Giant Magnetic Suppression of Tunneling Out of a 2D Electron System

L. Menna, S. Yücel, and E. Y. Andrei
Department of Physics and Astronomy, Rutgers University, Piscataway, New Jersey 08855
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We have measured tunneling rates out of an electron layer trapped at a liquid-helium–vacuum interface in the presence of a magnetic field. When the field is transverse to the escape direction we find a striking suppression of the tunneling rates: a field of 3000 G at 40 mK reduces the tunneling current by four orders of magnitude. As the temperature increases the magnetic suppression of tunneling diminishes until it disappears completely above 250 mK. By contrast parallel fields have no effect on the tunneling rates.

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The transport of electrons through potential barriers in the presence of a magnetic field is a powerful tool for understanding the physics of tunneling in the presence of correlations and in various applications such as the design and control of tunneling devices [1–5]. Experiments on magnetotunneling (MT) have thus far focused on semiconductor heterostructures [1,5] where the magnetic field has only a slight effect on the tunneling rates. The relative importance of the magnetic field may be estimated by comparing the time to complete a cyclotron precession, $\sim \omega_c^{-1}$, to a characteristic time scale of tunneling which is usually taken as [6] $\tau \simeq \int_{z_1}^{z_2} dz/\sqrt{2[U(z)-E]/m}$. Here $\omega_c = eB/m^*c$ is the cyclotron frequency, m, e, and E are the electronic mass, charge, and total energy, $U(z)$ is its potential energy, B is the magnetic field amplitude, and z_1, z_2 are the turning points of the potential.

In this Letter we report on measurements of MT for surface state electrons trapped at a liquid-helium–vacuum interface. In this system the barrier size and the electron densities are readily tuned *in situ* over a wide range of parameters that are not accessible in heterostructures [7–9]. The barriers can be made very wide and shallow so that $\tau \sim 20$ psec is two orders of magnitude larger than in heterostructures and the strong magnetic field limit ($\omega_c\tau \geq 1$) is reached for fields as low as 2500 G. For these barriers we observe a strong suppression of the tunneling rates in the presence of a transverse magnetic field B_t (transverse to the direction of tunneling and parallel to the electron layer). As illustrated in Figs. 1–3 the tunneling currents depend exponentially on B_t^2 and exhibit an unexpected temperature and density dependence.

Our experimental setup consists of a cylindrical cell of height 2.5±0.05 mm that is half filled with liquid helium supporting the electron disk. The electrons are prevented from escaping to the sides with a negatively biased guard ring of diameter 18±0.05 mm. A voltage V_t applied between the top plate of the cell and the grounded bottom plate creates an external field with which the electrons can be pressed toward or extracted from the surface. The cell is positioned at the center of a double split superconducting magnet that can provide up to 0.5 T of either transverse and/or parallel field. The electron density is determined from the plasmon spectrum of the electron disk which is measured with an rf spectrometer. From the known 2D screened plasmon dispersion and the geometry of the cell we calculate the density profile and the total number of electrons in the disk [7,8]. In a typical run the surface is charged to a density $n \leq 2\times10^8$ cm^{-2} at $T \sim 1$ K and then cooled to 40 mK where the rf spectrum is recorded to obtain the initial number of electrons N_i. For MT measurements, first the magnetic field is ramped up to a preselected value, and then an extracting voltage pulse of duration Δt, in the range 10^{-4}–10^3 sec, is applied to initiate a detectable escape rate. Subsequently the magnetic field is ramped down to zero and the rf spectrum is recorded again to determine the number of remaining electrons N_f, from which we obtain the tunneling rate: $W = 2(N_i - N_f)/(N_i + N_f)\Delta t$. We chose to determine the density from the zero field data so as to avoid analyzing the more complicated magnetoplasmon spectra. The tunneling rates can also be obtained from a direct measurement of the tunneling current. However, the spectroscopic technique is more sensitive because it can measure small density changes over long times, which is equivalent to detecting extremely small currents, $\simeq 10^{-18}$ A. When the rates are fast enough to produce a detectable current the two types of measurements are in good agreement.

In Fig. 1(a) we plot $W(B_t)/W_0$, the ratio of the finite field and zero field escape rates, against B_t^2 for $T \leq 250$ mK, in the regime where the zero field rates are due to tunneling [7,8]. We find that the rates can be fitted with the expression $W(B_t)/W_0 = e^{-\alpha(T)B_t^2}$ where $\alpha(T)$ is plotted in the inset of Fig. 1(a). For $T \leq 200$ mK, $\alpha(T) > 0$ corresponding to an exponential suppression of the escape with increasing magnetic field. For $T \geq 220$ mK the trend is reversed and the rates accelerate with increasing field. In Fig. 2, we compare the temperature dependence of the MT rates and the zero field tunneling rates. Below 250 mK the zero field tunneling rates do not vary with temperature whereas the MT rates depend on both temperature and magnetic field. For $T \leq 200$ mK

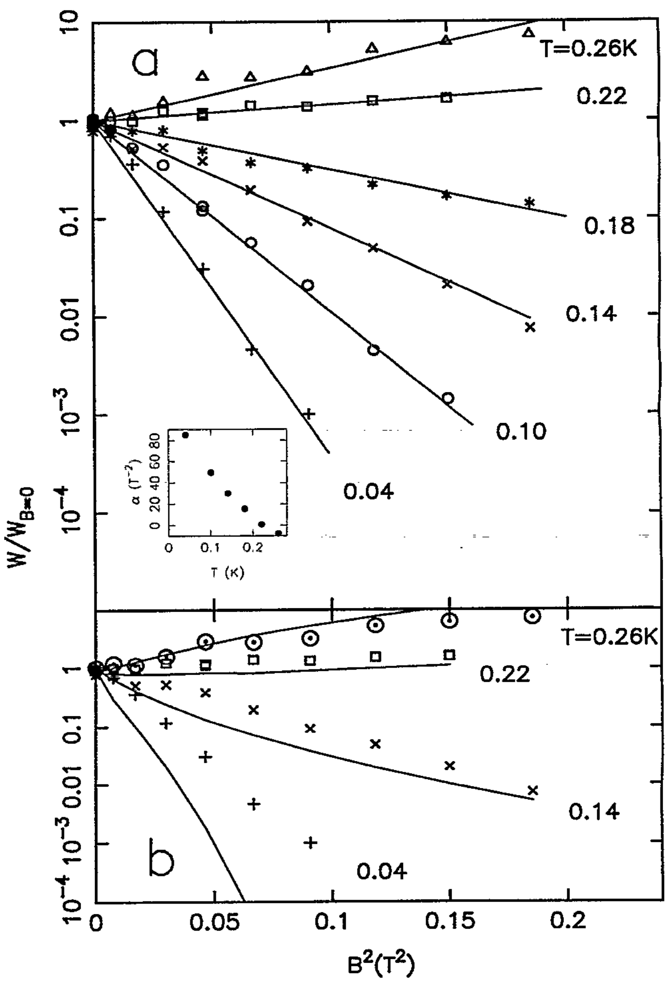

FIG. 1. Ratio of finite field escape rates to zero field rates vs B_t^2 for a series of fixed temperatures, at $n = (0.8 \pm 0.1) \times 10^8$ cm^{-2} and $V_t = 6.5$ V. (a) Measured escape rate ratios. The solid lines are guides to the eye. The inset shows the temperature dependence of α. (b) Comparison between calculated (solid lines) and measured escape rate ratios (symbols).

the MT rates decrease with increasing field: the magnetic suppression of tunneling is most pronounced at the lowest temperatures and diminishes as the temperature is raised. At higher temperatures, $T \geq 250$ mK, where the zero field rates are due to thermal activation, the temperature dependence of the rates becomes very steep and all the data approach the zero field thermally activated escape rates. As we show below, the rates for $T < 250$ mK are due to tunneling and their unusual temperature dependence can be attributed to the magnetic coupling between the in-plane and out-of-plane motions of the tunneling electron. Another unusual aspect of the MT rates is their density dependence. As shown in Fig. 3 the magnetic suppression of tunneling is most pronounced at the lowest densities and diminishes with increasing density until for $n \geq 10^8$ cm^{-2}, it appears to join the zero field rates. This signals a crossover from a regime of mag-

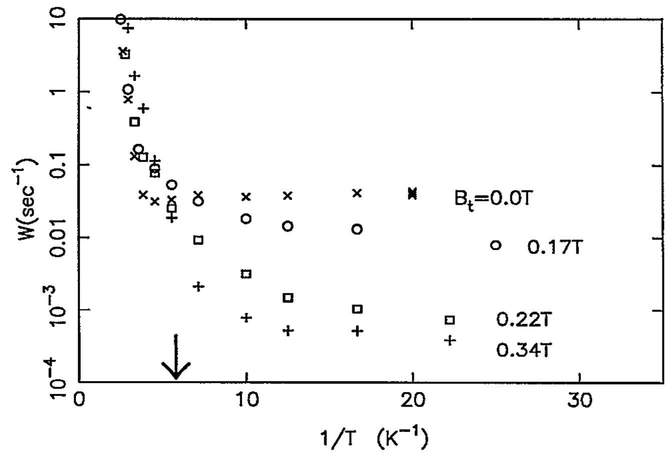

FIG. 2. Escape rates vs inverse temperature for several values of B_t. Here $n = (0.55 \pm 0.05) \times 10^8$ cm^{-2} and $V_t = 6.5$ V. The arrow marks the liquid-solid transition.

netically dominated tunneling to one that is correlation dominated as discussed below.

The measurements described here were carried out in the range of parameters where the escape mechanism in zero magnetic field has been identified [7,8,10] as tunneling through the potential barrier:

$$U(z) = \begin{cases} -\frac{Qe}{z+\beta} + e(\mathcal{E}_p + \mathcal{E}_i)z + U_c(z), & z > 0 \\ U_0, & z \leq 0 . \end{cases} \quad (1)$$

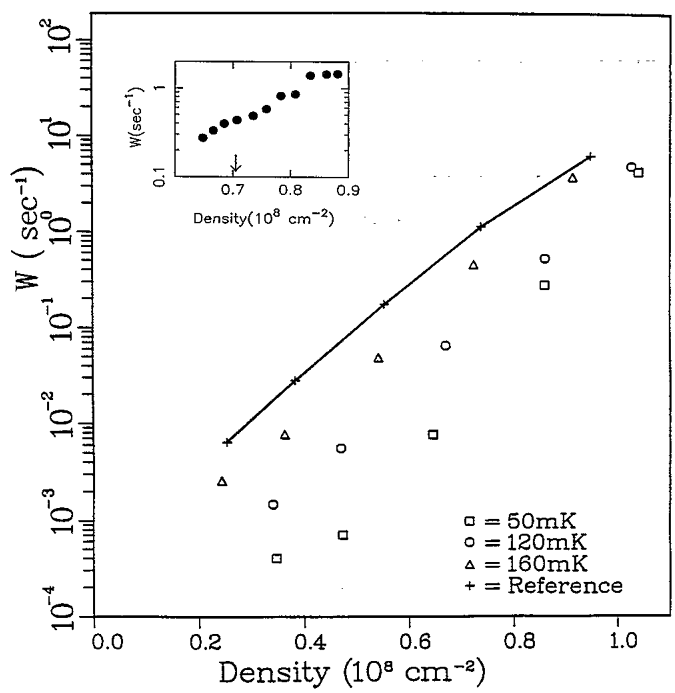

FIG. 3. Transverse MT rate vs electron density for several temperatures and $V_t = 6.5$ V. The solid line represents the zero field rates at 50 mK. For the rest of the data $B_t = 0.215$ T. The inset shows the density dependence of the MT rates near the liquid-solid transition, marked by the arrow.

Here $Q = (\epsilon - 1)/4(\epsilon + 1)$ is the image charge in the helium, $\epsilon = 1.057$ is the dielectric constant of liquid helium, $\beta = 1$ Å is a parameter obtained from fitting the potential to spectroscopic measurements [11], and $U_0 = 1$ eV is the energy of an excess electron inside the liquid helium [12]. $\mathcal{E}_p = -V_t/(h/\epsilon + d)$ is the external field, $\mathcal{E}_i = \mathcal{E}_c(d - \epsilon h)/(\epsilon h + d)$ is the field due to the image charges in the top and bottom plates located at distances d and h from the electron layer, and $\mathcal{E}_c = -2\pi ne$ is the field of a uniform charge sheet of density ne. The helium surface is kept close to the electrical midpoint of the cell, $d = \epsilon h = 1.28$ mm, so that $\mathcal{E}_i \simeq 0$. $U_c(z)$ is an effective potential describing the interaction between an escaping electron and the charge layer [13,14]:

$$U_c(z) = e\mathcal{E}_c \begin{cases} \frac{(z-z_0)^2}{2r_0}, & z - z_0 \le r_0, \\ z - z_0 - \frac{r_0}{2}, & z - z_0 \ge r_0 \,. \end{cases} \quad (2)$$

Here $z_0 = 114$ Å is the distance of the electron layer from the helium surface and $r_0 = 1.4/\sqrt{n\pi}$ is the radius of the correlation hole left by the escaping electron. If $\lim_{z\to\infty} U(z) < E_0$ where E_0 is the ground state energy, the electron can escape by tunneling. In these experiments the range of barrier widths is imposed by the window of measurable tunneling rates: $4 \times 10^{-4} < W < 10^4$ sec^{-1}. The barrier is always narrower than r_0 so that the tunneling can only probe short range correlations in the electron layer. For the zero field experiments this model, which takes into account only static correlations and leaves out the dynamic response of the charge layer, accounts well for the observed rates.

In the presence of a transverse magnetic field the Landau gauge, $\mathbf{A} = (0, -B_t z, 0)$, renders the problem separable into the in-plane and out-of plane motions. The potential for the out-of plane motion is

$$U(z, B, p_y) = U(z) + \frac{1}{2} m\omega_c^2 z^2 - \omega_c p_y z, \quad (3)$$

where $U(z)$ is the potential in the absence of the magnetic field and the canonical momentum p_y is a constant of motion. In the magnetic contribution to the potential we first consider the quadratic term which increases the barrier and tends to suppress tunneling. If the problem were strictly one dimensional the electron could never escape since this term causes the potential to diverge far from the helium surface. However, as shown in Fig. 4, the potential has a local minimum into which the electron can tunnel and subsequently escape sideways due to the Lorentz force. In the calculations described below we assume that the electron has escaped once it has crossed the initial barrier region into this local minimum, i.e., that it does not tunnel back [15]. When $p_y = 0$ the size of the barrier depends on the relative strength of the magnetic field and the correlations: the magnetic field increases the barrier while interactions with the other electrons decrease it. When $m\omega_c^2 \ll e\mathcal{E}_c/r_0 = 1.43e^2(\pi n)^{3/2}$

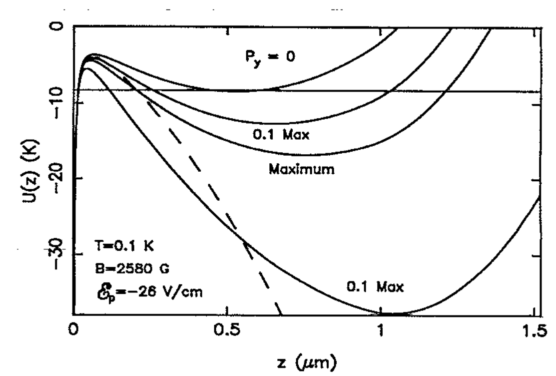

FIG. 4. Effective tunneling potential for an electron layer of density $n = 7 \times 10^7$ cm^{-2} at T=100 mK for several values of p_y: $p_y = 0$, p_y which maximizes the escape probability, and two values of p_y for which the escape probability is 0.1 of the maximum. The dashed line is the potential for $B = 0$.

the tunneling is correlation dominated while in the opposite limit it is magnetically dominated. The crossover between the two regimes becomes evident in Fig. 3. For $B_t = 0$ the tunneling increases rapidly with density indicating that it is dominated by correlations. A finite magnetic field suppresses the tunneling rates for any density, but its effect becomes negligible compared to the correlation enhancement of tunneling at sufficiently high densities.

When $p_y \neq 0$ the potential has an additional term which stems from the magnetic coupling of the in-plane and out-of-plane motions. This symmetry breaking term can either decrease or increase the potential depending on the relative orientations of the field and the momentum. Its effect for several values of $p_y \ge 0$ is illustrated in Fig. 4. It is this term that can introduce a temperature dependence of the escape rates through the statistics of the momentum distribution in the plane. To calculate the finite temperature tunneling rates we average $W(B, p_y)$, the rates at fixed p_y, over the initial momentum distribution: $W(B,T) = \int_{-\infty}^{\infty} \mathcal{P}(p_y,T) W(B,p_y) dp_y$. We chose a Maxwell-Boltzmann distribution for $\mathcal{P}(p_y, T)$ since the Fermi energy of the electrons is much smaller than the other energies in the system. The values of $W(B, p_y)$ are calculated numerically in a WKB approximation [16] with a prefactor that is adjusted to match the measured rates for $B_t = 0$. As shown in Fig. 1(b), the calculated rates agree with the measured ones for most of the data except at high fields and low temperatures where the calculated rates are too low. Some of this discrepancy may be due to an electronic temperature higher than that of the helium [17].

In the inset of Fig. 3 we again plot W vs n for parameters chosen so that the electrons undergo a phase transition from a Wigner solid to a correlated liquid. The liquid-solid phase boundary [18] is also crossed in Fig. 2 at $T = 173$ mK. We have not seen a signature of the transition in the tunneling rates with or without a magnetic field, indicating that on the short length scale probed by tunneling, $z \leq r_0$, the two phases are indistinguishable.

We have also measured tunneling in the presence of a longitudinal magnetic field and found that the longitudinal fields have no effect on the rates. This is not totally surprising since the longitudinal field, unlike its transverse counterpart, does not change the tunneling trajectory. However, it can influence the tunneling indirectly through its effect on the dynamics of the electrons and the correlation hole. Our data indicate that for the parameter range studied here dynamic correlation effects do not play a role in tunneling. This is consistent with the static picture of interactions used to arrive at the correlation potential.

In summary, the MT experiments have revealed a giant suppression of tunneling in transverse magnetic fields. The strong density dependence of the transverse MT rates is a manifestation of static correlations on the tunneling process. We attribute their unusual temperature dependence to the classical statistics obeyed by the in-plane electronic motion which couples to the out-of plane motion in the presence of a transverse field. By contrast, a longitudinal field does not change the tunneling rates, indicating that dynamical correlation effects do not play a role in the escape process. We expect the longitudinal MT rates will be sensitive to the magnetic field when the time scale of tunneling becomes comparable to the response time of the correlation hole.

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[15] The probability of tunneling back is $N_c W/\nu$, where W/ν is the probability of penetrating the barrier per collision and $\nu = E_b/\hbar \sim 10^{12}$ Hz is the escape attempt frequency. When the magnetic field is parallel to the helium layer the liberated electron moves in a skipping trajectory along the helium surface and encounters the surface $N_c \leq L\omega_c/v_d \sim 2000$ times before being absorbed by a wall, where v_d is the drift velocity and $L \sim 2$ cm. Thus for a perfectly aligned field the probability of tunneling back is $\leq 10^{-10}$, negligibly small compared to the observed suppression of tunneling. In reality the alignment is not perfect, so that the electron escapes at a slight angle eliminating the collisions altogether. Indeed, we find that the escape rates do not change when the field is slightly tilted.

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