Double Layer Two Dimensional Electron Systems at High Magnetic Field

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Received July 21, 1995

Two examples of the new physics afforded by double layer two dimensional electron systems are described. After a short introduction, I will briefly discuss the $\nu = 1$ quantized Hall effect in bilayer systems. This integer QHE state exhibits an interesting array of many-body phenomena, including a remarkable "textural" phase transition. In the second example I will describe recent experiments on the tunneling of electrons between the two layers. In particular, the dramatic suppression, at high magnetic fields, of the tunneling density of states at the Fermi level will be described.

I. Introduction

Double layer two dimensional electron systems (2DES) have been the subject of increasingly intense scrutiny. These systems have proven extremely interesting in two rather disparate ways. On the one hand, when the layers are close enough together that interlayer tunneling and/or Coulomb interactions are strong, new electronic configurations appear that depend critically on the combination of interlayer effects with the "old-fashioned" physics of a single layer 2D system. A good example of this is the recent discovery[1,2] of a fractional quantized Hall effect (QHE) at one-half filling of the lowest Landau level. This new QHE state is distinctly bilayer in nature and has no counterpart in single layer 2D systems. On the other hand, double layer 2D electron systems have also proven very valuable by providing new ways to study a single layer system. A example of this is the development of a remarkably sensitive tool for measuring the thermodynamic compressibility of a 2DES[3]. This new technique relies on double layer 2D electron systems in which the layers are only weakly electrostatically coupled. This paper briefly describes experiments from both of these categories. I will begin with the strongly coupled case by discussing the unusual physics of the $\nu = 1$ integer quantized Hall effect in double layer systems. After that I will turn to weakly coupled structures and review the application of double layer systems to the measurement of the tunneling density of states of 2D electron gases at high magnetic field.

The experiments described below rely upon the availability of double layer 2D systems in which each layer has high mobility and well-controlled density. Both of these requirements can now be met using MBE-grown double quantum wells (DQWs) in the GaAs/AlGaAs system. A typical sample consists of two $\approx$ 200Å -wide quantum wells separated by an undoped AlGaAs barrier. Depending upon the experiment, the composition and thickness of this barrier can be adjusted to achieve the desired levels of tunnel and Coulomb interlayer coupling. Doping is achieved by depositing Si delta-layers both above and below the DQW. The precise location of these layers is determined by the desired electron densities. A typical sample contains a density of $1.5\times10^{11}cm^{-2}$ in each well. While the electron mobility depends sensitively upon the quantum well thicknesses, values exceeding $3\times10^{6}cm^{2}$/Vs (in the dark) have been achieved.

The experiments which utilize a double layer 2DES in order to study a single layer system usually require that separate electrical connections be made to the individual layers. The technique we have developed for doing this has been discussed in detail elsewhere[4]. The basic idea is quite simple. Each contact consists of a standard diffused In dot (which contacts both 2DESs) and a pair of Schottky gates "surrounding" it, one on the sample top surface and one on its back side. These gates are positioned so that any current flowing in or out of the In dot must pass the gates. By applying an appropriate bias to one of the two gates we can fully deplete the closer 2DES without adversely affecting the remote 2DES. In this way a given In contact "sees" the rest of the sample through either one or the other of the 2DESs, but not both. These contacts are low resistance (typically 100Ω), are well isolated from the undesired 2DES (by more than 50MΩ), and can be switched between layers during an experiment. The ready availability of such contacts has open up a wide class of experiments on low dimensional electronic systems.

II. The bilayer $\nu = 1$ QHE

In strongly coupled double layer 2D electron systems the filling factor ν is defined in terms of the total electron density in the sample: $\nu = N_{tot}h/eB$. Consequently, for a sample which is balanced so that each layer has the same density, $\nu = 1$ corresponds to the individual layers being at $\nu_1 = \nu_2 = 1/2$. If the layers are so far apart that they are effectively decoupled, then at this filling factor no QHE would be expected since no Hall plateau has ever been observed at $\rho_{xy} = 2h/e^2$ in a single layer 2D system. Thus, as $\nu = 1$ quantized Hall states are readily observed in bilayer systems, they must arise from interlayer couplings. As it turns out, both single-particle tunneling and interlayer Coulomb interactions can produce this quantized Hall state. On the one hand, owing to tunneling the otherwise doubly degenerate lowest subband of the DQW confinement potential is split into a symmetric and an antisymmetric state. Provided the magnitude Δ_{SAS} of the splitting is large enough, then when the Fermi level is pinned in this gap a quantized Hall plateau will appear. In this case the $\nu = 1$ many-body ground state is just a fully filled Landau level of symmetric state electrons. On the other hand, even without tunneling, interlayer electron-electron interactions can produce[5] a QHE at $\nu = 1$. The situation is analogous to what one would have in a fictitious single layer 2D system which has zero g-factor. A QHE would still appear at $\nu = 1$, only the energy gap associated with flipping a spin would be due entirely to Coulomb interactions (i.e. exchange)[6]. In the present double layer case, while we will always assume that the spins of the electrons are polarized by the Zeeman energy, the interplay of intra- and interlayer Coulomb interactions is sufficient, when the layers are close enough together, to open up an energy gap at $\nu = 1$.

The interplay of these two distinct mechanisms for producing a $\nu = 1$ quantized Hall plateau in double layer 2D electron systems has proven to be quite fascinating. My purpose here is merely outline some of the more important experimental observations. The reader interested in more detail is referred to the chapters by myself and by S.M. Girvin and A.H. MacDonald in ref. [7] and to MacDonald's contribution[8] to these conference proceedings.

The experimental[9] phase diagram for the double layer $\nu = 1$ QHE is shown in Fig. 1. On the horizontal axis is the tunneling strength, parameterized by the ratio of the symmetric/antisymmetric gap Δ_{SAS} to the characteristic Coulomb energy $e^2/\varepsilon\ell_0$ (evaluated at $\nu = 1$). The vertical axis is the ratio of the spacing d between the two 2D layers (center-to-center between the quantum wells) to the magnetic length $\ell_0 = \sqrt{\hbar/eB}$. This is essentially the ratio of the average intra- to interlayer Coulomb interactions. Each symbol in the figure represents a different double quantum well sample. In each the GaAs quantum wells are 180Å wide but the barrier parameters (thickness and alloy profile) and the 2D sheet densities differ. Solid symbols denote samples which show a quantized Hall effect at $\nu = 1$; open symbols denote those which do not. The figure demonstrates that there exists a well defined boundary, estimated by the dashed line, between a QHE phase at small d/ℓ_0 and, as expected, a non-QHE phase at larger layer separations. Perhaps surprising, however, is the clear evidence that the $\nu = 1$ QHE evolves *continuously* from a regime dominated by tunneling (large Δ_{SAS}) to one where only Coulomb interactions are important ($\Delta_{SAS} = 0$). No compressible phase appears to separate the two regimes. This suggests that while the underlying physical mechanisms are quite different, the $\nu = 1$ ground states that they lead to are smoothly connected.

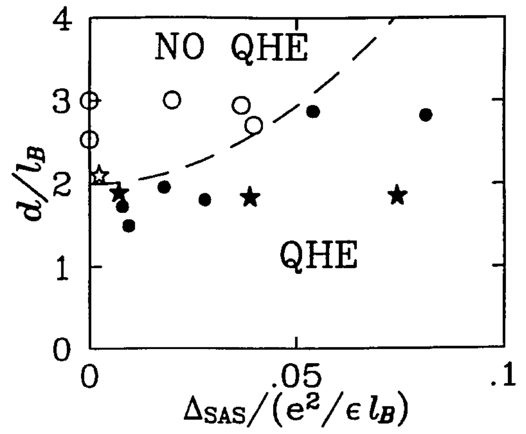

Figure 1. Phase diagram for $\nu = 1$ QHE in double layer 2D electron systems. Solid symbols denote samples which show a $\nu = 1$ QHE, open symbols denote those that do not.

More dramatic than the phase diagram in Fig. 1 is the dependence of the bilayer $\nu = 1$ QHE on an additional magnetic field component $B_{\|}$ aligned parallel to the 2D planes[9]. Fig. 2 shows the energy gap Δ at $\nu = 1$ as a function of the angle θ between the total applied magnetic field B_{tot} and the direction perpendicular to the 2D plane (the perpendicular field $B_{\perp}$, and thus the filling factor ν being kept fixed). (The gap Δ is determined from the thermally activated behavior of the resistivity ρ_{xx} of the 2D system. This quantity represents the energy required to create a well-separated quasielectron/quasihole pair. It should not be confused with the tunneling gap Δ_{SAS}, as the two are equal only under certain idealized circumstances.) The sample used for Fig. 2 is positioned deep in the many-body portion of the phase diagram in Fig. 1. It is clear from the figure that the $\nu = 1$ state is initially very sensitive to the added parallel magnetic field but then, beyond about $\theta = 8^0$, becomes insensitive to it.

The unusual angular dependence in Fig. 2 was entirely unexpected and has been attributed to a remarkable new phase transition in the $\nu = 1$ ground state. To visualize this transition it is convenient to introduce the pseudospin formalism. (Recall that the *real* electron spins are assumed to be aligned along the magnetic field by the Zeeman energy.) In this language an electron in one 2D layer is pseudospin "up" ($\sigma_z = +1$) while an electron in the opposite layer is "down" ($\sigma_z = -1$). From the Pauli spin algebra we know that an electron in a symmetric DQW state (|sym> = (|up> + |down>)/$\sqrt{2}$) is an eigenstate of pseudospin with $\sigma_x = +1$, while an antisymmetric electron has $\sigma_x = -1$. In the absence of Coulomb interactions, at $\nu = 1$ each electron is in a symmetric DQW eigenstate. Consequently, the ground state is fully pseudospin polarized along the $\hat{x}$-direction. The energy gap is just Δ_{SAS} since that is the amount of energy needed to flip an electron's pseudospin. For sufficiently small layer spacings, turning on the Coulomb interaction increases the energy gap by an exchange energy, but the ground state remains polarized along the x-axis. For the sample in Fig. 2, the gap is dominated by the exchange energy and the tunneling merely serves to orient the pseudospin. If a parallel magnetic field is now applied, it can be shown[10,11] that the tunneling matrix element acquires a phase which advances linearly across the sample in the direction perpendicular to the parallel field. As a consequence, the direction in pseudospin space which defines symmetric eigenstates is no longer $\hat{x}$ but instead rotates as one moves perpendicular to the parallel field component. Thus, if the $\nu = 1$ ground state is to maintain the energetic advantage of tunneling, then the pseudospin must acquire a twisted "texture". The wavelength λ of this texture is set by the parallel field required to inject one flux quantum between the two 2D layers: $\lambda = h/edB_{\|}$. Were it not for the Coulomb interaction, this twisting would have no observable effect. On the other hand, since neighboring electrons no longer have precisely parallel pseudospins, the twisting costs exchange energy. As the parallel field is increased, this cost rises and eventually exceeds Δ_{SAS}. the energetic benefit obtained by tracking the twisting texture. At this point the ground state becomes unstable against a phase transition[10,11] to a new, uniformly polarized, many-body state[12]. This scenario is schematically illustrated in Fig. 3. As Fig. 4 demonstrates, the calculated critical parallel field (or tilt angle θ) is in reasonably good agreement with the experiments. Interestingly, the twist wavelength at the critical angle in Fig. 2 is roughly $\lambda \approx 2500$Å. This is far larger than the magnetic length $\ell_0 \approx 120$Å and the layer separation $d = 210$Å. It suggests instead a remarkable long range phase coherence in this quantized Hall state[10,11,7,8].

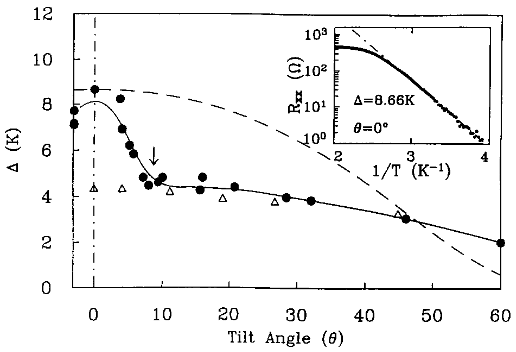

Figure 2. Energy gap Δ vs. tilt angle θ for $\nu = 1$ QHE at fixed perpendicular magnetic field $B_\perp = 5.2$T (solid dots). This sample is the leftmost solid star in the phase diagram in Fig. 1. The open triangles are for the $\nu = 2/3$ QHE and show that this QHE state exhibits no anomalous tilt behavior. The inset shows a typical Arrhenius plot of the resistivity ρ_{xx} at $\nu = 1$.

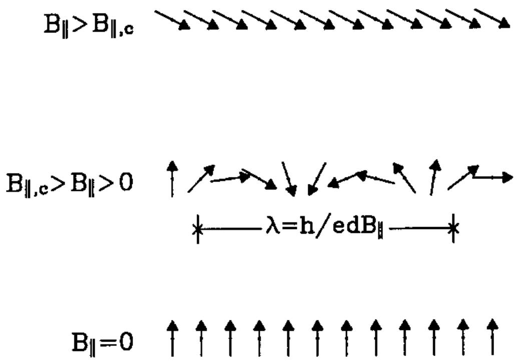

Figure 3. Schematic illustration of the textural phase transition at $\nu = 1$. Arrows represent the pseudospin field. The $\hat{x}$-axis of pseudospin space points "up" and has been arbitrarily taken to be parallel to the in-plane magnetic field $B_\parallel$. The wavelength λ of the textural twist is indicated.

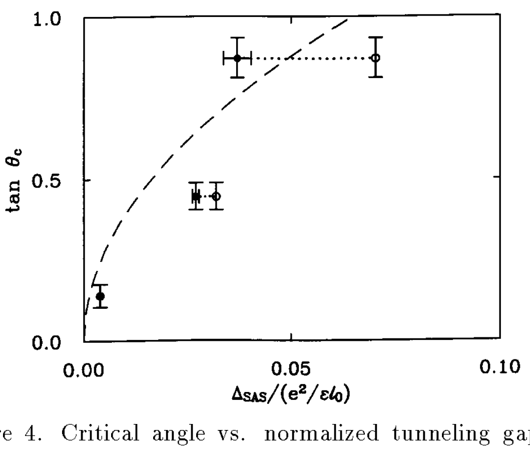

Figure 4. Critical angle vs. normalized tunneling gap for $\nu = 1$ tilted field phase transition. Dashed line is the theoretical prediction[10]. Open and solid symbols are the results of experiment. Solid symbols include a correction to the *magnitude* of the tunneling gap produced by the parallel field. Sample from figure 2 is the symbol in the lower left corner.

III. Tunneling at high magnetic field

Tunneling is an enormously powerful tool for examining many-body effects in electronic systems, as the classic studies of the density of states in superconductors demonstrate[13]. While fabricating tunnel junctions out of thin metal films and oxide layers is relatively straightforward, tunneling into a 2D electron system buried deep inside a semiconductor heterostructure presents a new challenge. Our approach to this problem has been to utilize double layer 2D systems having separate electrical contacts to the individual layers. If the layers are only weakly coupled then the measured tunneling I-V characteristic is determined by the convolution of the tunneling density of states in the two layers. Using this technique we have obtained a number of interesting results ranging from a quantitative measurement of the electron-electron scattering rate at zero magnetic field[14] to the detection of a final state interlayer exciton at high magnetic field[15].

In what follows, I will concentrate on the high field experiments. Fig. 5 compares the tunneling conductance dI/dV versus interlayer voltage V observed at zero and high magnetic field in a DQW sample consisting of two 200Å -wide quantum wells separated by a 175Å thick $Al_{0.3}Ga_{0.7}As$ barrier layer. The two 2DES are nominally identical, each having a density of about $1.5\times10^{11}cm^{-2}$ and a mobility of roughly $3\times10^{6}cm^{2}$/Vs. The 175Å tunnel barrier is sufficiently thick that the 2D layers are only very weakly coupled. Bilayer QHE states, like those discussed in the previous section, are not observable in such samples. Although the symmetric/antisymmetric gap Δ_{SAS} is typically less than 1μeV, the tunneling conductance itself is readily detectable.

As Fig. 5 shows, the tunneling conductance at B=0 exhibits a narrow resonance. This is a consequence of the conservation of energy and in-plane momentum. If these conservation laws are valid then tunneling between two 2D electron systems can only occur when there is a precise alignment of the subband energy levels in the two quantum wells[16]. For the sample used in Fig. 5, which has equal densities in the two 2D systems, this alignment occurs when the individual Fermi levels are aligned as well, i.e. at V=0. In general, however, the resonance appears at a voltage equal to the difference in the Fermi energies in the two quantum wells. The residual observed linewidth of the zero field tunnel resonance is determined by the lifetime of the electronic momentum states in the 2D systems. At low temperature this is set by the elastic scattering rate off of the static disorder in the sample, while at higher temperatures inelastic electron-electron scattering processes dominate[14].

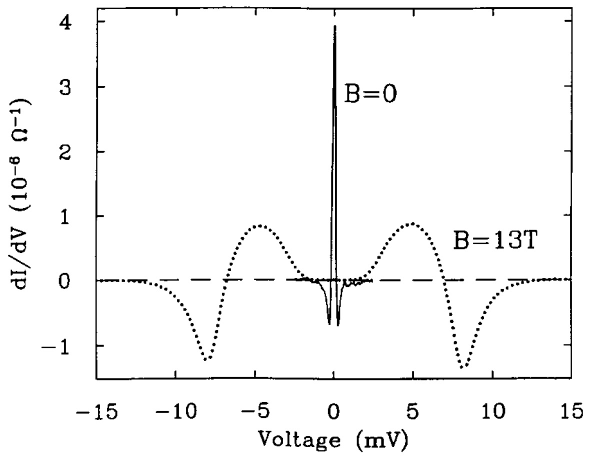

Figure 5. Low temperature tunneling conductance at zero and high magnetic field for a double layer 2D sample with a 175A $Al_{0.3}Ga_{0.7}As$ barrier.

At high magnetic fields the tunneling characteristics appear completely different[17]. For the data in Fig. 5, the 13 Tesla magnetic field is sufficient to put the Fermi level of each 2DES into the lowest spin branch of the lowest Landau level, with the individual filling factors being about $\nu_1 = \nu_2 \approx 1/2$. As the figure makes plain, in addition to being spread out over a very wide range in energy, the tunneling conductance around zero bias (V=0) is heavily suppressed. This suppression, which persists over wide ranges of magnetic field and is observed whether or not either 2DES is in a fractional quantum Hall state, represents a *pseudogap* tied to the Fermi level. Indeed, unlike the zero field tunnel resonance, this region of suppressed tunneling is always found around V=0, even if the 2DES densities are unequal. This is a purely many-body effect, and stems from the energetic penalty accompanying the rapid injection of a magnetically confined electron into an "interstitial" position in the strongly correlated electron liquid created by Landau quantization[17,18]. Even if the 2DES is thermodynamically gapless (i.e. compressible), a pseudogap appears in the tunneling spectrum so long as the injection (and extraction) process is fast compared to the time required for the 2DES to relax charge density fluctuations. In agreement with experiment, the gap magnitude is of order $e^2/\varepsilon\bar{r}$, where $\bar{r} \approx N^{-1/2}$ is the mean inter-electron separation. Above this gap the tunnel current rises to a broad peak (in Fig. 5 this is where dI/dV crosses zero near 7meV). This peak reflects all tunneling processes between the lowest Landau levels of each 2DES and its substantial width is due to electron-electron interactions, not disorder. At high magnetic fields this Coulomb broadening is, however, less than the cyclotron energy $\hbar\omega_c$ and so the tunnel current falls back to near zero at higher voltages in the (single-particle) gap between Landau levels.

This qualitative picture has been supported by recent studies[15] of the density dependence of the tunnel current. In order to change the densities of both 2DES's *in situ*, metal gate electrodes are deposited on both the top and bottom surfaces of the sample. If these gates are appropriately biased, then the densities in the two quantum wells can be kept equal but varied over a wide range. Fig. 6 shows tunneling I-V data taken with such a gated sample. The data traces shown were taken with individual 2DES densities ranging from about 0.5 to $1.3 \times 10^{11} cm^{-2}$. The magnetic fields were chosen so that each trace corresponds to Landau level filling fraction $\nu = 1/2$ in each layer. This was done since only at fixed filling factor does the qualitative model described above predict the density dependence of the tunneling spectrum be proportional to $\sqrt{N}$. Fig. 6a shows that, as expected, when the density is increased, the tunneling spectrum broadens and moves out to higher energy. In Fig. 6b the voltage axis has been divided by $\sqrt{N}$ and it is apparent that the anticipated collapse of the various data onto a single curve does not occur. Thus the simple model described above is inadequate.

In Fig. 6c the mean voltage $< V >$ and rms width Γ of the tunneling peaks shown in Fig. 6a are plotted versus $\sqrt{N}$. Both quantities are apparently linearly dependent on $\sqrt{N}$ but while Γ extrapolates to nearly zero at N=0, the mean voltage $< V >$ extrapolates to a significant negative value $V_{ex} \approx -1.4 mV$. We have recently argued[15] that this negative intercept reflects the Coulombic attraction, in the final state, between a tunneled electron and the hole it leaves behind in the source electrode. Such an effect is clearly plausible since

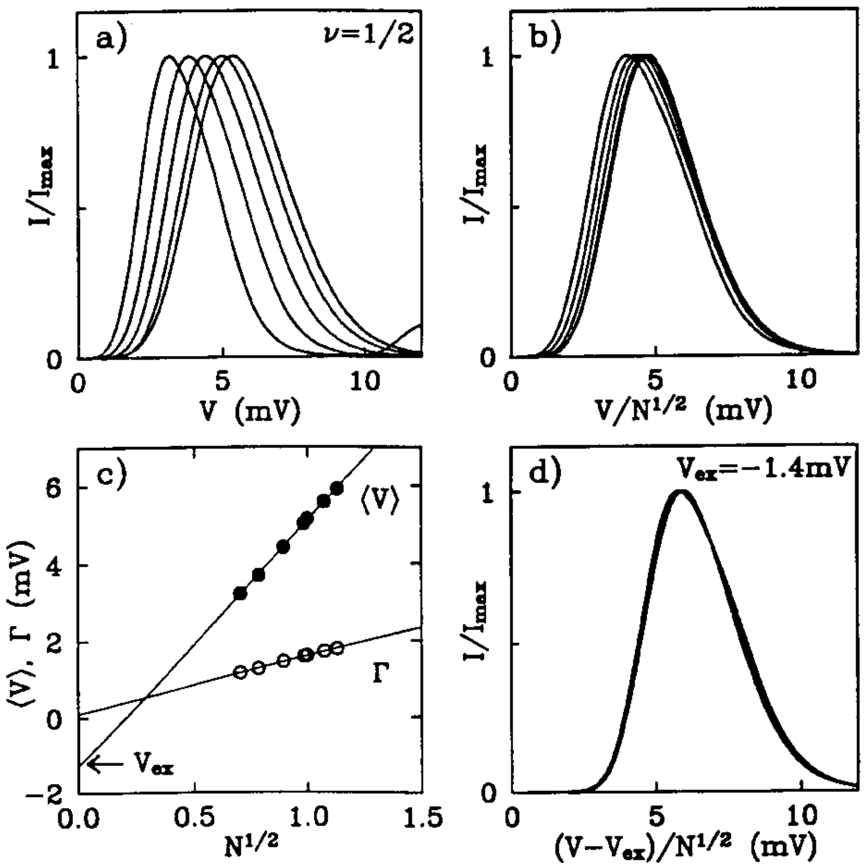

Figure 6. Density dependence of tunneling spectra at $\nu = 1/2$ from sample D. (a) Normalized tunneling I-V characteristics at T=0.6K. Left to right: N=0.62, 0.80, 1.0, 1.15, and 1.28 $\times$ $10^{11}cm^{-2}$. (b) Same as (a) but voltage axis divided by $\sqrt{N}$ (with N in units of 10^{-2}). Data does not collapse onto a single curve. (c) Mean voltage and rms width of $\nu = 1/2$ tunneling peak vs. $\sqrt{N}$ at several densities. (d) Collapse of I-V characteristics onto a single curve after subtracting $V_{ex} = -1.4mV$ from V and then dividing by $\sqrt{N}$.

the separation between the two 2D layers ($d = 400\mathring{A}$ in the sample used for Fig. 6) is comparable to the average spacing between the electrons in a given layer. Furthermore, the magnitude of the excitonic shift is quite sensible; in units of the interlayer Coulomb energy $e^2/\varepsilon d$, the shift V_{ex} comes out about 0.5. Finally, by including such a interlayer excitonic effect into the overall tunneling model, the observed density dependence makes sense. This is shown quite clearly in Fig. 6d where the data of Fig. 6a are replotted after first shifting the voltage axis of each trace by V_{ex} and then dividing by $\sqrt{N}$; now the traces do collapse onto a single curve.

It is interesting to consider why such an excitonic effect is detectable in tunneling between 2D electron systems at high magnetic field but is not observable using ordinary metal tunnel junctions. For typical 2D electron systems in GaAs, with densities in the $10^{11}cm^{-2}$ range, it is easy to apply a magnetic field large enough to force the Fermi level into the lowest Landau level, quench the kinetic energy, and thereby produce a very strongly correlated system. Under these conditions an electron tunneling in or out of the 2DES produces, initially at least, a strongly localized charge defect[18]. If both junction electrodes are two-dimensional, an interlayer exciton, analogous to a vacancy-interstitial pair, results. This exciton is relatively long-lived since not only can the electron and hole not escape into the third dimension, the magnetic field inhibits their radial spreading in the plane[18]. These conditions cannot be obtained with ordinary metal tunnel junctions. Their high electron density not only makes the screening length very short but also eliminates the possibility of reaching the lowest Landau level with laboratory magnetic fields.

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