Tunneling Between Two-Dimensional Electron Gases in a Weak Magnetic Field

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We have measured the tunneling between two two-dimensional electron gases (2DEGs) in weak magnetic fields, when the carrier densities of the two electron layers are matched. At zero magnetic field, B = 0, the lineshape of the equilibrium tunneling resonance is best fit by a Lorentzian, with a linewidth which is determined by the roughness of the tunnel barrier. For B ≠ 0 with filling factors \( \nu \gg 1 \), there is a suppression of the resonant tunneling conductance about zero bias, \( V_{sd} = 0 \). This low field signature of the high field Coulomb gap shows the same linear \( B \) dependence as previously measured for \( \nu < 1 \).

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Recent experimental and theoretical studies of tunneling between two-dimensional electron gases (2DEGs) have concentrated on the high magnetic field regime. At fields when the filling factor \( \nu < 1 \), there is a strong suppression of the resonant equilibrium tunneling which is observed at zero field. This suppression has been labelled the “Coulomb gap” and is thought to be caused by electron-electron interactions. Various theoretical descriptions of the Coulomb gap are only valid in certain regimes, conveniently described by relationships between three lengths: the magnetic length \( l_B \) (determined by the applied magnetic field \( B \), \( l_B = \sqrt{h/eB} \)), the average in-plane electron spacing \( a \) (fixed by the carrier density \( n \), \( a \propto n^{-1/2} \)), and the distance \( d \) between the centers of the two quantum wells (set during sample growth).

In an earlier paper, we explored the high magnetic field regime, \( \nu < 1 \), when \( l_B/a \) is (by definition) less than unity, measuring the gap width as a function of \( n \) and \( B \). In experiments at both constant \( n \) and fixed \( \nu = 1/2 \), we observed a high field gap which was independent of carrier density, and depended linearly on \( B \). These results do not agree with many-body theories which predict that the energy gap should scale with either \( e^2/\epsilon l_B \sim \sqrt{B} \) (Refs. 3, 4), or \( e^2/\epsilon a \sim \sqrt{B} \) (Refs. 5, 10). Recently Eisenstein et al. 11 have also measured high field \( I-V \) characteristics at \( \nu = 1/2 \), investigating the gap as a function of \( d \) and \( a \). At high fields electrons can be considered to be point-like, and the interlayer electron-hole pair created by a tunneling event is expected to be reasonably long-lived. The width of the gap was argued to be determined by \( e^2/\epsilon a \) modified by a smaller excitonic term proportional to \( -e^2/ed \). Their experiments show evidence for an interlayer exciton formed by the tunnelled electron and the hole left behind. It is also proposed that the average voltage position \( \langle V \rangle \) of the tunneling current peak should scale as some function \( f \) of \( \nu \), as \( \langle V \rangle = f(\nu) \times (e^2/\epsilon a - e^2/ed) \).

It has been suggested that our high field data that we have described as having a linear \( B \) dependence, might better be described by a \( \sqrt{B} \) dependence offset from the origin. In this paper we present data for higher filling factors \( \nu > 1 \) which corroborates our previous measurement for \( \nu < 1 \), and distinguishes conclusively between gaps with \( \sqrt{B} \) and linear \( B \) dependencies. We show that for layers with matched carrier densities there is a suppression of tunneling, where the width of the gap depends linearly on \( B \) alone over the range of filling factors \( 0.4 < \nu < 10 \). In our samples we see no evidence for a change in the functional form of the width of the gap as we move from \( l_B/a < 1 \) to \( l_B/a \gg 1 \).

The results presented in this paper are from measurements on a sample from wafer C751 (the same wafer, but a different sample to that used in Ref. 3). The two GaAs quantum wells are 180 Å wide, and separated by a 125 Å wide AlGaAs barrier. The as-grown carrier densities and mobilities of the upper and lower 2DEGs are \( n_1 = 3.1 \times 10^{11} \text{ cm}^{-2} \), \( \mu_1 = 8 \times 10^5 \text{ cm}^2/\text{Vs} \) and \( n_2 = 1.8 \times 10^{11} \text{ cm}^{-2} \), \( \mu_2 = 2 \times 10^5 \text{ cm}^2/\text{Vs} \). The carrier densities in the 100 μm × 150 μm tunneling area are controlled by voltages \( V_g \) and \( V_{sd} \) to a front gate 800 Å above the upper 2DEG, and a buried backgate 3500 Å below the lower 2DEG. The proximity of the backgate to the lower 2DEG, compared with a backgate evaporated onto a thinned sample, allows us to accurately control \( n_2 \) and ensure that the carrier density is uniform over the entire tunneling area. References 3, 4 give details of the design and fabrication of our devices.

Before discussing data taken in a magnetic field, we present zero-field tunneling measurements of our sample. Figure 1 shows the differential conductance \( G = dI/dV_{sd} \) as a function of top gate voltage \( V_{g1} \), when the carrier density in the lower well was fixed at \( n_2 = 1.6 \times 10^{11} \text{ cm}^{-2} \). Resonant tunneling between the two 2DEGs occurs when their carrier densities are equal, \( n_1 = n_2 \), corresponding to simultaneous conservation of energy and momentum of the tunneling electrons. After subtracting a weak parabolic background \( (0.15 - 4.2V_{g1} + 10V_{g1}^2) \), the gate sweep characteristic in Fig. 1 has been fit to the Lorentzian lineshape

\[
G(V_{g1}) = \frac{G_0}{1 + \left( \frac{V_{g1} - V_0}{\delta V_g} \right)^2},
\]

with \( V_0 = -0.1968 \text{ V}, \delta V_g = 0.01874 \text{ V}, \text{ and } G_0 = \)
6.70 \mu S. \ V_0 is the gate voltage at which \ n_1 = n_2; the resonance can be shifted to a different value of \ V_0 by applying a different backgate voltage \ V_{BG}, thereby changing the carrier density \ n_2 of the lower 2DEG. We have also tried fitting the resonance to the derivative of a Fermi function (\mathcal{F} \sim \text{sech}^2[\delta E/(kT)]) and a Gaussian. Of the three functional forms, a Gaussian fit is the poorest, \mathcal{F} is considerably better, but the overall shape, and especially the tails of the resonance are best fit by a Lorentzian. Further evidence that \mathcal{F} is inappropriate comes from the temperature dependence of the resonance width; the width of \mathcal{F} scales with \ kT, whereas the measured width of the resonance changes little between 50 mK and 4.2 K.

The shape of the resonance as a function of \ V_{sd}, the DC bias applied between the two layers, makes us more confident of our Lorentzian fit. Assuming that equilibrium tunneling has been calculated from \mathcal{F} shows deeper regions of negative differential conductance and fits the experimental data less well than Eq. 2 derived from a Lorentzian. The full width at half maximum (FWHM) of Eq. 2 is 1.06\delta E_g (which we will call \delta E_V) and can be compared with the FWHM (2\delta E_g) of Eq. 1, after the latter has been converted to energy units via the relation \delta E_g = \delta E_{V}(dE_{V}/dV_{sd}). From Figs. 1 and 2 we measure \delta E_{V} \approx 0.54 meV and \delta E_{V} \approx 0.32 meV. Eisenstein et al. have measured \delta E_{V} \approx 0.3 meV in a much higher mobility sample, having \mu \approx 3 \times 10^6 \text{cm}^2/\text{Vs}.

The close agreement of our value of \delta E_{V} with other measurements even though our 2DEGs are more disordered by a factor of 10-20, leads us to believe that the linewidth is not caused by scattering processes within either 2DEG.

The line shape for equilibrium tunneling has been calculated to be Lorentzian, under the assumption that the width is determined by the “quantum lifetime” of the electrons, \delta E = \hbar/\tau, where \tau is a lifetime derived from disorder scattering. The experimental width \delta E_{g} of our resonance corresponds to a quantum lifetime of \tau \approx 1.2 ps. The transport lifetime in the more disordered 2DEG is \tau_{tr} = m^*\mu/e \approx 7.6 ps. Experimentally we find the width of the resonance peak (and hence \tau) to be independent of \nu over the range 0.6 < \nu < 3.5 \times 10^{11} \text{cm}^{-2}, a fact which is at odds with theoretical predictions for the \nu dependence of both the transport and single-particle relaxation (small-angle scattering) times in 2D. If scattering is predominantly due to remote ionised impurities, the scattering times are predicted to vary as \tau_{tr} \sim \nu^{3/2} and \tau_{sp} \sim \nu^{1/2}. As the linewidth does not depend on the in-plane properties of the 2DEGs, the most likely broadening mechanism is non-uniformity in the width of the tunnel barrier. If the well width \omega of 180 \text{Å} varies by up to one monolayer (\delta \omega \approx \pm 3 \text{ Å}) the broadening \delta E_g of the ground state energy \mathcal{E}_{0}, assuming a hard-wall square-well potential, will be \delta E_g = \mathcal{E}_0 \times 2\delta \omega/\omega \approx 0.44 meV, which is consistent with the measured value of \delta E_g \approx 0.54 meV.

Resonant tunneling between 2DEGs in a magnetic field is expected to occur when Landau levels (LLs) of the same index are aligned in the two wells. For wells with matched carrier densities, the resonance should occur at \ V_{sd} = 0. Figure 2 shows typical \mathcal{G}(V_{sd}) data, when the carrier densities in the two layers were matched at \ n_{1,2} = 0.95 \times 10^{11} \text{cm}^{-2}, at four magnetic fields; for clarity the traces are offset vertically. The \mathcal{B} = 0 data is plotted as points and compared with Eq. 2, shown as a solid line. There is a single peak centered about \ V_{sd} = 0, with a width determined by the scattering time \hbar/\delta E_{V}, surrounded by regions of weak negative differential conductance. At \mathcal{B} = 0.9 \text{ T} (corresponding to \nu = 4.3) there is a clear suppression of tunneling at \ V_{sd} = 0, the single zero-field peak has split into two distinct peaks. The top trace, taken at \mathcal{B} = 2.6 \text{ T}, \nu = 1.5, shows an increased separation of the split peaks and \mathcal{G}(V_{sd} = 0) is reduced. The suppression of the zero bias conductance is a manifestation of the Coulomb gap, eventually causing \mathcal{G}(V_{sd} = 0) to vanish when \nu < 1.

The splitting \Delta of the conductance peaks shown in Fig. 2 is a measure of the width of the Coulomb gap. Moreover, \Delta defined in this way is equivalent to the gap parameter in the expression \ I = I_0 \exp(-\Delta(\nu_{sd})) introduced by He et al. and which has been used to fit high field \mathcal{I}−V_{sd} characteristics. Figure 3 shows \Delta measured from peak-to-peak splittings as a function of \mathcal{B}, showing a basically linear behaviour over most of the range of magnetic field, except at particular values of magnetic field.

The measured values of \Delta at \mathcal{B} = 0.9 and 2.6 \text{ T} are in good agreement with the linear fit of Fig. 3, but the value of \Delta from the curve in Fig. 2 at \mathcal{B} = 0.4 \text{ T} (\nu = 9.8) is too large. As the LL filling approaches integral and certain fractional (2/3, 1/3) values the simple picture of tunneling from one bulk 2DEG to another is distorted by the presence of edge state transport in the two layers, and our understanding of the gap is complicated. As these filling factors are approached the magnitude of the tunneling conductance is reduced and the measured value of \Delta increases rapidly with \mathcal{B}, producing the bumps identified in Fig. 3. The high field data in Fig. 3 reproduces our previous work for \nu < 1, and in combination with the low field data in this paper convincingly supports a linear \mathcal{B} dependence over a wide magnetic field range, in spite of the complications caused by the apparent hardening of the gap at integral filling factors.

Previously, using only the high-field data, we were un-
able to determine whether the linear fit passed through the origin, or through some point close to the origin (the error in the intercept parameter being very much larger than the value of the parameter itself). It is clear from Fig. 2 that the high field points have considerably less scatter than those at low fields. This makes any link between the negative intercept and a screened exciton binding energy slightly tentative.

In the only available theory predicting a low field \((1g/a \gg 1)\) gap, it is assumed that the tunneling electron acquires extra Coulomb energy \(e^2/ea\) on entering the electron liquid. However, this is not the final state of the system and the increased local charge density relaxes back to its background value by the emission of virtual magnetoplasmons. Lorentz forces acting on the spreading carriers cause the formation of a stable vortex of energy

\[
E_0 = \frac{\hbar \omega_c}{2 \nu} \ln \left( \frac{e^2}{\hbar \nu \nu_F} \right),
\]

where \(\nu_F\) is the Fermi velocity. The energy \(E_0\) sets the scale for the gap, and Eq. 3 (ignoring the weak logarithm) has a quadratic dependence on magnetic field at fixed \(n\), \(E_0 \sim B^2\). In contrast, our experimental results show that the low \(B\) gap is linear in magnetic field, with the same constant of proportionality as that measured at high fields. In a slightly different system Ashoori et al. have measured the temperature dependence of the tunneling conductance deduced from the capacitance. They observe a suppression of tunneling which depends only on \(B\), independent of LL filling, and deduce the gap width from the activation of the conductance. Their results were interpreted using a model with a tunneling density of states with a linear gap at the Fermi energy, with a magnitude of \(0.047 \hbar \omega_c\), one tenth of our measured value. The measurements presented in this paper differ from these previous experiments in the low \(B\) regime, in that the tunneling reported here is 2D–2D tunneling (from one 2DEG to another) as opposed to 2D–3D (from a 2DEG to an \(n^+\) bulk layer). Therefore, it might be expected that we should measure a wider gap (if there is a gap in both the emitter and collector density of states) than Ashoori et al.

In principle inter-LL tunneling, between LLs in the two 2DEGs with different indices, is forbidden. In practice scattering of electrons within the tunnel barrier makes these transitions weakly allowed. For wells with matched carrier densities, this gives rise to conductance peaks at \(eV_{sd} = \pm \hbar \omega_c, \pm 2\hbar \omega_c \ldots\) where \(\omega_c = eB/m^*\) is the cyclotron frequency. In addition to the high field suppression of tunneling, it has been observed that the inter-LL tunneling peaks were shifted to slightly higher energies than the usual \(\hbar \omega_c\) spacing. The downward pointing arrows in Fig. 2 mark the positions of the inter-LL tunneling peaks. For fields in the range \(0.5 < B < 3\) T the separation of the inter-LL peaks in the two bias directions is measured to be \(2.2\hbar \omega_c\), which is larger than the expected value of \(2\hbar \omega_c\). In high fields, the inter-LL tunneling has been observed as a peak in \(I-V_{sd}\) characteristics at \(eV_{sd} \approx 1.3\hbar \omega_c\). The equivalent measurement in conductance \(dI/dV_{sd}\) has a peak at \(eV_{sd} \approx 1.1\hbar \omega_c\). Here we observe an inter-LL tunneling peak at the enhanced value of \(eV_{sd} = 1.1\hbar \omega_c\) all the way down to \(B = 0.5\) T, below which inter-LL tunneling is not clearly discernible. In the interpretation that the enhanced LL spacing is due to a reduced effective mass, our observations would suggest a constant mass reduction which is independent of filling factor, in contradiction with the predictions of Smith et al.

In conclusion, we have shown that the electron tunneling lifetime is governed by the interface roughness of the tunnel barrier, rather than by scattering within either of the 2DEGs. While we have seen some weak evidence that the linear gap does not pass through the origin \((\Delta(B = 0) \neq 0)\), and hence that there may be some excitonic effects, we have presented much stronger evidence that the linear relationship \(\Delta \sim B\) (rather than \(\sqrt{B}\) or \(B^2\)) fits the data over a wide field regime, supporting our previous measurements and contradicting available theories.

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1. J. P. Eisenstein, L. N. Pfeiffer, and K. W. West, Phys. Rev. Lett. 69, 3804 (1992).
2. J. P. Eisenstein, L. N. Pfeiffer, and K. W. West, Surf. Sci. 305, 393 (1994).
3. K. M. Brown et al., Phys. Rev. B 50, 15465 (1994).
4. K. M. Brown et al., Physica (1994), in press.
5. S.-R. E. Yang and A. H. MacDonald, Phys. Rev. Lett. 70, 4110 (1993).
6. Y. Hatsugai, P.-A. Bares, and X. G. Wen, Phys. Rev. Lett. 71, 424 (1993).
7. S. He, P. M. Platzman, and B. I. Halperin, Phys. Rev. Lett. 71, 777 (1993).
8. P. Johansson and J. M. Kinaret, Phys. Rev. Lett. 71, 1435 (1993).
9. A. L. Efros and F. G. Pikus, Phys. Rev. B 48, 14694 (1993).
10. C. M. Varma, A. I. Larkin, and E. Abrahams, Phys. Rev. B 49, 13999 (1994).
11. I. L. Aleiner, H. U. Baranger, and L. I. Glazman, preprint (1994).
12. J. P. Eisenstein, L. N. Pfeiffer, and K. W. West, Phys. Rev. Lett. 74, 1419 (1995).
13. E. H. Linfield, G. A. C. Jones, D. A. Ritchie, and J. H. Thompson, Semiconduc. Sci. Technol. 9, 415 (1993).
14. K. M. Brown et al., Appl. Phys. Lett. 64, 1827 (1994).
15. K. M. Brown et al., J. Vac. Sci. Technol. B 12, 1293 (1994).
J. P. Eisenstein, L. N. Pfeiffer, and K. W. West, Appl. Phys. Lett. 58, 1497 (1991).

Due to differences in the definition of the linewidth, there may be a factor of two discrepancy between our value of $\delta E_V$ and that of Eisenstein et al. However, this difference is much less than that due to the relative mobilities.

L. Zheng and A. H. MacDonald, Phys. Rev. B 47, 10619 (1993).

A. Gold, Phys. Rev. B 38, 10798 (1988).

P. T. Coleridge, Phys. Rev. B 44, 3793 (1992).

R. C. Ashoori, J. A. Lebens, N. P. Bigelow, and R. H. Silsbee, Phys. Rev. Lett. 64, 681 (1990).

R. C. Ashoori, J. A. Lebens, N. P. Bigelow, and R. H. Silsbee, Phys. Rev. B 48, 4616 (1993).

A. P. Smith, A. H. MacDonald, and G. Gumbs, Phys. Rev. B 45, 8829 (1992).

FIG. 1. Differential conductance $G(V_{g1})$, where the gate voltage $V_{g1}$ controls the top 2DEG carrier density $n_1$. The carrier density of the bottom 2DEG was fixed at $n_2 = 1.6 \times 10^{11}$ cm$^{-2}$. After subtracting a weak parabolic background from the raw data, the tunneling peak fits the Lorentzian lineshape given in Eq. 1.

FIG. 2. Differential tunneling conductance $G(V_{sd})$ at matched carrier densities $n_1 = n_2 = 0.95 \times 10^{11}$ cm$^{-2}$, measured at $B = 0, 0.4, 0.9, 2.6$ T (the four curves are offset vertically). The data for $B = 0$ (some points are not shown for clarity) is plotted with Eq. 2 (solid line). The splitting of the zero field resonance into two peaks defines the gap width $\Delta$.

FIG. 3. The measured gap width $\Delta$ plotted versus $B$. Open circles ($\circ$) show the gap hardening due to the presence of edge states, solid circles ($\bullet$) are unaffected by edge states. The dashed line is a least-squares fit through the points unaffected by the presence of edge states, $\Delta = 0.45h\omega_c - 0.19$.
Conductance $G$ ($\mu$S) = $\frac{G_0}{1 + \left(\frac{V_{g1} - V_0}{\delta V_g}\right)^2}$

$V_0 = -0.1968$ V
$G_0 = 6.70 \mu$S
$\delta V_g = 0.01874$ V

C751
$T = 40$ mK
\[ \Delta B = 0 \]

\[ B = 0.4 \text{ T} \]
\[ \nu = 9.8 \]

\[ B = 0.9 \text{ T} \]
\[ \nu = 4.3 \]

\[ B = 2.6 \text{ T} \]
\[ \nu = 1.5 \]

\[ \Delta = 2\delta E_V \]

\[ G \text{ (}\mu\text{S)} \]

\[ V_{sd} \text{ (mV)} \]
