phase coherence which is at the heart of our understanding of this remarkable collective electronic state.

Another perspective on the enhanced tunneling can be gleaned from the notion of counter-flow superfluidity. Crudely speaking, when an electron tunnels through the barrier, a localized charge deficit is produced in one layer and a charge excess is created in the other layer. In order to relax these defects, the electron gases must move toward the defect in one layer and away from it in the other. If the layers are far apart, these two processes are independent and, most importantly, very slow. In essence, as charge trys to flow radially in or out from the defects, the large magnetic field forces it to spiral around in circles. This effect, which is a really just a reflection of the very small conductivity σ_{xx}, vastly slows down the approach to equilibrium. On the other hand, in a closely-spaced bilayer at $\nu_T = 1$, the two layers are not independent. Furthermore, since relaxing the two defects requires an in-flow in one layer and an out-flow in the other, it is the conductivity for counterflow which is relevant. Since this is the anticipated superfluid mode of the system, the tunneling event can be relaxed very rapidly. From this point of view, the giant tunneling enhancement at $\nu_T = 1$ offers strong indirect evidence for excitonic superfluidity.

Although the appearance of Josephson-like tunneling in the $\nu_T = 1$ QHE phase is qualitatively understood, it has proven difficult to develop a quantitative picture [23,26,27]. For example, the magnitude of the step in the tunneling current around $V = 0$ (about 30 pA) is several orders of magntitude smaller than simple estimates would suggest. Although these estimates depend upon some poorly-known parameters (e.g. the tunnel splitting Δ_{SAS}), it seems likely that the discrepancy implies that the pseudospin field in the ground state of the system is heavily disordered on mesoscopic scales. This is not surprising, since random fluctuations in the Si donor populations induce long-range fluctuations in the 2D densities and, in the $\nu_T = 1$ QHE regime, these charge fluctuations are accompanied by quenched vorticity (merons and anti-merons) in the pseudospin field.

In spite of the lack of a quantitative understanding of the height and width of the tunneling peak, it has nonetheless proven possible to use it to extract other key information about the $\nu_T = 1$ excitonic phase. Most importantly, tunneling experiments in the presence of an adjustable in-plane magnetic field $B_\parallel$ have allowed for a direct observation of the expected linear dispersion of the pseudospin collective modes. Ignoring disorder, the tunneling conductance at $B_\parallel = 0$ is sensitive to the spectral density only at zero wavevector, $q = 0$. Indeed, one may view the observation of the giant zero bias peak in the tunneling as direct detection of the anticipated $\omega = q = 0$ Goldstone collective mode in the system. On the other hand, when $B_\parallel > 0$ tunneling provides access to spectral features at the non-zero wavevector $q = edB_\parallel/\hbar$. Consequently, the expectation [26,27] was that a finite $B_\parallel$ would split the zero bias peak into two 'derivative-shaped' resonances at voltages $eV = \pm\hbar\omega(q)$.

Fig. 8 demonstrates that the experimental results are more complex than this theoretical scenario would suggest [28]. It is clear from the figure that the tunnel spectrum is quite sensitive to the in-plane field. The giant zero bias peak shrinks rapidly as $B_\parallel$ is increased, falling a factor of 2 by $B_\parallel \approx 0.15$ T. At the same time, weak sideband features appear at small $B_\parallel$ and move to higher energies as the field is increased. These sidebands have the derivative-shape expected of the dispersing collective modes, but for modest in-plane fields they are much smaller than the zero bias peak. It is also clear that the sidebands are broader in energy than the zero bias peak and that their width increases with $B_\parallel$. Beyond about $B_\parallel \approx 0.5$ T the sidebands are lost in the residual non-resonant background tunneling conductance.

Fig. 9 shows the voltage location of the sideband features (identified as the inflection points depicted in the inset to Fig. 8) versus the wavevector $q = edB_\parallel/\hbar$. Data for three different effective layer separations d/ℓ are shown. The linear dispersion of the data evident. From the slope of the solid line in the figure we infer a velocity of about 14 km/s.

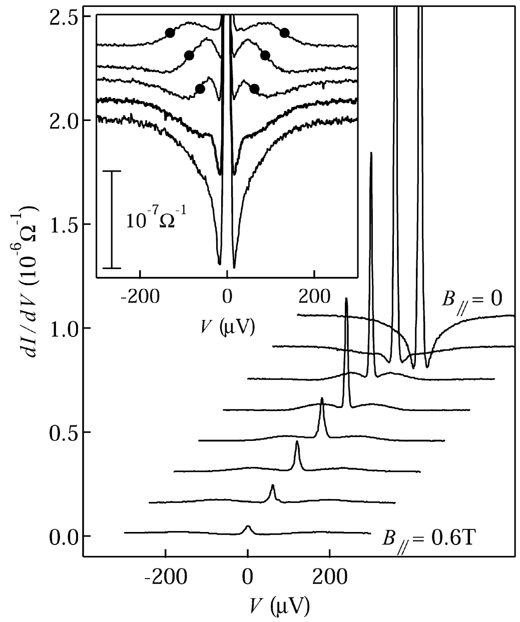

Fig. 8. Tunneling conductance spectra at $\nu_T = 1$ at $T = 25$ mK and $d/\ell = 1.6$ for various parallel magnetic fields. Main panel: $B_{\|} = 0$, 0.11, 0.24, 0.29, 0.35, 0.43, 0.49, and 0.59 T. Inset: Expanded view of spectra for $B_{\|} = 0.07$, 0.11, 0.15, 0.24, and 0.35 T. Dots indicate the positions of the sideband resonances in dI/dV. After Ref. [28].

This is in reasonable agreement with the recent theoretical estimate indicated by the dashed line [29]. These data strongly suggest that the expected linearly dispersing collective mode of the $\nu_T = 1$ bilayer QHE state has been detected.

The persistence of the zero bias peak to finite $B_{\|}$ and the relative weakness of the dispersing sidebands are puzzles. As always, disorder in the 2D system may be the source of these mysteries. For example, if the strongly correlated

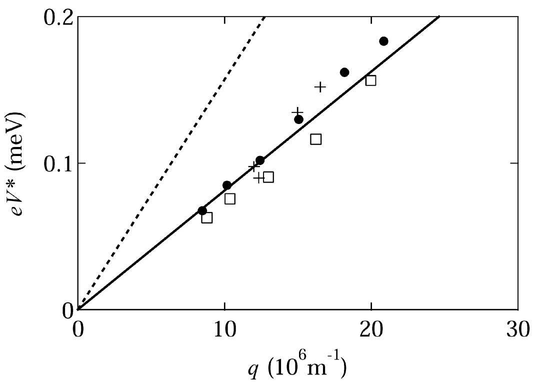

Fig. 9. Energy, eV, of sideband resonances vs. the wavevector $q = eB_{\|}d/\hbar$ induced by the in-plane field. Crosses, $d/\ell = 1.77$; empty squares $d/\ell = 1.72$; filled circles, $d/\ell = 1.6$. Dashed line is theoretical estimate [29]. Solid line is guide to the eye and corresponds to a collective mode velocity of 14 km/s. After Ref. [28].

$\nu_T = 1$ fluid only exists in mesoscopic regions of the sample, then the zero bias peak might survive to finite in-plane field. A rough analogy to the famous Fraunhofer pattern observed in small Josephson junctions can be made. A magnetic field threading flux through such a junction causes the phase to wind up transverse to the field and the net result is a reduced critical current [30]. If there were a distribution of junction sizes then the Fraunhofer interference pattern would be smeared out, but the average junction size could still be estimated. Setting aside the question of whether a 'true' Josephson effect exists at $\nu_T = 1$ or not, such an analysis of the $B_{\|}$ dependence of the zero bias tunneling peak suggests that the size of the strongly correlated regions in our samples is on the order of 1 μm. Further theoretical work is needed in order to assess the relevance of this number.

4. Conclusion

This paper has summarized the main results of our tunneling study of the bilayer quantized Hall state at $\nu_T = 1$. Tunneling has proven to be an ideal way to investigate the microscopic nature of the electronic correlations in this system. Most importantly, a giant and very sharply resonant enhancement of the zero-bias tunneling conductance has been found. This enhancement, which appears when the boundary separating the strongly correlated QHE phase from the weakly-coupled compressible phase is crossed, signals the development of spontaneous interlayer phase coherence in the system. The behavior of the tunneling in applied in-plane magnetic fields has allowed for a direct measurement of the velocity of the predicted Goldstone collective modes in the system. Taken together, these observations offer compelling support for the theoretical picture of the system as a pseudospin ferromagnet or, equivalently, a Bose condensate of interlayer excitons.

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