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Shear Flow and Kelvin-Helmholtz Instability in Superfluids


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The first realization of instabilities in the shear flow between two superfluids is examined. The interface separating the A and B phases of superfluid 3He is magnetically stabilized. With uniform rotation we create a state with discontinuous tangential velocities at the interface, supported by the difference in quantized vorticity in the two phases. This state remains stable and nondissipative to high relative velocities, but finally undergoes an instability when an interfacial mode is excited and some vortices cross the phase boundary. The measured properties of the instability are consistent with the classic Kelvin-Helmholtz theory when modified for two-fluid hydrodynamics.

Instabilities in the shear flow between two layers of fluids [1] belong to a class of interfacial hydrodynamics which is attributed to many natural phenomena. Examples are wave generation by wind blowing over water [2], the flapping of a flag in the wind [3,4], and even flow in granular beds [5]. In the hydrodynamics of inviscid and incompressible fluids the transition from calm to wavy interfaces is known as the Kelvin-Helmholtz (KH) instability [2,6]. Since Kelvin’s treatise in 1871, difficulties have plagued its description in ordinary fluids, which are viscous and dissipative. They also display a shear-flow instability, but its correspondence with that in the ideal limit is not straightforward. The tangential velocity discontinuity in the shear flow is created by a vortex sheet. In a viscous fluid a vortex sheet is not a stable equilibrium state and not a solution of the hydrodynamic equations [7].

Superfluids provide a close variation of the ideal inviscid limit considered by Kelvin and thus an environment where the KH theory can be tested. The initial state is a nondissipative vortex sheet—the interface between two superfluids brought into a state of relative shear flow. So far the only experimentally accessible case where this can be studied in stationary conditions is the interface between 3He-A and 3He-B [8], where the order parameter changes symmetry and magnitude but is continuous on the scale of the superfluid coherence length $\xi \sim 10$ nm. We discuss an experiment, where the two phases slide with respect to each other in a rotating cryostat: 3He-A performs solid-body-like rotation while 3He-B is in the vortex-free state and thus stationary in the laboratory frame. While increasing the rotation velocity $\Omega$, we record the events when the AB phase boundary becomes unstable—when some circulation from the A phase crosses the AB interface and vortex lines are introduced into the initially vortex-free B phase. On increasing the rotation further, the instability occurs repeatedly. Such a succession of instability events can be understood as a spin-up of 3He-B by rotating 3He-A.

Our experimental setup is shown in Fig. 1. The AB boundary is forced against a magnetic barrier in a smooth-walled quartz container, by cooling the sample below $T_{AB}$ at constant pressure in a rotating refrigerator. The number of vortices in both phases is independently determined from the simultaneously measured nuclear magnetic resonance (NMR) spectra [10,11]. The state of the sample can be changed from an all A phase to an all B phase or to a two-phase configuration. The evolution of the quantized vorticity as a function of $\Omega$ is then observed to be radically different when the AB interface is present.

The quasi-isotropic 3He-B supports singly quantized vortices with a core size comparable to $\xi$ [10]. In the anisotropic 3He-A we form vortex lines [11] with continuous Skyrme topology: Inside its central part the order-parameter amplitude remains constant, but the orientations of the axis of the orbital anisotropy cover a solid angle of $4\pi$. Such a “soft-core” structure carries continuous vorticity with two circulation quanta and is 3 orders of magnitude larger in radius than the “hard” core of the B-phase vortex. Thus converting an A-phase vortex into a B-phase vortex requires a large concentration of the flow energy.

The large difference in core radii is also the origin for the much lower rotation at which vortices start forming in the A phase at the outer sample circumference, compared to the B phase [10]. If the sample consists of only B phase, the critical velocity for forming the first vortex line in the
setup of Fig. 1 is \( v_{cB} > 7 \text{ mm/s} \) [12]. In the A phase the
critical velocity \( v_{cA} \) is a factor of 20 smaller [13],
independently of the presence of the AB boundary.

Vortex lines are thus easily created at low rotation in
the A-phase section of the sample, while no vortices are
detected in the B-phase section [Fig. 2(a)]. This is the
ideal nondissipative initial state where the two superfluid
phases slide along each other without friction. When \( \Omega \)
is increased further, sudden bursts of vortex lines are
observed in the B-phase section (Fig. 3). The onset \( \Omega_c \)
depends on temperature and on the current in the barrier
magnet. Overall, the measured characteristics of the
bursts fit a shear-flow instability, which provides a
mechanism for the cross to the AB interface.

\( ^3\text{He}-A \) and \( ^3\text{He}-B \) are states of the same order-
parameter manifold. One of the conditions on their phase
boundary is that the phase of the order parameter has to be continuous [14]: If the circulation is not continuing
across the interface [Fig. 2(a)] then the existing A-phase
vortex lines have to bend and form a vortex layer on the
AB interface. This is the only hydrodynamically stable state,
with vortex-free flow seen by the B-phase spectrometer
and a large cluster of vortex lines detected by the
A-phase spectrometer. The buildup of a sheelike vortex layer means that the A-phase circulation cannot
easily penetrate into the vortex-free B phase and the AB
interface remains stable up to the measured critical velocity \( u_c = \Omega_c R \sim 2-4 \text{ mm/s} < v_{cB} \).

At \( \Omega > \Omega_c \), some vortex lines have broken through the
AB phase boundary. Thus there exists also a stable config-
uration in which vortices from the A phase continue
into the B phase, where they form a cluster in the center
[Fig. 2(b)]. The likely topology of the intersection is illustrated in Fig. 2(c). Thus, although the AB interface
is not directly monitored by the two NMR spectrometers,
the hydrodynamically stable states below and above \( \Omega_c \)
can only be the configurations in Fig. 2. By comparing
B-phase vortex-line creation in the presence and absence
of the AB interface, we conclude that the events in Fig. 3
originate from the AB phase boundary.

In the classical KH instability the interface between fluids with densities \( \rho_1 \) and \( \rho_2 \) becomes destabilized by inertial effects, which are normally balanced by gravity \( g \)
and surface tension \( \sigma \). The relative velocity \( |v_2 - v_1| \) acts as a drive. When it reaches a critical value [2] given by

\[
\frac{\rho_1 \rho_2}{\rho_1 + \rho_2} (v_2 - v_1)^2 = 2\sqrt{\frac{g}{\sigma}} F,
\]

where \( F = g(\rho_1 - \rho_2) \) is the gravitational restoring force, waves with wave vector \( k = \sqrt{F/\sigma} \) are created on the interface.

This approach can be generalized for superfluids in
terms of two-fluid hydrodynamics [15]. For superfluid
The dissipative, damping modifications of the interface lead to a negative free energy in textbooks [1], or from the thermodynamics when perturbation instability can be derived from the dynamics of small before at a drive given by

\[ \frac{\partial}{\partial t} \chi^B(T, \rho) = \Omega \chi^B(T, \rho), \]

where \( \Omega \) is the rotation rate. Instead of gravity, the restoring force is now produced by the magnetic barrier \( H_b(z) \), owing to the difference in the susceptibilities \( \chi_A \) and \( \chi_B(T, H) \):

\[ F = \Omega \chi_A^B(T, H) \nabla (H_b^2), \]

The restoring force can thus be calculated as a function of temperature and current in the barrier solenoid. The motion of the \( AB \) interface with respect to the normal component is subject to a finite damping [16]. Since the initial state is non-dissipative, damping modifies, but does not explicitly appear in, the stability condition [17]. The onset of the instability can be derived from the dynamics of small amplitude perturbations, as Eq. (1) is generally derived in textbooks [1], or from the thermodynamics when perturbations of the interface lead to a negative free energy in the rotating frame. The mode for which the interface first becomes unstable has the same wave vector \( k = \sqrt{F/\sigma} \) as before at a drive given by

\[ \frac{1}{2} \rho_A (v_{aA} - v_A)^2 + \frac{1}{2} \rho_B (v_{aB} - v_A)^2 = \sqrt{\sigma F}, \]

where \( \rho_A, \rho_B \) and \( v_A, v_B \) are the densities and velocities of superfluid and normal components.

In Fig. 4 we plot the measured \( \Omega_c \) for the first appearance of \( B \)-phase vortex lines in the presence of the \( AB \) boundary as a function of temperature, while \( \Omega \) is slowly increased (\( \Omega = 5 \times 10^{-4} \text{ rad/s}^2 \)). The current in the barrier solenoid is constant but different for the three sets of data. It controls the magnetic force \( F \) which is largely responsible for the shape of the curves: As a function of decreasing temperature the curves start at \( T_{AB}(H_{AB} = 35 \text{ mT}) \) and end at the temperature at which \( H_{AB} \) equals the maximum value of the barrier field \( H_b \) at the given current. The steep slopes at both ends of the curves are caused by the rapidly changing value of \( \nabla H_b \). The solid curves represent Eq. (2) if one sets \( \Omega_c = |v_B - v_n|/R \) and \( v_A = v_n \). No fitting parameters are used. The values for \( \sigma(T), \chi_A = \chi_B(T, H), \rho_B(T, H) \) are obtained from accepted references [8,18–20], while \( H_b \) and \( \nabla H_b \) apply for the profile \( H_b(z) \) at the location of the \( AB \) boundary: \( H_{AB}(T) = H_b(z) \).
remain valid even in the low temperature limit, when \( p_n \to 0 \) and \( p_{SA} \approx p_{SB} \to p \), if one waits long enough.

A further comparison to Eq. (2) is obtained from the number \( \delta N \) of new \( B \)-phase vortex lines after an instability event. In Fig. 3 it is seen that \( \delta N = 10 \). At 0.77 \( T_c \) (with \( u_c = \Omega, R = 0.39 \text{ cm/s} \) at 29.0 bars) the measured average for more than 100 events is \( \delta N = 11 \). We interpret \( \delta N \) to correspond to the number of circulation quanta \( \kappa = h/2m \), which fit in one corrugation of the interface mode of size \( \lambda_k = 2 = \pi/\kappa \). If we set \( v_{AB} = 0 \), then there are in solid-body rotation \( N = \pi R^2 \Omega/\kappa \) vortex lines in the \( A \) phase which all flare out into the lateral sample boundary at the \( AB \) interface [Fig. 2(a)]. Measured along the perimeter of the sample there are \( \Omega, R/\kappa \) circulation quanta per unit length and thus in one corrugation \( \delta N = (\kappa u_c)/(kk) \). From Eq. (2) this is seen to be \( \delta N = (2\pi \sigma)/(ku_c \rho_{SB}) \), which at 0.77 \( T_c \) gives nine vortex lines in agreement with the measured number.

Earlier rotating experiments have all been performed in the absence of a strong magnetic field gradient. It is then not possible to localize the \( AB \) interface at a fixed position. Instead one can record the motion of the interface through the sample in a slow \( A \to B \) transition during cooling at constant \( \Omega \) [21]. An \( A \)-phase vortex layer and tangential vortex-free \( B \)-phase counterflow are moving with the interface. Depending on the interface velocity, part of the \( A \)-phase vortex lines are swept with the interface to the sample boundary and are annihilated there, while part of the lines break through the boundary into the \( B \) phase. As a result the total circulation is found to be reduced after the \( A \to B \) transition. As the \( AB \) interface is in motion, the state is dissipative and Eq. (2) is not directly applicable. Nevertheless, applying dimensional arguments we find that the time for the KH instability to develop is \( \sim (ku_c)^{-1} \), which from Eq. (2) is seen to be \( \sim (\sigma / (\rho_{SB} t_{AB})) \). To form vortices, this time must be less than the transition time \( t_{AB} \) for the \( A \to B \) interface to sweep through the entire sample. This leads to an effective critical velocity \( u_c \sim [(\sigma / (\rho_{SB} t_{AB}))^{1/3}] \). It fits the measurements in Ref. [21] in the limit of slow interface velocities and shows that the \( AB \) interface becomes stiffer and less permeable at increasing velocity. At suitable velocities the removal of vortices from the system in front of the moving \( AB \) phase boundary becomes quite effective. Similar mechanisms have been discussed for sweeping away the abundantly created magnetic monopoles from the early Universe [22].

To summarize, in uniform rotation the magnetically stabilized \( AB \) phase boundary has been found to give rise to a state with tangential superfluid shear flow. The reason for this is a large energy barrier which prevents the nucleation of point and line defects on the small length scale \( \xi \). This situation evolves into the first example of a shear-flow instability where the initial state is nondissipative. The outcome of the instability is the transmission of a burst of vorticity across the interface. The instability creates corrugations in the interface and its vortex coating, upon which the vortex lines in the deepest trough are pulled by the Magnus force through the boundary. The process represents a new intrinsic mechanism of defect formation, applicable to interfaces between superfluids or to the free surface of a superfluid [23].

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[12] With only $B$ phase in the sample, $v_{AB}$ is determined by single-vortex formation at the sample boundary or by the leakage of vortices through the bottom orifice. Both processes differ in appearance from the events in Fig. 3.

[13] In this experiment the $A$-phase critical velocity is $v_{AB} = (0.36 \pm 0.06)$ mm/s.


