warmed to ≈ 270 °K in order to initiate a cross-channel jump. After cooling to below 50 °K, the atom appearing in the adjacent channel is pulse desorbed and its charge-to-mass ratio is measured. It is apparent from the data in Fig. 3 that after a cross-channel event, an iridium atom rather than the tungsten atom initially deposited on the surface is present in the adjacent channel.

Our measurements thus lead us to conclude that for tungsten atoms on Ir(110), cross-channel motion occurs by exchange of the tungsten adatom with an atom from the substrate, presumably from the protruding rows forming the channels. If this is indeed what happens, then we should also expect to find a tungsten atom incorporated in the substrate on analyzing the first lattice layer. These measurements have been done. After cross-channel motion of a tungsten adatom and subsequent field evaporation of the adatom remaining on the surface we do indeed find tungsten incorporated in the lattice. In a comparable number of blank runs on the substrate, tungsten was never detected.

Several other observations, which will be reported fully elsewhere, also support the conclusion that cross-channel motion of tungsten atoms on the (110) plane of iridium occurs by exchange with lattice atoms. We presume that a similar mechanism accounts for channel crossings in self-diffusion on Ir(110), and also on Pt(110). It is of interest that for both iridium[8] and platinum,[9, 10] LEED measurements on (110) yield a 1×2 pattern suggesting a channel spacing twice the normal value. In contrast, for rhodium (110), no rearrangement is indicated by LEED data,[11, 12] and neither are there any indications of cross-channel diffusion.[13] This suggests that the forces responsible for the restructuring of the (110) plane may also favor cross-channel motion by an exchange mechanism.

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Ripplons, ^{3}He, and Heat Conduction on the Surface of Superfluid ^{4}He

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A gradient in the concentration of adsorbed ^{3}He, produced by a current of ripplons, has been observed in heat-conduction measurements on the surface of superfluid ^{4}He. The ^{3}He spreading pressure and ^{3}He-ripplon collision rate have been determined from the data.

According to the phenomenological theory of Andreev and Kompaneets[1] the free surface of superfluid ^{4}He is governed by a two-dimensional (2D) version of the classic Landau theory of superfluidity. "Ripplons" (quantized capillary waves) are the elementary excitations of the surface, in analogy to phonons in the bulk liquid. ^{3}He may be added to form quasiparticles which are bound to the surface, and which, at low ^{3}He concentration, behave like a 2D Fermi gas. The 2D hydrodynamic equations of the surface are similar to the 3D, two-fluid equations in bulk helium, but with a reduced number of kinetic coefficients. So far none of the 2D kinetic coefficients have been measured.[2]

In this Letter we describe experiments on the surface thermal conductivity and, in particular, the effect of very small concentrations of adsorbed ^{3}He. The heat is carried by "two-fluid convection" in which a hydrodynamic flow of ripplons produces a heat current

$$\vec{q} = S_R T \vec{v}_R , \tag{1}$$

with a corresponding opposing flow of superfluid. In (1), $\vec{v}_R$ is the drift velocity of the ripplon gas, assumed to be in internal equilibrium via the three-ripplon process.[3] The ripplon entropy per unit area S_R is $-d\alpha_R/dT$, where $\alpha_R = -6.5\times10^{-3} \times T^{7/3}$ erg cm^{-2} is the ripplon contribution to the surface tension.[2]

With small amounts of ^{3}He on the surface we observe an effect whose 3D analog is well known, the "heat flush", in which the heat flow produces a ^{3}He concentration gradient. By determining when the heat flush is complete, i.e., when all ^{3}He has been swept off the surface, we can measure the spreading pressure of the adsorbed ^{3}He with extremely high sensitivity. Analysis of the conductivity data gives the ^{3}He-ripplon scattering time t_{3R} and its temperature dependence.

To make evaporation, and conduction by phonons, negligible, the measurements[4] were made at low temperatures and on a saturated helium film. The substrate is a vertical, free-standing, Mylar sheet, 13 μm thick, 4.1 cm wide, and 10 cm high. There is a heater at the top, and two thermometers, Acheson DAG 580 colloidal graphite films, at heights of 4.2 and 7.2 cm. The bottom of the sheet is in thermal contact with the copper cell. With enough liquid to give a saturated film, the free surface of all the film in the cell has an area $A = 1.28$ m^2 for the adsorption of ^{3}He. This was determined in earlier experiments[2] using a capacitor to measure the reduction in the surface tension by large concentrations of ^{3}He. Before cooling, the cell was flushed with neon so that the substrate of the film is actually neon-plated Mylar. The thickness of the saturated ^{4}He film was measured by flash evaporation from a heater.[4] It is approximately 100 Å at the height of the thermometers, in agreement with the van der Waals field of neon.[5] The thickness is large enough for corrections to the bulk values of α_R and S_R to be very small.

Figure 1 shows the measured thermal conductance between the thermometers, as a function of their mean temperature T. The data are for the substrate alone, the substrate with a saturated film of specially purified ^{4}He, then the ^{4}He film and substrate with two different amounts of adsorbed ^{3}He. When divided by A the latter correspond to mean ^{3}He surface densities $\bar{N}_3{}^s$ of 2.27 $\times 10^{12}$ and 8.82×10^{12} cm^{-2}. The data were taken with the cell and the bottom of the Mylar at various constant temperatures T_0, low enough for all the ^{3}He to remain on the surface. The 2D thermal conductivity of the He film is found by subtracting the conductance of the Mylar, then multiplying by the ratio of the thermometer separation to twice the width of the Mylar. For the Mylar and for the pure ^{4}He film the conductivity is accurately independent of the heat current (verified over one order of magnitude) but with ^{3}He, because of the heat flush, it depends on $\vec{q}$ and therefore on T_0.

Let us first consider the conductivity of the pure ^{4}He surface, $K_4(T)$. In bulk ^{4}He the conduction of heat by two-fluid convection is limited by

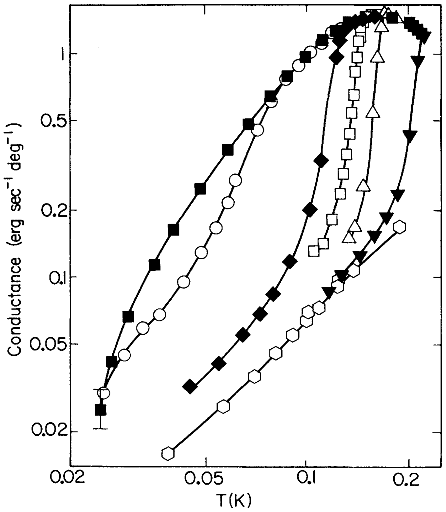

FIG. 1. Thermal conductance of neon-plated Mylar (hexagons); of Mylar covered with saturated pure ^{4}He film (solid squares); then ^{4}He with two mean surface concentrations of ^{3}He, $\bar{N}_3{}^s$, and various cell temperatures T_0: $\bar{N}_3{}^s = 2.27\times10^{12}$ cm^{-2}, T_0 = 22 mK (open circles); T_0 = 100 mK (open squares) and T_0 = 125 mK (open triangles); $\bar{N}_3{}^s = 8.82\times10^{12}$ cm^{-2}, T_0 = 22 mK (solid diamonds) and T_0 = 101 mK (solid triangles).

the viscosity of the normal fluid η_n and the boundary condition $\vec{v}_n = 0$ at the walls of the channel. The free surface, on the other hand, has no boundaries and the momentum loss per unit area, $\nabla\alpha = -S_R\nabla T = (S_R/K_4)\vec{q}$ for the pure ^{4}He surface, must be transferred to the Mylar substrate. We note from (1) that $\nabla\alpha$ and $\vec{v}_R$ are linearly related:

$$\nabla\alpha = (S_R/K_4)\vec{q} = (S_R^2T/K_4)\vec{v}_R. \qquad (2)$$

The mechanism for momentum transfer to the substrate is not yet clear, but one idea is as follows. Ripplons of energy near $k_B T$ are confined within about 20 Å of the surface, since this is their average inverse wave number k^{-1}. However, they continually emit and absorb ripplons of much smaller k by the three-ripplon process.[3] The velocity field of the low-k ripplons penetrates the full thickness of the film and so they can create or absorb phonons in the substrate. The momentum transfer and the conductivity are expected to vary with the film thickness, and from one Mylar sample to another. This is observed experimentally,[4] although there is no quantitative theory to compare with the data.

In analyzing the conductivity with ^{3}He on the surface we assume that, for a given $\vec{q}$, the momentum transfer to the substrate is the same as for the pure ^{4}He film. This is reasonable provided that all the heat is transported convectively by the ripplons, so that (1) remains valid. (Note that in the steady state the ^{3}He drift velocity is zero.) We thereby obtain a relation between $\nabla\alpha_3$, the gradient in the ^{3}He contribution to the surface tension ($-\alpha_3$ is the ^{3}He spreading pressure), and the heat current. Using $\nabla\alpha = \nabla\alpha_3 + \nabla\alpha_R$ and (2), we find

$$\nabla\alpha = \nabla\alpha_3 - S_R\nabla T = (S_R/K_4)\vec{q}. \qquad (3)$$

So

$$\nabla\alpha_3 = S_R(1 - K/K_4)\nabla T, \qquad (4)$$

where $K(T, T_0)$ is the conductivity with ^{3}He. In Fig. 1 we see that K becomes equal to K_4 when the thermometers are sufficiently hot ($T > T_s$ say) compared to the cell temperature T_0; this happens when the heat current is large enough for all the ^{3}He to be swept clear from the space between the thermometers. Integrating (4) between T_0 and T_s, using the empirical $K(T, T_0)$ and $K_4(T)$, gives the ^{3}He spreading pressure at the cell temperature, $\alpha_3(T_0)$. The results are in Fig. 2, which compares $\alpha_3(T_0)$ with values calculated from the known surface concentrations $\bar{N}_3^s$ and the ^{3}He surface effective mass and interaction.[2] (The concentrations are so small that α_3 is close to the ideal-gas value $-\bar{N}_3^s k_B T_0$.) Except for the lowest point, the agreement is very good, since there are no adjustable parameters. The lowest point may be in error as a result of an extraneous constant heat leak; this would cause additional depletion of the ^{3}He on the surface.

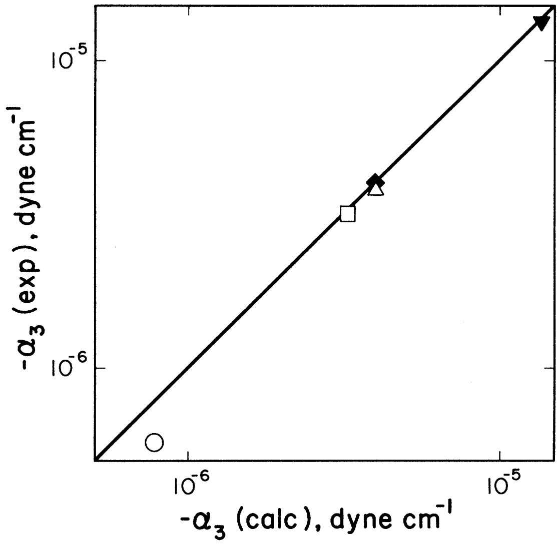

FIG. 2. Spreading pressure of ^{3}He from conductivity, $-\alpha_3$(exp), compared with theory calculated from the cell temperature T_0 and surface concentration $\bar{N}_3^s$. The symbols are the same as in Fig. 1.

The same analysis can be used to "correct" the measured conductivity K for the effect of the substrate to obtain an idealized surface conductivity. We call this K_{eff} in analogy with the corresponding 3D quantity in bulk helium.[5] It clearly depends on the local values of the ^{3}He surface number density N_3^s, and T. To obtain $K_{\text{eff}}(T, N_3^s)$ we subtract from the measured thermal resistance the resistance due to momentum transfer to the substrate:

$$1/K_{\text{eff}}(T, N_3^s) = 1/K(T, T_0) - 1/K_4(T). \qquad (5)$$

To interpret K_{eff} we need to obtain N_3^s from T_0 and T. This is done by integrating (4) to find $\alpha_3(T) - \alpha_3(T_0)$ and using the theoretical formula for $\alpha_3(T)$ to give N_3^s.

The results for K_{eff} are in Fig. 3 which shows K_{eff} vs N_3^s on a log-log graph. The temperature-dependence of K_{eff} is very slight, and we see that $K_{\text{eff}} \propto 1/N_3^s$. This is exactly as expected for scattering of the ripplons by the ^{3}He, since the ^{3}He Fermi temperature is very small and each ^{3}He atom acts independently.

We now define t_{3R}, the mean scattering time

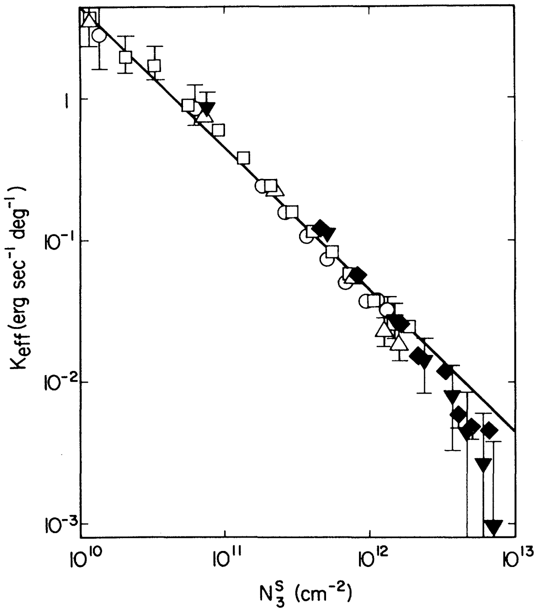

FIG. 3. The 2D thermal conductivity K_{eff}, after correcting for momentum transfer to the substrate, as a function of ^{3}He surface density. The straight line represents $K_{eff}\, N_3^s = 4.5\times 10^{10}$ erg K^{-1} cm^{-2} sec^{-1}. The symbols are the same as in Fig. 1.

for a ^{3}He atom in the ripplon gas, by the equation

$$K_{eff}N_3^s = S_R^2 T t_{3R}/m_3, \tag{6}$$

where m_3 is the bare ^{3}He mass. In bulk He the analogous 3D formula has been derived from the phonon-^{3}He Boltzmann equations.[6] For the surface we use the following argument. Suppose that, in each collision, the momentum acquired by a ^{3}He atom is $m_3\vec{v}_R$. Then the momentum transfer per unit area is $\nabla\alpha_3 = -N_3^s m_3 \vec{v}_R/t_{3R}$, and (6) follows with use of (1), (3), and (5).

If we neglect the slight temperature dependence of K_{eff}, the straight line through the data in Fig. 3 corresponds to $t_{3R} = 4.5[(0.1\text{ K})/T]^{3.67}$ μsec. Recently Saam[7] has calculated the rate at which a moving ^{3}He of momentum p emits and absorbs ripplons. With use of the mean thermal value of p his theory gives a collision time of $1.8[(0.1\text{ K})/T]^{3.5}$ μsec. This is only an approximation for t_{3R}, since a precise result requires solution of the ^{3}He-ripplon Boltzmann equations; however, the T dependence, which agrees with the experiment, should be correct.

Finally we consider the possibility that the ^{3}He quasiparticle gas has a measurable 2D conductivity K_3, requiring an additional term $-K_3\nabla T$ in (1). For a constant effective interaction $V^s(k) = V^s(0)$, effective mass M, and averaging over spin states, the 2D He^3-He^3 cross section is $\sigma_v = M|V^s(0)|^2/(4\hbar^3 v)$ for relative velocity v. Substituting in the approximate formula $K_3 \sim k_B v/(4\sigma_v)$ with $v \sim v_{rms}$ gives $K_3 \sim 7\times 10^{-5}T$ erg sec^{-1} K^{-2}. This is too small to be measurable in the present experiment.

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