

Quasi-two-dimensional spin-polarized atomic hydrogen

Yu. Kagan, G. V. Shlyapnikov, I. A. Vartan'yants, and N. A. Glukhov
I. V. Kurchatov Institute of Atomic Energy

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It is shown that quasi-two-dimensional spin-polarized hydrogen adsorbed on a surface remains an atomic gas for $T \rightarrow 0$. It is established that on a helium surface strong delocalization of an atom in a direction perpendicular to the surface accounts for the three-dimensional nature of the collisions between particles and, hence for a fundamental change in the decay kinetics of the polarized phase.

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1. In analyzing spin polarized atomic hydrogen, the properties of the quasi-two-dimensional phase, arising when atoms are adsorbed on a surface, is of special interest. Interesting in itself, in particular, from the point of view of the superfluid properties, this phase plays a critical role in the recombination kinetics of a polarized atomic hydrogen gas (see Refs. 1-3). For low adsorption energy ϵ_0 , characteristic for a surface covered with helium, hydrogen atoms move freely along the surface, forming a quasi-two-dimensional gas. The first question that arises, therefore, is: Does a bound state appear in the real triplet interaction potential between two hydrogen atoms with the transition to the two-dimensional case? If the answer is no, then a second question arises: Do drops of two-dimensional fluid form on the surface, even if the average surface density of hydrogen atoms is low?

For low adsorption energy, the characteristic localization radius of an atom l in a direction perpendicular to the surface turns out to be large. If $l \gg R_0$ ($R_0 \approx 3.5 \text{ \AA}$ is the characteristic interatomic interaction radius), then the collisions between the particles acquire largely a three-dimensional nature. This fundamentally changes the limiting filling of the surface and the decay kinetics of the polarized phase. Such a situation occurs when the surface is covered by a film of the mixture He^3 - He^4 ($\epsilon_0 \approx 0.35 \text{ K}^{4.5}$), for which $l \approx 2.5R_0$.

2. In order to analyze the question of the presence of a bound state for a pair of particles in the two-dimensional case, we use d the direct solution of Schrödinger's equation with the Kolos-Wolniewicz triplet potential,⁶ viewing the mass of the atom as a parameter. In the region of very low energies, when $kR_0 \ll 1$, only the distance s is important and the corresponding phase has the form

$$\delta_0 = -\frac{\pi}{2} \frac{1}{\ln \frac{1}{k\rho_t}} + n\pi. \quad (1)$$

Here, $\rho_t > 0$ is the two-dimensional analog of the scattering length, depending both on the interaction potential and on the mass of the particles, and n is the number of discrete levels. Figure 1 shows the dependence of δ_0 on the reduced mass m^* . The arrows indicate the values that correspond to hydrogen, deuterium, and tritium. It is evident that for tritium there is a bound state, while for $H \uparrow - H \uparrow$ and $D \uparrow - D \uparrow$, there is no bound state. The two-dimensional triplet scattering length for hydrogen turns out to be equal to $\rho_t \approx 0.9 \text{ \AA}$.

If $\delta_0(k)$ is known, it is possible to obtain directly an expression for the energy of a pair of particles

$$E(k) = \frac{2\pi\hbar^2}{mS} \frac{1}{\ln \frac{1}{k\rho_t}} \quad (2)$$

Here, S is the surface area, m is the mass of an atom, k is the relative momentum of the particles ($kR_0 \ll 1$), and a small term corresponding to the kinetic energy is omitted.

If we now take as the characteristic value of k in (2) the inverse average distance between particles ($k \sim \sqrt{\pi n_s}$, where n_s is the surface density of atoms), then the expression following from (2) for the chemical potential of a rarefied dimensional Bose gas will coincide, to within a numerical factor in the logarithm, with the exact expression ($T \rightarrow 0$, $n_s R_0^2 \ll 1$). The latter can be obtained with the help of the diagrammatic technique developed by Belyaev,⁷ and has the form

$$\mu_s = \frac{4\pi\hbar^2 n_s}{m} \frac{1}{\ln \frac{1}{4\pi n_s \rho_t^2}} \quad (3)$$

An equation of the form (3) was found previously for a model of solid disks by Schick.⁸

It is significant that in the absence of a bound state for the pair of particles, the sign

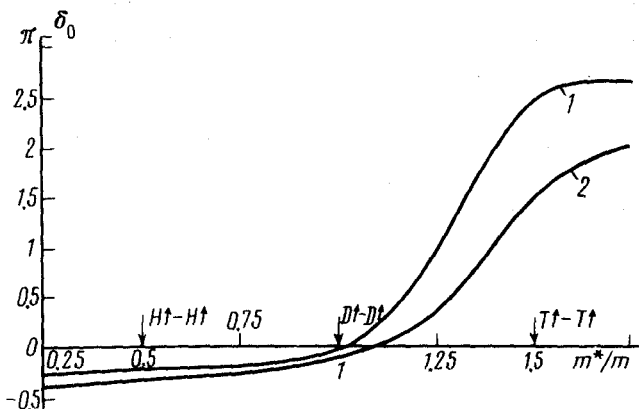


FIG. 1. The two-dimensional scattering phase δ_0 as a function of the reduced mass m^* . Curves 1 and 2 correspond to $k = 2 \times 10^{-3} \text{ \AA}^{-1}$ and $k = 2 \times 10^{-2} \text{ \AA}^{-1}$.

of the phase (1), and also the sign of the energy of the pair (2) and chemical potential (3), are strictly defined independent of the form of the interaction potential. Repulsion is also effective and this determines the stability of the gas phase at $T=0$. The results of approximate numerical calculations for the two-dimensional system of polarized atomic hydrogen examined here for high density^{9,10} also indicate that the chemical potential is positive and that it increases monotonically with increasing density. This suggests that the gaseous state of such a system is truly stable (rather than metastable).

When the gas in the bulk and on the surface is in equilibrium, keeping in mind the fact that for $n_v < 10^{20} \text{ cm}^{-3}$ the bulk chemical potential $\mu_v \ll \epsilon_0$, we obtain the following limiting values of the surface density for atomic hydrogen adsorbed on a helium film (compare Ref. 11):

$$n_s^0 \approx \begin{cases} 5 \times 10^{13} \text{ cm}^{-2}, \text{He}^4, \epsilon_0 \approx 0.9K^{1,2,12} \\ 2.5 \times 10^{13} \text{ cm}^{-2}, \text{He}^3 - \text{He}^4, \epsilon_0 \approx 0.35^{4,5} \end{cases} \quad (4)$$

3. It may be concluded from (4) that polarized atomic hydrogen is always adsorbed by a helium surface in the form of a comparatively rarefied gas phase. This assertion is all the more strengthened as ϵ_0 decreases. However, in this case, as we have already noted, the gas phase ceases to be purely two-dimensional.

Let us examine the case of quite small values ϵ_0 . Then, at distances from surface z , greater than R^* ($R^* \sim R_0$ is the effective radius of interaction of particles at the surface), the wave function of the adsorbed particle will have the form

$$\phi(z) \approx \left(\frac{2}{l}\right)^{1/2} \exp\left(-\frac{z}{l}\right),$$

where $l = \hbar / \sqrt{2m\epsilon_0}$. In this case, for a pair of particles moving along the surface with momenta $k_{1,2} \rightarrow 0$ and having an interaction potential $U(r)$, the expression for the energy can be written in the following form:

$$E = \int \Psi(\vec{\rho}, z_1, z_2) U(r) \Psi_0(\vec{\rho}, z_1, z_2) d\vec{\rho} dz_1 dz_2. \quad (5)$$

Here, $\vec{\rho}$ is the projection of the relative-distance vector between the particles \mathbf{r} on the surface, $\Psi(\vec{\rho}, z_1, z_2)$ is the true wave function of the pair, and $\Psi_0 = 1/\sqrt{S} \phi(z_1)\phi(z_2)$ is the wave function for free motion of the particles, taking into account their interaction only with the surface.

In the limit $l \gg R_0$, the main contribution to the integral (5) comes from distances from the surface $z_{1,2}$ much greater than R_0 . Explicitly separating the motion of the center mass of the pair, we have

$$\Psi \approx \frac{2}{\sqrt{Sl}} \exp\left(-\frac{z_1 + z_2}{l}\right) \chi(\mathbf{r}). \quad (6)$$

Here, $\chi(\mathbf{r})$ is the wave function of the three-dimensional relative motion of the particles.

For r appreciably greater than R_0 , the function $\chi(r) \rightarrow 1$.

Substituting (6) into (5), we find

$$E = \frac{4\pi\hbar^2 a_t}{mSl}, \quad (7)$$

where a_t is the three-dimensional scattering length. In this case, we have the following expression for the chemical potential of adsorbed gas instead of (3):

$$\mu_s = \frac{4\pi\hbar^2 a_t n_s}{ml}. \quad (8)$$

It is interesting to note that in contrast to the purely two-dimensional case, the sign of the chemical potential μ_s (8) generally can be both positive and negative, depending on the sign of a_t . For polarized atomic hydrogen, however, $a_t = 0.72 \text{ \AA} > 0$ (Refs. 3 and 13) and μ_s increases as before with increasing n_s , which corresponds to the stability of the quasi-two-dimensional gaseous phase at $T = 0$.

From the condition $\mu_s = \epsilon_0$, it follows that as ϵ_0 decreases, the limiting surface density decreases as

$$n_s^0 \sim \sqrt{\epsilon_0} \quad (9)$$

i.e., much more slowly than in the purely two-dimensional case. On a surface covered with a He^3 - He^4 film, instead of (4), we obtain

$$n_s^0 \simeq 7 \times 10^{13} \text{ cm}^{-2}. \quad (10)$$

4. When polarized hydrogen is separated in a pure spin state (the spins of the nucleus and of the electron are parallel) and the surface density is close to n_s^0 (10), the leading recombination channel on the surface is recombination with collisions of three hydrogen atoms due to the spin dipole-dipole interaction (for greater detail, see Ref. 3). If $l \gg R_0$, then this process occurs almost in the same manner as in the bulk. The main difference stems from the fact that now the wave function of the initial state has the form

$$\Psi_i \cong \frac{1}{S} \left(\frac{2}{l} \right)^{3/2} \exp \left\{ - \frac{z_1 + z_2 + z_3}{l} \right\} \chi_i \left(r_1 - r_2, r_3 - \frac{r_1 + r_2}{2} \right)$$

(χ_i is the wave function of three-dimensional motion in the center of mass system). When the particles separate to a distance much greater than R_0 , the function $\chi_i \rightarrow 1$. The final wave function can be represented as a product of the wave function of the motion of the center of mass, normalized to unity, and the wave function of the relative motion of the particles χ_f , coinciding with the analogous function for volume recombination. As a result, the rate constant for surface dipole recombination is

$$\alpha_s^d \approx \frac{4}{3l} \alpha_v^d \sim \epsilon_0, \quad (11)$$

where α_v^d is the rate constant of volume dipole recombination obtained in Ref. 3. Hence, the number of surface recombination events per unit time per unit surface area for $T \rightarrow 0$ ($n_s \approx n_s^0$) are

$$\nu_s = \alpha_s^d n_s^0 \sim \epsilon_0^{5/2}.$$

As shown in Ref. 3, the optimum method for achieving superlow temperatures, in particular, those corresponding to Bose condensation in the bulk, is accumulation of polarized hydrogen precisely in a pure spin state. In this case, the inverse effective decay time of the system ($1/\tau_{\text{eff}} \approx (\nu_s S + \nu_v V)/(n_s^0 S + n_v V)$) turns out to be minimum for volume density ($n_s^0 S \ll n_v V$)

$$\tilde{n}_v \approx \left(\frac{\alpha_s^d S}{2\alpha_v^d V} \right)^{1/3} n_s^0 \sim \epsilon_0^{5/6} \quad (12)$$

and has the magnitude

$$\frac{1}{\tau_{\text{eff}}} \approx 3\alpha_v^d \tilde{n}_v^2 \sim \epsilon_0^{5/3}. \quad (13)$$

In the case of a surface covered by a He^3 - He^4 film, the density $\tilde{n}_v \approx 4 \times 10^{18} \text{ cm}^{-3}$ and, using the quantity $\alpha_v^d \approx 4 \times 10^{-39} \text{ cm}^6/\text{s}$ ($H \approx 10^5 \text{ Oe}$) found in Ref. 3, we have $\tau_{\text{eff}} \sim 10 \text{ s}$. This time can be increased either by further decreasing ϵ_0 or due to a nonequilibrium situation leading to the condition $n_s < n_s^0$.

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