Evidence for an Interlayer Exciton in Tunneling between Two-Dimensional Electron Systems

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Experiments on the tunneling of electrons between parallel two-dimensional systems at high magnetic field present strong evidence for a final state excitonic attraction between a tunneled electron and the hole it leaves behind. This evidence is obtained from the observed dependences of the tunneling spectra upon the thickness of the separating barrier and the densities of the electron gases. We suggest that both two dimensionality and confinement to the lowest Landau level greatly enhance the importance of this unusual exciton.

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Measurements of the quantum tunneling of electrons between two solids separated by a potential barrier date back to the seminal work of Esaki on *p-n* junctions [1]. A vast amount of literature [2] now exists, recording the numerous and varied applications of electron tunneling spectroscopy, ranging from the classic studies of superconductor densities of states to the invention of scanning tunneling microscopy. Nevertheless, one fundamental, yet conceptually simple, effect that remains elusive is the Coulombic attraction, in the final state, between a tunneled electron and the hole it leaves behind in the source electrode [3]. The absence of this excitonic effect in tunnel junctions with metal electrodes reflects, among other things, the very short screening length characteristic of ordinary metals. In this Letter, however, we report on tunneling experiments which give direct evidence for just such an exciton.

The tunnel junctions used in this experiment are fabricated from GaAs/AlGaAs double quantum well (DQW) heterostructures grown by molecular beam epitaxy. The two GaAs quantum wells are separated by a thin undoped barrier layer of $Al_xGa_{1-x}As$. Each quantum well contains a two-dimensional electron system (2DES) and forms one of the junction electrodes. The experiment consists of measuring the low temperature tunneling current-voltage (I-V) characteristics of such bilayer structures with a large magnetic field applied perpendicular to the 2D planes [4].

In each of our four samples the GaAs quantum wells are 200 Å wide. For samples A, B, and C the thickness of the barrier separating these wells is $d_b = 175$, 246, and 340 Å, respectively. In order to produce roughly equal tunneling resistances, the barrier heights V_b were adjusted. This height is controlled by the Al mole fraction x_b in the barrier, with $V_b(x_b) \approx 0.8x_b$ (eV) for $x_b \lesssim 0.4$. For samples A, B, and C, $x_b = 0.33$, 0.20, and 0.10, respectively. Silicon doping sheets deposited well above and below the DQW produce nearly equal 2DES densities in the lowest subband of each quantum well. Residual density imbalances are removed by applying, at low temperature, a small dc voltage to a gate electrode deposited on the sample backside. At balance the individual layer densities are $N = 1.55 \times 10^{11}$, 1.49×10^{11}, and 1.55×10^{11} cm^{-2} in samples A, B, and C. Sample D (having a 200 Å, $x_b = 0.33$ barrier) was grown with lower densities, $N = 1.0 \times 10^{11}$ cm^{-2}, but these could be varied from about 0.5×10^{11} to 1.3×10^{11} cm^{-2} via gates on the top *and* bottom of this wafer. Samples A–D exhibit low temperature mobilities of 3.0×10^6, 0.65×10^6, 1.0×10^6, and 1.4×10^6 cm^2/V s, respectively. Square mesas 250 μm on a side were fabricated on each sample, and separate Ohmic contacts [5] to the individual 2DES layers were placed at the ends of narrow arms extending from the central mesa square. These contacts allow direct measurement of the tunneling characteristics.

Figure 1 contains tunneling I-V characteristics for samples A and C taken at $T = 0.6$ K, with a $B = 8$ T

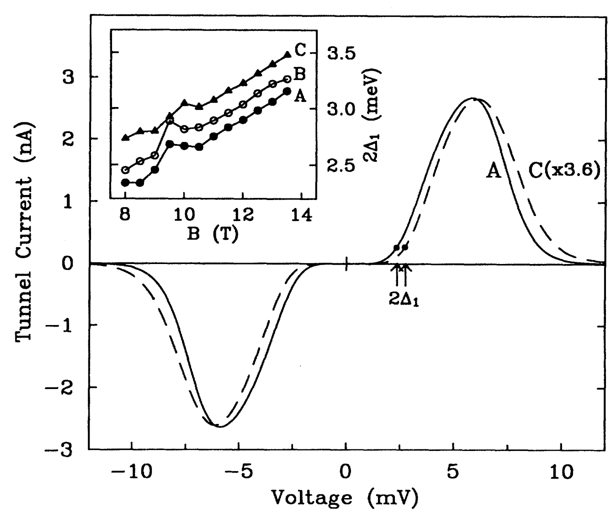

Fig. 1. Tunneling current-voltage characteristics at $B = 8$ T and $T = 0.6$ K for samples A and C. For sample C the current has been multiplied by 3.6 to aid comparison. The onset voltage $2\Delta_1$ is indicated [13]. Inset: magnetic field dependence of $2\Delta_1$ for samples A, B, and C. Tunnel barrier thicknesses are 175, 246, and 340 Å, respectively.

magnetic field applied perpendicular to the 2D layers. (Note that the current for sample C has been multiplied by 3.6.) At this magnetic field the Fermi level for each 2D layer is in the lower spin branch of the lowest Landau level, corresponding to filling fractions $\nu = hN/eB \approx 0.8$. Aside from the different magnitude of the tunneling current, the two traces are very similar. Both samples exhibit the broad region of suppressed tunneling around zero bias ($V = 0$), as reported earlier [6]. This feature, which persists over wide ranges of magnetic field and is observed whether or not either 2DES is in a fractional quantum Hall (FQH) state, represents a *pseudogap* tied to the Fermi level [7]. This purely many-body effect 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 [6,8–12]. 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 $e^2/\varepsilon\langle a\rangle$, where $\langle a\rangle \approx N^{-1/2}$ is the mean interelectron separation. Above this gap the tunnel current rises to a broad peak. 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. (Although not shown, a second peak in the current, corresponding to inter-Landau-level tunneling events, is observed [6] at higher voltages.)

Beyond the qualitative similarity of the two traces in Fig. 1 there is an important quantitative difference: The width of the tunneling gap around zero bias is larger in sample C ($d_b = 340$ Å) than in sample A ($d_b = 175$ Å). This observation, which is at odds with the assumption that the measured I-V characteristics reflect a simple convolution of independent tunneling densities of states in the two layers, suggests that interlayer Coulomb interactions play an important role in the tunneling process.

Figure 1 suggests, roughly, that the entire tunneling peak is shifted to higher energy in sample C. For simplicity, we parametrize this effect by defining a tunneling onset energy $2\Delta_1$ to be the voltage at which the current first reaches 10% of its peak (intra-Landau-level) value [13]. The arrows in Fig. 1 indicate the value of $2\Delta_1$. The inset to the figure displays the magnetic field dependence of $2\Delta_1$ from $B = 8$ to 13.5 T for samples A, B, and C (which have nearly equal 2D densities). The increase of $2\Delta_1$ with field reflects the shrinking cyclotron radius of the electrons and the concomitant increase of the *intralayer* Coulomb energy. The enhancement of the onset energy around $B = 9$–10 T is related, we believe, to the $\nu = \frac{2}{3}$ fractional quantum Hall state. Most importantly, the data also show that $2\Delta_1$ is systematically larger in samples with thicker tunnel barriers. Although values of $2\Delta_1$ depend upon the definition of the onset energy, the *shifts* in $2\Delta_1$ between samples do not. In particular, they remain essentially unchanged if $2\Delta_1$ is redefined to be at 1%, 2%, or 5% of the peak current instead of the present 10%. This current independence is good evidence that the shifts reflect genuine spectral differences and are not artifacts of unknown series resistances in the measurement circuit.

A simple explanation for these results follows from considering the Coulomb attraction of the tunneled electron and the positively charged hole it leaves behind. Such an interlayer effect is not obviously negligible, since the quantum well separations (375–540 Å, center to center) are comparable to the distance between electrons in each layer (~270 Å). The effect of this excitonic interaction will be to *reduce* the net barrier to tunneling. In qualitative agreement with the experiment, the tunneling gap will be larger in samples with thicker barriers. In the simplest model one expects

$$2\Delta_1(d) = 2\Delta_1(\infty) - \alpha e^2/\varepsilon d\,, \qquad (1)$$

where $d = d_b + d_w$ is quantum well center-to-center spacing, and α is a numerical constant. The term $2\Delta_1(\infty)$ represents the energy cost of extracting an electron from one layer and injecting it into another in the limit that the layers are infinitely separated. For a finite layer separation this cost is reduced by the final-state attraction of the electron and hole. From the $2\Delta_1$ data shown in the inset of Fig. 1 we can estimate the coefficient α. For$B \geq 11$ T (above the structure related to the $\nu = \frac{2}{3}$ FQH effect), we find $\alpha \approx 0.4$, showing that the exciton model is not quantitatively unreasonable.

In the absence of complete control over all potentially relevant sample parameters (e.g., the level of disorder) the data in Fig. 1 offer only qualitative support for the exciton picture. To build a stronger case we have examined the tunneling characteristics of sample D for which we could make *in situ* changes in the densities in each quantum well by simply adjusting gate voltages [14]. Since the origin of the tunneling gap Δ lies in the Coulomb repulsion between electrons, its magnitude should depend on the density N. At extremely high magnetic field (i.e., $\nu \ll 1$), where the vanishing magnetic length $l_0 = \sqrt{\hbar/eB}$ renders the electrons effectively pointlike, this dependence must be $\Delta \propto N^{1/2}$. At lower fields, where l_0 is comparable to the mean electron separation $\langle a\rangle \approx N^{-1/2}$, the dependence will be more complex. A simple result is obtained, however, if we compare tunneling spectra at *fixed filling factor*, $\nu = 2\pi l_0^2 N$. Within the lowest Landau level, dimensional arguments alone imply that all relevant Coulomb energies are of the form $f(\nu)e^2/\varepsilon\langle a\rangle$, where $f(\nu)$ is a function only of ν. Thus, at fixed ν, the intra-Landau-level tunnel current should only depend on voltage through the ratio V/V_0, where $V_0 \propto N^{1/2}$. Implicit in this simple scaling relation is the assumption [15] that interlayer Coulomb interactions can be ignored.

Figure 2(a) shows representative tunneling spectra taken from sample D at different (yet balanced) densities, but all at filling factor $\nu = \frac{1}{2}$. For the traces shown, N ranges from 0.62×10^{11} to 1.28×10^{11} cm^{-2}. As expected, when the density is increased, the peak in the tunnel current broadens and moves to higher voltages. In Fig. 2(b) the voltage axis for each trace has been divided by $N^{1/2}$ (with N in units of 10^{11} cm^{-2}). Since the data do not collapse onto a single curve, the argument of the previous paragraph is apparently not valid. To examine this further, we have calculated the average voltage $\langle V \rangle$ and rms width $\Gamma = (\langle V^2 \rangle - \langle V \rangle^2)^{1/2}$ of the tunneling peak [16]. Figure 2(c) shows that both $\langle V \rangle$ and Γ *are* linearly dependent upon $N^{1/2}$, but, while Γ extrapolates to nearly zero at $N = 0$, $\langle V \rangle$ extrapolates to a significant *negative* value. This suggests that the basic density dependence of the tunneling spectrum is $N^{1/2}$ as expected, but the I-V curves are shifted downward in voltage. To illustrate this more clearly, Fig. 2(d) shows that the same data *do* collapse onto a single curve if the traces are first shifted by the constant voltage $V_{ex} = -1.4$ mV and then scaled by $N^{1/2}$.

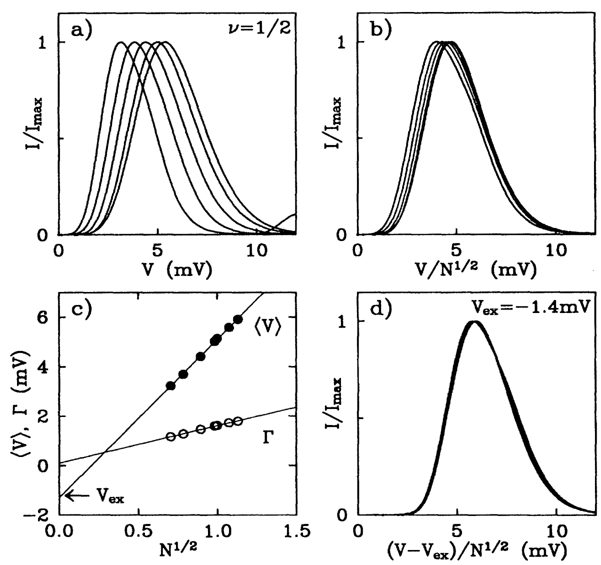

FIG. 2. Density dependence of tunneling spectra at $\nu = \frac{1}{2}$ from sample D. (a) Normalized tunneling I-V characteristics at $T = 0.6$ K. Left to right: $N = 0.62 \times 10^{11}$, 0.80×10^{11}, 1.0×10^{11}, 1.15×10^{11}, and 1.28×10^{11} cm^{-2}. (b) Same as (a) but voltage axis divided by $N^{1/2}$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

The observation of the negative shift V_{ex} in the density dependence experiment provides independent evidence for the excitonic effect invoked above to explain the barrier thickness experiment. The data in Fig. 2 have revealed that the tunnel current is a function of the combination $(V - V_{ex})/V_0$ with $V_0 \propto N^{1/2}$. This dependence provides *a fortiori* justification of Eq. (1) and identifies eV_{ex} with the excitonic term $-\alpha e^2/\varepsilon d$. Since $V_{ex} = -1.4$ mV and $d = 400$ Å in sample D, we find $\alpha = 0.5$, in good agreement with the value $\alpha \sim 0.4$ estimated independently from the barrier thickness experiment. The *in situ* nature of the density dependence experiment eliminates much of the uncertainty inherent in comparing different samples. Taken together, the two experiments offer compelling support for the existence of a final-state exciton. We remark that our findings are consistent with those theoretical discussions of 2D-2D tunneling at high magnetic field that have considered the role of interlayer excitonic interactions [9,11,12]. In fact, recent numerical studies [17] of the dependence of the I-V characteristics on layer separation are in semiquantitative agreement with the data in Fig. 1.

The density dependence of the tunneling spectrum at other filling factors $\nu < 1$ shows the same general behavior as that described above for $\nu = \frac{1}{2}$. Figure 3(a) shows the mean voltage $\langle V \rangle$ of the tunneling peak vs ν in sample D at three different densities. In Fig. 3(b) the three data sets are shown to again collapse reasonably well onto a single curve if $V_{ex} = -1.4$ mV is subtracted from $\langle V \rangle$ and the results are divided by $N^{1/2}$. The slight nonscaling behavior around the $\nu = \frac{1}{3}$ and $\frac{2}{3}$ FQH effect

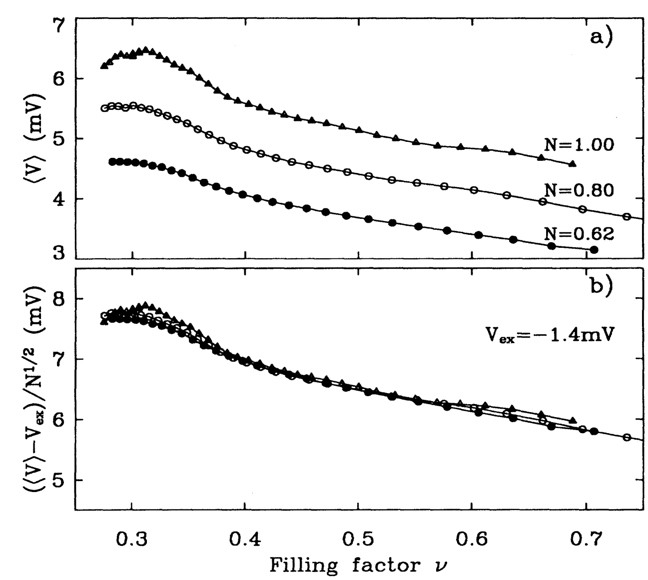

FIG. 3. (a) Mean voltage $\langle V \rangle$ of tunneling peak at $T = 0.6$ K in sample D vs filling factor for three densities (in units of 10^{11} cm^{-2}). (b) Same data, but after subtracting $V_{ex} = -1.4$ mV from V and dividing by $N^{1/2}$.

states may reflect a change in the excitonic interaction or be simply due to the influence of sample disorder on these relatively fragile correlated states. The density dependence reported here supports the view [6,8–12] that the high magnetic field tunneling gap is due, fundamentally, to the Coulomb barrier an electron confronts on injection into a strongly correlated 2DES. The magnitude of this barrier, however, is significantly reduced by interlayer excitonic interactions. In contrast to this, Brown *et al.* [18] have very recently reported finding no significant density dependence to the tunnel spectrum. They argue, instead, that the spectrum depends linearly on magnetic field alone. Our data are inconsistent with such a dependence. We can only speculate that the excitonic effect, which Brown *et al.* [18] neglect, will help to resolve this discrepancy.

The primary effect of the interlayer exciton appears to be a rigid shift of the tunneling spectrum to lower voltages. Over the ranges examined here, this shift is independent of density and magnetic field, and depends only on the interlayer separation d. This suggests that the lateral size R of the electron and hole charge density defects is small compared to d. Following He, Platzman, and Halperin [9], we can qualitatively estimate R at the peak in the tunneling by setting $e\langle V\rangle \approx e^2/\varepsilon R$. Since the measured value of $|V_{\rm ex}|$ is much less than $\langle V\rangle$, we conclude that R is much less than d. This argument breaks down deep inside the tunneling gap where R becomes large [9]. Thus, it should be interesting to examine the excitonic effect in low density, closely spaced bilayer systems.

In summary, we have presented evidence for an interlayer exciton in the tunneling of electrons between two 2D electrodes. We believe that two dimensionality and high magnetic fields have greatly enhanced the importance of this effect. 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 [9]. 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 radial spreading [9,19] in the plane. 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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