Hydrogen in the Submillikelvin Regime: Sticking Probability on Superfluid ^{4}He

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We have measured the sticking probability of atomic hydrogen on a superfluid ^{4}He surface for atom energies between 100 μK and 1 mK. The sticking probability shows little variation with energy over this interval, and is close to the value 0.2. The studies were carried out using a magnetic trap that produced a density of 8×10^{13} cm^{-3} at a temperature of 100 μK.

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Advances in laser [1] and evaporative [2] cooling have made it possible to experiment with matter in the gaseous state at temperatures in the μK regime. One of the new opportunities presented by these gases is the chance to study their interaction with cold surfaces. The probability that an atom is adsorbed in a surface collision—the sticking probability $s(E)$—depends on the atom-surface potential, the coupling between the atom and the surface excitations, the temperature of the surface, and the energy E of the atom. The regime of weak interactions and low energy is of particular interest. When the coupling is weak and the energy is sufficiently low [3,4], the sticking probability for a finite-range potential must vary as $E^{1/2}$. At higher energies, however, the sticking probability can be extremely sensitive to the atom-surface potential.

Atomic hydrogen incident on a superfluid ^{4}He surface is a particularly attractive system to study: The excitations of the surface—ripplons—are well understood and the atom-surface potential is so shallow that there is only one bound state for motion perpendicular to the surface. We have measured the sticking probability for atom energies between 100 μK and 1 mK, and have observed its qualitative behavior between 1 and 20 mK. An unexpected feature of the results is that we find that the values for $s(E)$ are contrary to the well established trend [5,6] at higher energies. Our data vary only slightly as E is changed over 2 order of magnitude, and we see no evidence of the expected $E^{1/2}$ low-energy behavior.

In these experiments hydrogen atoms are trapped and evaporatively cooled in an apparatus similar to one we have used previously [7], but which has been redesigned to achieve lower temperatures and higher densities [8]. The experimental cell is a vertical cylindrical tube 4.4 cm in diameter and 65 cm long. The inner surface is covered with a saturated film of superfluid ^{4}He. The temperature of the film can be varied from 20 to 600 mK. Initially the atoms are confined in the upper part of the cell by static magnetic fields. A quadrupole trapping field confines the atoms radially. Two solenoids confine the atoms longitudinally, forming a pencil-shaped trapping region 20 cm long. The upper solenoid prevents atoms from returning to the rf-discharge atom source. The lower solenoid produces a potential-energy barrier of height E_t separating the trapped atoms from the lower section of the cell. Atoms that manage to pass over this barrier and flow into the lower section of the cell are free to collide with the walls. They are adsorbed and then rapidly recombine into H_2. A sensitive bolometer [9] is located at the bottom of the cell. The recombining atoms give up a fraction of their energy to it. The bolometer signal provides an accurate measure of the flux of atoms from the trap.

By suitably varying E_t it is possible either to cool the trapped atoms or to measure their temperature. The equilibrium temperature of the trapped gas is governed by a balance between cooling due to high-energy atoms escaping over the field barrier E_t (evaporation) and heating due to two-body dipolar decay [10] of atoms in the trap. Experimentally one finds that the equilibrium temperature is a well-defined fraction of E_t. This fraction, which varies smoothly with changes in the operating conditions of the trap, is typically between 0.05 and 0.15. If E_t is lowered at a rate that is slow compared to the equilibration rate in the trap, the gas is cooled to a new equilibrium temperature at the expense of the loss of some atoms. If E_t is lowered much faster than the equilibration rate, the trapped atoms are unable to readjust their energies. In this case the energy distribution $N(E)$ can be determined by simultaneously measuring the flux of escaping atoms and the corresponding value of E_t as the entire contents of the trap are dumped. The temperature of the gas can be found by comparing $N(E)$ with calculated distributions based on the known field profile in the trap.

Figure 1(a) displays the energy distribution of 7.4×10^{12} atoms evaporatively cooled to a temperature of 1.1 mK and central density of 1.6×10^{13} cm^{-3}. Using these conditions as a starting point, an identical sample was further cooled to a temperature of 190 μK as displayed in Fig. 1(b). During that cooling the central density increased by a factor of $(2.5\text{-}4.0)\times10^{13}$ cm^{-3}. Figure 1(c) shows the results of another evaporation in which 3×10^{11} atoms were cooled to 100 μK at a density of 8×10^{13} cm^{-3}. At this density, the critical temperature for Bose-Einstein condensation is 30 μK, only a factor of 3.5 lower.

We used our cold, trapped atoms to measure $s(E)$ at

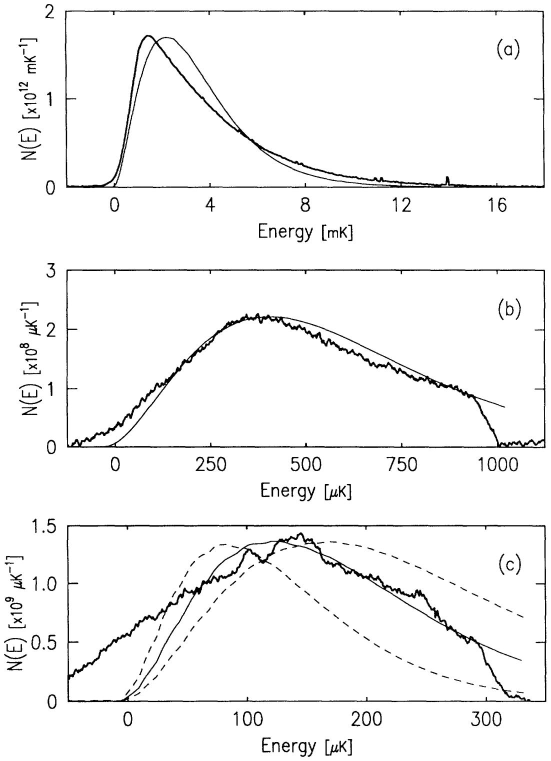

FIG. 1. Energy distributions for evaporatively cooled atomic hydrogen. The solid curves give the calculated distributions for atoms at (a) 1.1 mK, (b) 190 μK, and (c) 100 μK. The dashed curves in (c) are for temperatures 30% higher and lower than the best-fit value. We believe the finite atom signals at $E < 0$ are due to a small axial field inhomogeneity (< 1 G) that slightly deforms the bottom of the trap. Although such a deformation could influence the determination of the density, measured densities were verified with sample decay measurements.

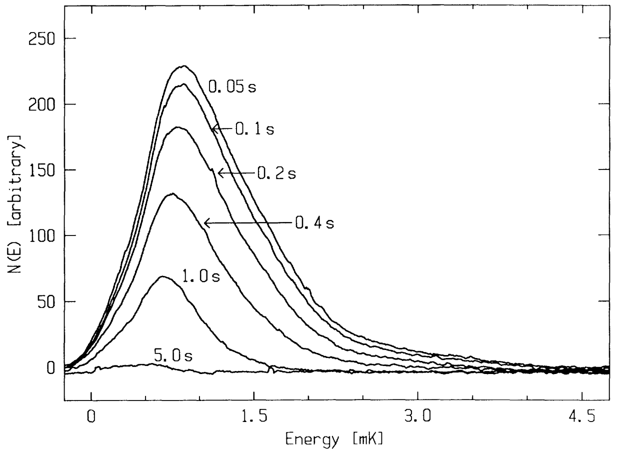

FIG. 2. Evolution of the energy distribution resulting from lowering the radial confining fields to zero for an interval of time and then raising them again. The down times for the different curves are indicated. These data are for an initial starting temperature of 500 μK. Data were taken for initial starting temperatures between 500 μK and 3 mK. The agreement between sets was good.

temperatures below 1 mK. The method starts with atoms at equilibrium in the trap with a known distribution $N(E)$. The radial confining field is rapidly lowered to zero, allowing the trapped atoms to interact with the wall. Axial confinement is maintained by the upper and lower solenoids. After a specified interval, typically 0.1 to 5 s, the radial confining field is returned to its initial value and the energy distribution is measured (see Fig. 2). While the fields are down atoms may stick to the walls where they quickly recombine [8] and are lost. By measuring the distribution after different times, we confirmed that atoms in each energy interval decayed exponentially and we determined their corresponding decay time.

In analyzing the results we assume that surface sticking is the only source of atom loss. Theoretical calculations [11,12] indicate that at these low energies the probability of inelastic nonsticking collisions is negligible. We have verified this assertion experimentally. The energy distribution of the recovered atoms does not show a contribution from atoms whose energy has been raised by their encounter with the wall. Another conceivable source of atom loss is electronic spin relaxation. If the spin flip occurs during a nonsticking wall collision, the resulting high-field-seeking atom will be lost from the trap without having been adsorbed on the wall. Analysis [8] of data taken in earlier experiments [7] indicates that the spin relaxation takes place only after the atoms have been adsorbed. Thus, the process does not influence our results.

The decay rate of the atoms depends on both the collision rate with the surface and $s(E)$. The number of atoms should decay exponentially with a rate

$$\tau(E)^{-1} = \eta(2E/MD^2)^{1/2} s(E) . \qquad (1)$$

D is the cell diameter, M is the atomic mass, and η is a dimensionless factor relating the atom's mean free path to D. For monoenergetic atoms scattering diffusely in an infinite cylinder, η would be exactly equal to 1.0. In practice, adjustments must be made for the finite length of the trapping region and the fact that some of the atoms may stick to the wall in a location where their magnetic potential energy is still finite. We have used a Monte Carlo simulation which calculates atomic trajectories to determine that $\eta = 0.95$ for our situation.

To determine $s(E)$ from the decay times involves one additional step. Because adiabatic cooling occurs as the field is lowered, the mapping of initial energies in $N(E)$ into the final atom energies at the wall must be determined. The potential energy at the wall, E_w, is proportional to the strength of the quadrupole fields responsible for the radial confinement of the atoms. As E_w is lowered, the energy E of a given trapped atom is lowered by an amount that depends on E_w. The mapping for an ideal cylindrical quadrupole potential is $E \propto E_w^{4/7}$. This is expected to be a good approximation to our situation provided that the atoms cannot reach the wall, that is, when $E < E_w$. If $E > E_w$, the wall itself helps to confine the atoms and the energy mapping is more complex, although calculable. For the conditions under which data were taken, from 50% to 80% of the adiabatic cooling occurred before E_w passed through the atom energy as E_w was reduced to zero. The time for raising or lowering the fields was > 30 times longer than the time for the atom to make one radial oscillation in the trap.

We verified that the changes associated with cycling the quadrupole fields were reversible by comparing $N(E)$ before and after cycles that were too shallow to permit any atoms to touch the walls. We found no measurable difference between the distributions before and after such cycles. We measured the energy mapping under the specific conditions used in these sticking-probability measurements and found that for $E < E_w$ it was close to the pure-quadrupole result [8]. Because the atoms begin to stick and recombine when $E > E_w$, we could not extend these tests to the fully expanded gas. Thus, when determining $s(E)$ by this method we have had to assume that our calculations for the energy mapping remain valid down to $E_w = 0$. The resulting data are plotted as circles in Fig. 3. If the energy reduction in the final stage of expansion were not as large as predicted, typically a factor of 2.5, the value of the atom energy assumed in the calculation of the sticking probability would be in error. We have computed this possible systematic error and find that it is unimportant. This method is useful in the range 100 μK to 1 mK. Below 100 μK it is limited by uncertainities in the residual magnetic fields. Above 1 mK it is limited by the rapid decay rate of the atoms.

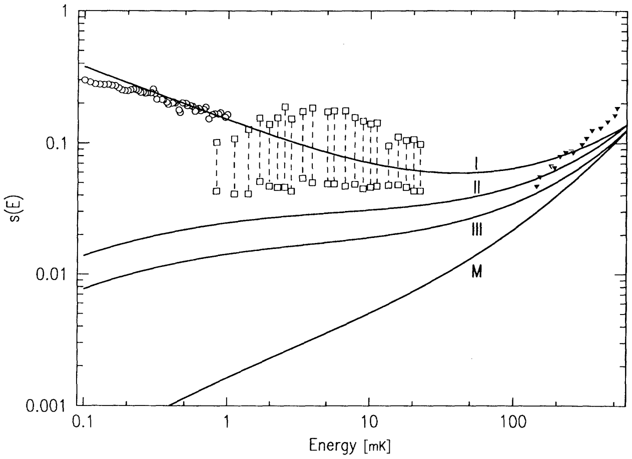

FIG. 3. Sticking probability vs atom energy. The open symbols represent our measurements. Our higher-temperature results (open squares) depend on the nature of the elastic atom-surface scattering and are plotted for the extremes of diffuse scattering, lower symbols, and specular reflection, upper symbols. The results at low temperature (open circles) do not depend on this distinction. The typical total error (including random and systematic errors) is ±35%. The solid triangles are the earlier results of Berkhout *et al.* [5]. The solid lines represent various theoretical models discussed in the text.

Another method for determining $s(E)$ is to measure the recombination time of atoms released into the detection region at the bottom of the cell. This method is applicable for energies above 1 mK. A pulse of atoms is released from the trap by lowering E_t by a small amount, typically a 10% reduction, and then quickly raising it again. The released atoms "slide" down the magnetic-field gradient and bounce about in the detection region. They are confined in one direction by the field of the lower solenoid, and in all others by ^{4}He-covered walls. The energy distribution of the released atom remains narrow and nonthermal, although the distribution of incident angles becomes random.

The response time of the bolometer to a pulse of atoms released into the lower cell was measured to be 60 ms. An exponential decay curve for the atoms, convolved with the system response function, was fitted to the bolometer signal. The resulting values of the lifetime $\tau(E)$ ranged from 60 ms at 20 mK to 30 ms at 1 mK. As in the previous method, Eq. (1) is used to determine $s(E)$ from the observed values of $\tau(E)^{-1}$. In this case, however, the value of the factor η depends on the macroscopic surface roughness of the helium film. The atoms released from the trap gain a significant amount of momentum along the axis as they slide down the potential hill; essentially, a beam is formed. If the atoms reflect specularly from the film on the bottom surface of the cell, they may oscillate up and down for some time before beginning to collide with the cylindrical walls as well. For this case our Monte Carlo calculation gives $\eta = 0.36$. If the reflection is diffuse, perhaps caused by an accumulation of solid molecular hydrogen between the helium film and the wall, the distribution of atomic velocities will randomize more quickly. In this case we calculate that $\eta = 0.93$. We cannot independently determine the reflection conditions that apply in this experiment. For this reason we have plotted in Fig. 3 the values of $s(E)$ that would result in each of the two limiting cases. In the method used below 1 mK the distribution of initial incident angles is essentially random to begin with and the question of surface roughness is unimportant. Comparison of the results of the two methods for energies where they overlap seems to indicate that the reflections are closer to specular than diffuse.

The data plotted in Fig. 3 were taken at a wall temperature of 50 mK. It has been predicted that $s(E)$ should be independent of the wall temperature as long as that temperature is well below the energy of the ripplon created when the atom is adsorbed [13]. Conservation of energy and momentum in the adsorption process set the ripplon energy to be 0.6 K. We have measured the sticking probability at several energies for wall temperatures between 40 and 80 mK. No change in $s(E)$ with wall temperature was observed within this range.

Earlier measurements of sticking probabilities at higher temperatures [5,6] yielded $s(T)$ rather than $s(E)$. Those results are plotted in Fig. 3. They are consistent with the expression [5] $s(T)=0.33T$. The lowest value of $s(T)$ obtained for pure ^{4}He in those experiments was 0.05 at 145 mK.

The solid curves in Fig. 3 are calculations of $s(T)$ by Goldman [12] for several models of atom-surface potential. Parameters in each model were adjusted to give a minimum energy of 4.5 K at a distance 0.42 nm from the surface. By way of comparison, the de Broglie wavelength of the atoms, which is proportional to $E^{-1/2}$, increases from 22 nm at 20 mK to 310 nm at 100 μK. These theoretical results, and the earlier high-temperature measurements, assume a thermal distribution of incident atoms. In Fig. 3 they are plotted at the energies of the equivalent temperatures.

The curve M results from a Morse potential, originally investigated because some of the required calculations can be done analytically. Curves I, II, and III represent more realistic approximations. The major difference between the potentials used in I and II is the energy μ_0 necessary for the hydrogen atom to penetrate into the bulk liquid: I uses $\mu_0=37$ K and II uses $\mu_0=75$ K. Recent measurements of μ_0 for deuterium [14] can be used to infer that μ_0 in hydrogen should be less than 75 K. However, considering the extreme sensitivity of the theory to the form of the potential in this energy range, we feel that it is premature to speculate on the merits of one model over another.

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