Decay of a 2-D electron crystal

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Abstract

We have observed tunneling of electrons out of the two dimensional layer trapped at the liquid-helium vacuum interface and studied the effects of electron-electron interactions and magnetic fields on the tunneling process.

Two dimensional electron layers trapped at the surface of liquid helium or at a semiconductor heterojunction are amongst the most versatile experimental tools in condensed matter physics. They have been used extensively to study the in-plane dynamics of the electrons leading to a new understanding of the physics of correlated systems and to the discoveries of the Wigner crystal¹, the quantum Hall effect² and magneto-edge plasmons¹. In the experiments described here we focus on another aspect of these systems, the motion of electrons *out* of the layer, which allows the study of tunneling with an exceptional degree of control. Some of the unusual features in these tunneling experiments are a well defined and tunable potential barrier, interactions which can be dialed-in gradually and an observable single electron tunneling limit.

In the absence of external fields the physics of the 2D electron layer depends on the interplay between the correlation energy $V_c = e^2/\sqrt{\pi n}$, the thermal energy kT and the Fermi energy $E_F = \hbar^2 \pi n/m$, where n, m and e are the 2-D electronic density, mass, and charge, respectively. In our experiments $E_F \ll kT$, and the collective properties of the electrons are determined by the classical ordering parameter $\Gamma = V_c/kT$. For $\Gamma \ll 1$, they are in an uncorrelated gas phase. As Γ increases correlations become more important and at $\Gamma = 127$ they solidify into a Wigner Crystal.

The 2D electron layer discussed here is trapped at a liquid helium-vacuum interface. The trapping of an electron is a result of the combined effects of the attraction to the image charge in the helium substrate $Q = \frac{(\epsilon-1)}{4(\epsilon+1)}e$, and a repulsive barrier for penetrating it $U_0 = 1\text{eV}$, where $\epsilon = 1.057$ is the dielectric constant of the liquid helium. If this barrier is made infinitely high and the helium-vacuum interface perfectly sharp, the potential is that of a 1-dimensional hydrogen atom with a scaled nuclear charge Q . The resulting motion perpendicular to the helium surface is quantized in a Rydberg series with a ground state energy of $E_B = 7.6\text{K}$. This energy shifts by several percent if the finite values of U_0 and the thickness of the helium interface are taken into account.³ However, since the uncertainties in the measured parameters are too large to distinguish the details of the potential near the surface,⁴ we use the simplified form of the 1-D hydrogen atom. The wavefunction in the ground state describes an electron free to move in a plane parallel to the helium surface at a distance $z_0 = 114\text{\AA}$ above it. In an actual experiment many electrons are present. If the temperature is kept sufficiently low they all occupy the ground state forming a two dimensional electron layer. The interaction between the layer and an electron at distance z above it contributes a correlation term to the potential, $U_c(z)$. At distances $z \gg n^{-1/2}$, the potential is that of a uniform charge sheet and $U_c(z) = \mathcal{E}_c z = -2\pi n e z$, where $\mathcal{E}_c$ is the repulsive field of the layer. To stabilize the layer we apply an external field $\mathcal{E}_p$, which presses it into the helium surface. The total potential is thus:

$$U(z) = -Qe/z + e\mathcal{E}_p z + U_c(z).$$

When the asymptotic field is repulsive, $\lim_{z \rightarrow \infty} (\mathcal{E}_p + \mathcal{E}_c) < 0$, this represents a finite potential barrier through which an electron can escape by tunneling.

The experimental cell⁵ consists of two parallel metallic plates separated by $D = 2.5\text{mm}$ and a circular guard ring of diameter 18mm. The liquid helium level is $d = \epsilon h = 1.28\text{mm}$ above the bottom plate, where $h = D - d$. The electrons are deposited by a glow discharge from a tungsten filament. During the discharge the applied voltage V_t between the top and grounded bottom plate controls the

density of the electron sample. After the surface is charged V_t can be tuned to control the barrier size which depends on the applied field: $\mathcal{E}_p = -V_t/(h/\epsilon + d)$.

We used swept rf spectroscopy to determine the density of the electron layer⁸ from the resonance frequencies of the plasmon modes shown in Fig. 1. From the density together with the values of the applied voltages and the known geometry we can obtain the total number of electrons $N(t)$ to an accuracy of better than 0.5%. In a typical experiment the surface is charged to a density $n \leq 10^8 \text{ cm}^{-2}$ at $T \sim 1 \text{ K}$ and then cooled to 40mK where $N(t)$ is measured. Subsequently we apply an extracting voltage pulse of duration

Figure 1. Radio Frequency spectrum of 2-D electron disc for $n = 1.5 \times 10^8 \text{ cm}^{-2}$. The arrows indicate the positions of calculated plasmon resonances.

$\Delta t \sim 10^{-4} - 10^3$ sec to initiate a detectable escape rate. Following the pulse the rf spectrum is recorded again to determine the density and the number of remaining electrons $N(t + \Delta t)$, from which the escape rate, $W = [N(t) - N(t + \Delta t)]/N(t)\Delta t$, is obtained. The escape rates can also be obtained from a direct measurement of the current. However the spectroscopic technique is more sensitive because it can detect the loss of a small number of electrons over a long time, the equivalent of currents as low as 10^{-18} A . When the rates are fast enough to produce a detectable current the two types of measurement are in good agreement.

We measured the escape rates as a function of density, temperature and external electric and magnetic fields.^{5,7} The measurable rates span 11 orders of magnitude over a density range $0.02 - 2.2 \times 10^8 \text{ cm}^{-2}$ and temperature range of 35-450mK. The top plate voltage was varied in the range $V_t = +10 \text{ V}$ to -15 V corresponding to an external electric field in the range $\mathcal{E}_p = -40 \text{ V/cm}$ to $+60 \text{ V/cm}$. In Fig. 2 we display the density dependence of the escape rates at 40 mK, for a series of fixed values of V_t . Most of the density range explored at this temperature is in the Wigner crystal phase as indicated by the values of Γ . We distinguish between three regimes where the escape is qualitatively different. Region A: (The tunneling regime) is defined by $W > 4 \times 10^{-4} \text{ sec}^{-1}$ and a density below the critical value $n_c(V_t)$, indicated by vertical arrows. In this region the escape depends exponentially on density and external field. At low temperatures the escape rate in region A is independent of temperature as expected of tunneling, while at higher temperatures $T \geq 200 \text{ mK}$, it is activated. In the limit of zero density where correlations are negligible the barrier is determined only by the voltage difference between the top and bottom plates and by the image potential so that the rates can be calculated precisely to provide a check on the experimental parameters. For $V_t \geq 6.5 \text{ V}$ the zero density limit of the measured escape rates agrees with the calculated rates within the uncertainty of the measured parameters.

Figure 2. Density dependence of the escape rates at 40mK for a series of fixed top plate voltages. The solid lines are calculated tunneling rates.

The largest source of uncertainty in the comparison with experiment is the actual value of the voltage difference between the top and bottom plates. This has an undetermined contribution from the contact potential difference caused by adsorbed species and trapped charges⁴ which can be as large as $V_c = 1 \text{ V}$. By assuming $V_c = +0.5 \text{ V}$ we can obtain an equally good fit of the tunneling data to the potential of reference 3. We conclude that the measured tunneling rates cannot provide a definitive test of the form of the potential near the surface. In a more recent experiment⁸ the tunneling data was also fit to the potential of reference 3. But since V_c was not measured directly, again no conclusions can be drawn about the actual form of the potential near the surface.

To interpret the results at finite density we must know U_c , the contribution of correlations with electrons in the layer. We adopt a quasistatic model, the Correlation Hole Model (CHM),^{9,10} in which correlation effects are described by the potential of a continuous charge sheet pierced by a correlation hole of radius $r_0 = 1.4n^{-1/2}$ which closes at a distance r_0 above the plane:

$$U_c(z) = e\mathcal{E}_c \begin{cases} \frac{z^2}{2r_0}, & z \leq r_0 \\ z - \frac{r_0}{2}, & z > r_0, \end{cases}$$

The solid lines in Fig. 2 are results¹¹ of numerical calculation of the tunneling rates through the CHM potential. For $V_t \geq 6.5V$ where the barriers are narrowest, the CHM accurately accounts for the density dependence of the measured escape rates. For $V_t < 5.5V$ and low densities, this model becomes less satisfactory, though it still predicts the correct order of magnitude for the rates. The activation energy obtained from an Arrhenius plot of the thermally activated rates coincides with the barrier height calculated in the CHM model. We conclude that in region A, the data is consistent with a tunneling escape mechanism and that the correlation effects can be described by the single particle effective potential of the CHM.

In Region B: (background region) defined by $W < 4 \times 10^{-4} \text{sec}^{-1}$ the rates depend weakly on density and external field and are independent of temperature. For $V_t > 0$ and low densities all the data converge to a single rate. The escape in this regime is up to fifty orders of magnitude faster than calculated tunneling rates and is not consistent with tunneling. At present we cannot attribute these rates to any known escape mechanism and they are likely of spurious origin. The temperature independent escape rates reported previously¹² were also situated in this regime and could not have been due to tunneling.

In region C (the instability regime) defined by $n > n_c(V_t)$, the rates depend strongly on the pulse duration indicating the onset of a nonlinear process leading to an instability. We have found that the instability occurs when the energy of the escaping electron is sufficient to resonantly excite a helium atom upon hitting the upper electrode, $E \geq 19.5eV$. The instability is most likely connected to the photon emitted when the helium atom de-excites.⁴

The liquid-solid phase boundary is crossed in Fig. 1 at a density of 10^7cm^{-2} . We observe no clear signature of the transition in the tunneling rates indicating that the short range order of the two phases is the same.

We have measured the effect of a magnetic field on the tunneling rates⁷ in region A, for fields $B < 0.5T$. At the lowest temperatures where the zero field escape is due to tunneling we observe a dramatic suppression of the escape rates when the magnetic field is transverse to the direction of tunneling. The escape rates, shown in Fig. 3, are suppressed exponentially with the square of the transverse field amplitude, B_x . By contrast, a longitudinal field has no effect on the tunneling rates. The magnetic suppression of the tunneling rates diminishes with increasing temperature and completely disappears above 250mK. The temperature dependence of the rates below 250mK is non activated suggesting that the escape is not due to thermal activation above the barrier but rather to a temperature dependence of the tunneling process itself. Above 250mK the escape becomes activated with an activation energy that is independent of magnetic field indicating that the field does not change the barrier height significantly.

It is easiest to understand the magnetic suppression of tunneling by considering the problem in the Landau gauge, $\vec{A} = (0, -B_x z, 0)$, which renders it separable into the in-plane and out-of plane motions.

The out-of-plane potential is then:

$$U(z, B, p_y) = U(z) + \frac{1}{2}m\omega_c^2(z - z_c)^2 - \frac{1}{2}m\omega_c^2 z_c^2, \quad z > 0$$

where $U(z)$ is the potential in the absence of the field, $\omega_c = eB_x/mc$, $z_c = p_y/m\omega_c$ is the cyclotron radius and p_y is the in-plane canonical momentum of the electron which is a constant of motion. The magnetic contribution to the potential is a parabola whose location depends on the initial momentum: for $p_y < 0$ it increases the barrier everywhere, while for $p_y > 0$ the barrier is lowered in the region $z \leq 2z_c$ and increased elsewhere. Physically this reflects the transfer of kinetic

Figure 3. Escape rate vs B^2 at $n = 8 \times 10^7 \text{cm}^{-2}$ and $V_t = 6.5V$. a) Measured rates. The solid lines are guides to the eye. b) Comparison between calculated (solid lines) and measured rates (symbols).

energy from the in-plane to the out-of-plane motion due to the magnetic coupling. At $z = z_c$ where the entire kinetic energy of motion in the plane has been transferred to the escape motion, the barrier is effectively lowered by $-p_y^2/2m$. Thus, as is illustrated in Fig. 4, the more positive p_y the lower the barrier and the more likely the escape.¹³ For a distribution of initial momenta there is a corresponding distribution of barrier sizes. In our experiments the initial momentum distribution is Maxwell-Boltzman since the Fermi temperature of the electrons is negligibly small $T_F \leq 1\text{mK} \ll T$. It is this distribution which can give rise to a temperature dependence in the tunneling rates. In Fig 3b we compare the measured escape rates with finite temperature tunneling rates calculated by averaging $W(B, p_y)$, the rates at fixed p_y , over the initial momentum distribution. The values for $W(B, p_y)$ were calculated

numerically in a WKB approximation¹¹ with a field independent prefactor that is adjusted to match the measured rates at $B_x = 0$. The agreement with the measured rates is best at low fields and high temperatures. At low temperatures and high fields as the the magnetic energy and the Fermi energy are no longer negligible a more refined model may be needed. Some of this discrepancy may also be due to an electronic temperature higher than that of the helium.

These results raise several interesting issues. The barrier dependence on initial momentum clearly favors the escape of the hottest electrons while leaving behind the colder ones. Could one then get a cooling effect by allowing the electrons to tunnel? If so this could be used to cool the electrons at low temperatures where the electron-ripplon interaction ceases to provide efficient coupling to the helium bath.

The higher temperature of the escaping electrons could also explain the puzzling absence of oscillations in the field dependence of the tunneling rates. These are expected when $\hbar\omega_c > k_B T$ and the density of states on the other side of the barrier is quantized into Landau levels.¹⁴ If T is taken as the temperature of the electron layer this condition is satisfied in most of our experiments. However the relevant temperature is that of the electrons that tunnel, and that is much higher. For example for a layer temperature of $T=100\text{mK}$ and $B_x = 2580\text{G}$ the momentum of most electrons, $p_y = \sqrt{mk_B T}$, corresponds to $z_c = 300\text{\AA}$ which gives a large barrier and low escape probability. But those electrons that are most likely to escape have $z_c = 1100\text{\AA}$ which corresponds to a temperature of 1.6K . In other words for the electrons that escape $k_B T > \hbar\omega_c$ and the Landau level quantization is smeared out unless the field is much larger, $B_x \geq 4.2\text{Tesla}$. Another new aspect of this model arises from the asymmetry of the magneto-tunneling which favors the escape of electrons with initial momentum $p_y > 0$ so that the remaining electrons are left on the average with a negative momentum $p_y < 0$. It appears that the magneto-tunneling causes the entire electron layer to recoil.

In summary we have observed tunneling of electrons out of a 2-D electron layer and demonstrated that the effect of interactions on the tunneling process can be modeled by a single particle effective potential. Our experiments have revealed a large suppression of tunneling rates in transverse magnetic fields. Their unusual temperature dependence may be understood in terms of the classical statistics obeyed by the in-plane electronic motion which couples to the tunneling motion in the presence of a transverse field.

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Figure 4. Effective potential barrier for $n = 7 \times 10^{17}\text{cm}^{-2}$ for several values of initial momentum: $p_y = 0$; p_y which maximizes the escape probability and two values of p_y for which it is 0.1 of the maximum. The dashed line is the potential for $B=0$.