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NEW RESEARCH ON THE PROPERTIES OF LIQUID HELIUM

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1. RESEARCH ON He^4

THE superfluidity of the liquid He^4 has been the object of most thorough theoretical and experimental research from the instant of its discovery by P. L. Kapitza [1] in 1938. At the present time it can be assumed that the essential aspects of the behavior of a superfluid liquid have been sufficiently well studied. This has been considerably aided by the fact that as early as 1941 L. D. Landau developed a consistent theory of superfluidity [2], and all subsequent work was done with a clear understanding of the nature of the phenomenon. (As is known, the situation was entirely different in the case of superconductivity, the microscopic theory of which was developed only many years after the phenomenon was experimentally observed.) If we add also the fact that only a single object—liquid He^4—is presently available for research on superfluidity, it is quite natural to assume that this field of research has been completely exhausted. Nonetheless, quite a few theoretical and experimental investigations were carried out in recent years and have added to our knowledge of superfluidity and led to the formulation of new interesting problems. The purpose of the present review is to report on some of these results, published in 1964-1965. No attempt will be made here at a complete exposition, and we merely confine ourselves to a brief analysis of the papers that in our opinion are of the greatest interest. We shall not describe experiments connected with the so-called Josephson effect in helium, since these can hardly be discussed without reference to the corresponding phenomenon in superconductors. We shall assume that the reader is familiar with the main premises of the theory of superfluidity and refer him for a background to previously published review papers [3,4].

As is well known, the description of the properties of superfluid helium is based on knowledge of its energy spectrum. From data on the specific heat and on the temperature dependence of the density of the normal component we can determine the spectrum at sufficiently low energies. It turns out that in helium there are excitations of two kinds. One comprises sound quanta—phonons—and the dependence of the energy on the momentum is of the form

$$\varepsilon(p) = pc \tag{1}$$

(c—the velocity of sound in helium), and the other constitutes rotons, for which

$$\varepsilon = \Delta + \frac{(p-p_0)^2}{2\mu}, \tag{2}$$

where $\Delta/k = 8.6°K$, $\mu = 1 \times 10^{-24}$ g, $p_0/\hbar = 1.9$ Å^{-1}, and k is Boltzmann's constant. Landau, who introduced the concepts of rotons and phonons, proposed that they pertain to different parts of the same energy curve, i.e., that the curves (1) and (2) go continuously into each other at $p \sim p_0/2$. Experimental spectra obtained in 1957-1961 with the aid of inelastic scattering of slow neutrons [5,6] have confirmed this prediction and have made it possible to establish the form of the $\epsilon(p)$ curve for all $p \lesssim p_0$. Of special interest is, however, the question of subsequent behavior of the curve, particularly the way it terminates. The point is that with increasing momentum the spectral curve can reach a certain threshold point, above which the elementary excitation becomes unstable. In other words, starting with this point, the conservation laws admit of the de-

cay of the excitation into two or more excitations with lower energies and momenta. Thresholds of different types, differing in the properties of the excitations produced upon decay [7], are theoretically possible. The character of the spectrum at the point of decay can be explained theoretically in general form, but it is impossible to indicate on the basis of purely theoretical considerations precisely which type of threshold is actually realized in the spectrum of helium. Experimental data on the behavior of the spectrum has indicated that with increasing momentum the $\epsilon(p)$ curve should probably reach an energy value 2Δ at some momentum value $p_c < 2p_0$. It is clear that starting with this point the excitation can break up into two excitations, each with energy Δ. The excitations produced upon decay are emitted at a certain angle and the sum of their momenta is equal in absolute magnitude to p_c. According to the theory, the spectral curve should then terminate at the point 2Δ, the dispersion law near this point being

$$2\Delta - \varepsilon = \alpha e^{-\frac{\beta}{p_c - p}}, \tag{3}$$

where α and β are certain constants. The $\epsilon(p)$ curve is thus a horizontal tangent of infinite order at the point $\epsilon = 2\Delta$. The intensity of the neutron scattering accompanied by production of an excitation of energy $\epsilon \approx 2\Delta$ vanishes when $\epsilon \to 2\Delta$ in accordance with the law

$$(2\Delta - \varepsilon)\ln^2(2\Delta - \varepsilon). \tag{4}$$

The behavior of $\epsilon(p)$ when $p > p_0$ was experimentally investigated by Woods [8]. (Certain preliminary data are contained also in [9].) The experiment consisted of irradiating a vessel with helium at a temperature 1.6°K with a monoenergetic neutron beam. The scattering of the neutrons, accompanied by production of a single excitation in the helium, was manifested by a peak on the energy distribution of the neutrons scattered through a given angle. The position of the peak determined the energy of the produced excitation.

Knowing the initial neutron momentum p_1, the final momentum p_2, and the scattering angle ϑ, we can use the formula

$$p^2 = |\mathbf{p}_2 - \mathbf{p}_1|^2 = p_1^2 + p_2^2 - 2p_1p_2\cos\vartheta$$

to determine the momentum of the excitation p. The results obtained are shown in Fig. 1. The solid curve shows the function $\epsilon(p)$ in accordance with the data of [9], the points show the results of measurements with neutrons of wavelength $\lambda = 2.77$ Å, and the triangles—with wavelengths 2.48 Å. The abscissas are the momenta of the excitations (in $Å^{-1}$) and the ordinates the energies (in °K). We see from Fig. 1 that in accordance with the theory the $\epsilon(p)$ curve goes into a horizontal tangent when $\epsilon \approx 2\Delta$. The intensity of the scattering with creation of a single excitation as $\epsilon \to 2\Delta$ decreases, as in qualitative agreement with (4), and the probability of creation of an excitation with $p/\hbar = 3.35$ $Å^{-1}$ is approximately 1% of the probability of creation of an excitation with $p = p_0$. It is also seen from the figure that the experimental points rise somewhat above the $\epsilon = 2\Delta$ line. It is hardly expected from theory that the curve $\epsilon(p)$ with $p < 2p_0$ would be continued in any form in the region $\epsilon > 2\Delta$. It must be borne in mind, however, that the probability of creation of two excitations with $\epsilon \approx \Delta$ can have at $|p + p'| \approx p_c$ a spread-out maximum, which can be regarded as a peak corresponding to creation of one excitation. We note also that the excess of energy over 2Δ has in fact the same order of magnitude as the observational error, and that a more detailed discussion of this question is for the time being premature. Figure 1 contains also a point at $p \approx 3.6$ $Å^{-1}$. The author of [8] is not completely certain of the reliability of this result, since the intensity of the corresponding peak in the distribution of the scattered neutrons is very low. If, however, the existence of this peak is confirmed, this will signify the observation of a rather interesting phenomenon. The point is that the spectral curve, which is made discontinuous by the decay at $\epsilon = 2\Delta$, can be again continued in the region $p > 2p_0$. Indeed, the excitation with $p > 2p_0$, $\epsilon \approx 2\Delta$ cannot break up into two excitations with $\epsilon \approx \Delta$, since these excitations

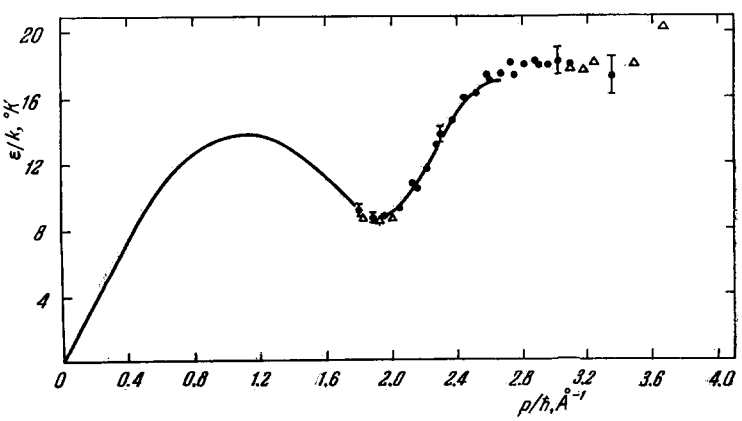

FIG. 1.

cannot carry away a momentum larger than $2p_0$. (This circumstance was noted earlier in [10].) Furthermore, it is clear beforehand that excitations with sufficiently large p actually do exist. These are vortex rings. As is well known (see, for example, [11]) vortex filaments with quantized circulation can exist in a superfluid liquid. These are singular lines near which the superfluid part of the liquid rotates like $\hbar/mr$, where r is the distance from the filament and m is the mass of the liquid atom. These filaments can form closed rings. Such a ring has energy and momentum

$$\varepsilon = 2\pi^2\rho_s\frac{\hbar^2}{m^2}\ln\left(\frac{R}{a}\right)R, \quad p = 2\pi^2R^2\rho_s\frac{\hbar}{m}$$

(R is the radius of the ring, ρ_s is the superfluid component of the liquid, and a is of the order of atomic length; it is assumed that $R \gg a$). We confine ourselves to the case of circular rings; it is readily seen that only such rings are stable. Rings of any other form go over into circular rings with emission of sound. Thus, a vortex ring can be regarded as an elementary excitation with dispersion law

$$\varepsilon(p) = \sqrt{2}\pi\sqrt{\rho_s}\left(\frac{\hbar}{m}\right)^{3/2}\ln\left(\sqrt{\frac{pm}{2\pi^2\rho_s a}}\right)\sqrt{p}. \qquad (5)$$

This formula is suitable, of course, only for sufficiently large p. This, however, raises the question: how does this branch of the spectrum behave with decreasing p, and in particular, where does it begin? One is tempted to assume that the point at $p/\hbar \approx 3.6$ Å^{-1} marked in Fig. 1 is the start of the vortex spectrum, that is, that it corresponds to the creation by the neutron of a vortex with minimal possible dimension. We emphasize once more, however, that there are no experimental grounds for making such a statement at present.

Much more effective at present is the study of the spectrum of vortex rings by an entirely different method, based on the creation of such rings by ions moving in helium. (We shall not stop to discuss model representations concerning the structure of matter in the direct vicinity of the ion, since these questions are dealt with in a recently published article [12].) We have in mind here the work by Rayfield and Reif [13], who investigated the motion of ions in helium in the temperature interval from 0.3 to 0.6°K. The ions, like all other impurities in the helium, can be regarded as elementary excitations with definite energy and momentum. Rayfield and Reif determined the dispersion law of such excitations. It turns out that in the energy interval from 1.5 to 45 eV this dispersion law coincides exactly with the dispersion law of the vortex rings (5). This means that in this energy interval the ion in the helium is closely linked with the vortex. At a certain energy the ion creates a vortex ring and subsequently moves together with it. The energy and the momentum of the compound excitation produced in this manner actually coincides with the energy and momentum of the ring, and the charge, naturally, is equal to the charge of the ion.

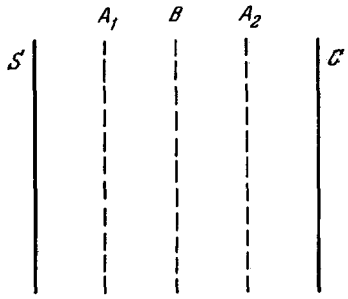

FIG. 2.

The principal scheme of these properties is exceedingly simple (Fig. 2). The ions are produced near an electrode S by ionizing the helium with α particles from radioactive Po^{210}. Between the grid A_1 and S the ions are accelerated to a definite energy eV_1. The remainder of the setup serves to measure the velocity of the excitations having the same energy. The grid A_1 and A_2 has the same potential. Between C and A_2 there is applied a decelerating potential $V_2 \approx V_1$, which prevents the ions from falling on the electrode C. Between the grid B and the grids A_1 and A_2 there is applied a small potential $V_B \ll V_1$, which reverses its sign periodically with a period τ. If the time of flight of the ion between the grids A_1 and B is equal to $\tau/2$, then these ions will be accelerated by the potential V_B, if the initial phase of the motion is properly chosen, both between A_1 and B and between B and A_2, since the potential reverses sign precisely when the ion passes through the grid B. As a result, the dependence of the current to C on τ has a resonant maximum at

$$\tau = \frac{2l}{V}$$

(l is the distance between A_1 and B), making it possible to measure the velocity of the perturbation V. The values of U obtained in this manner are plotted in Fig. 3 against the energy $E = eV_1$. The solid curve in this figure is the velocity of the vortex rings as a function of their energy. This velocity can be calculated from a formula which follows directly from (5):

$$V = \frac{\partial\varepsilon}{\partial p} = \sqrt{\frac{\rho_s}{2}}\pi\left(\frac{\hbar}{m}\right)^{3/2}\ln\left(\sqrt{\frac{pm}{2\pi^2\rho_s a}}\right)p^{-1/2}. \qquad (6)$$

(The formula given in [13] differs from (6) in an inessential factor under the logarithm sign.) The splendid agreement between the experimental points on the curve leaves no doubt that the interpretation of the phenomenon proposed by the authors is correct. We note that when the energy changed from 1.5 to 45 eV the radius of the vortex ring changed from 5×10^{-6} to 10^{-4} cm. No difference was observed in the behavior of the positive and negative charges in these experiments.

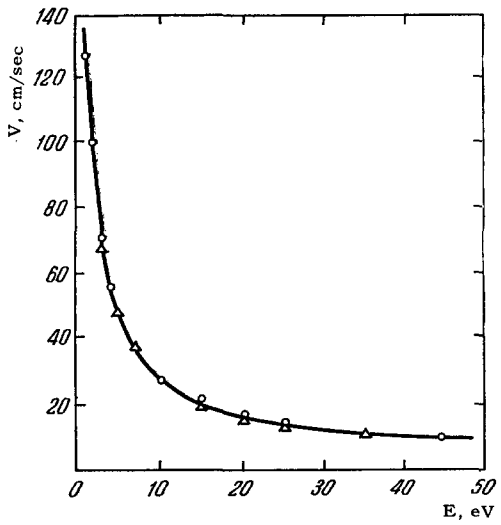

FIG. 3.

A very interesting experiment, also connected with the interaction between vortices in helium and ions, was carried out by Douglass [14]. This case pertains to capture of negative ions by straight-line filaments present in a rotating superfluid helium. The instrument in which the experiment was carried out is shown in Fig. 4. It consisted of a vessel in which the helium was placed, made up of several metallic electrodes separated by teflon liners. The vessel could rotate about a vertical axis. The maximum angular velocity of rotation was 45 rad/min. The cylindrical electrode S was covered with an α-active source (Po^{210}), which produced ionization in a thin layer near S. The experiment that proved most convincingly the existence of negative ions captured by the vortex filaments was performed as follows. First, a positive potential (relative to the source S) was applied to the grid G and to the central electrode C. This produced a current of negative ions to the electrode C and space charge in the space between G and C. Then a negative potential relative to S was applied to the grid G, stopping the flow of new ions from the source. Since a horizontal electric field (intensity ~ 20 V/cm) existed in the space between G and C, all the free ions were drawn off from the helium to C after a certain transient

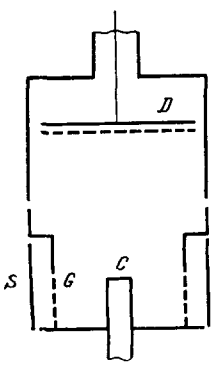

FIG. 4.

time (~1 sec), and the current stopped flowing. If, however, the ions were then allowed to flow to the electrode D, then following a certain time an ion current started to flow to D, in spite of the apparent absence of ions in the interelectrode space. This current appeared only in the rotating helium and did not appear when the helium was at rest.* The only explanation for this phenomenon is that some of the ions of the rotating helium are captured by the vortex filaments that are parallel to the rotation axis. These ions cannot move transversely to the filaments and are not removed from the space by the horizontal electric field. On the other hand, as they move along the filaments, they strike the electrode D. The time that the captured ions stay on the filament is, of course, finite. This time can be estimated by measuring the dependence of the number of ions reaching D on the time elapsed between the removal of the free ions and the instant when the ions are started towards D. In the temperature interval between 1.60 and 1.72°K this time is described by the formula

$$\tau \sim \exp\left(\frac{\varepsilon_0}{kT}\right),$$

where $\epsilon_0 = 0.012$ eV. The energy ϵ_0 can be regarded as the depth of the potential well in which the ion captured by the filament is situated. As should be the case, the number of captured ions is proportional to the angular velocity of rotation, that is, to the number of vortex filaments. The mobility of the captured ions along the vortex filament turns out to be, in accordance with preliminary estimates, about one third the mobility of the free ions. Therefore, a plot of the current to D against the time, obtained without first removing the free ions, shows two regions where the current increases. One corresponds to the instant when the free ions arrive at D, and the other to the instant of arrival of the captured ions.

The foregoing phenomena are observed only with negative ions, apparently because a rather large cavity is produced around a negative ion (for more details see [12]). The negative ion turns out to be, as it were, "lighter than the liquid" and the centripetal forces drag it toward the center of the filament. Another study of rotating helium was made by Reppy and Depatie [16], who observed undamped flow of the superfluid part of helium in an annular channel. The feasibility of such a motion is obvious, since the superfluid motion is not connected with energy dissipation. The work is nonetheless of interest, since, unlike earlier investigations [17,18], it deals with the dependence of the angular momentum of the superfluid part on the temperature.

The experiment consisted of rotating a helium-filled vessel placed between two cylinders at a velocity

*The capture of ions by vortex filaments was first observed in [15].

larger than critical. (We disregard specific experimental conditions, which are immaterial to us.) The vessel was then slowly stopped. The superfluid part of the liquid, on the other hand, continued to rotate. The vessel was then left free to rotate and the entire system was heated to a temperature above the temperature of the λ transition, so that the liquid and the vessel could start rotating as a unit at a common angular momentum equal to the momentum of the superfluid part prior to the heating, making it possible to measure the latter. On the other hand, when the vessel was secured and the temperature varied slowly, then, as can be readily understood, the angular momentum of the superfluid part would change. Indeed, let us consider for simplicity a superfluid liquid in a narrow annular channel. Then the angular momentum of the liquid will be (per unit height of the ring)

$$L = M_s R V_s = \varrho_s \cdot 2\pi R V_s d,$$

where M_S is the total mass of the superfluid part, R the radius of the ring, d its thickness, and $2\pi RV_S$ the circulation of the velocity along the contour of the ring, which is conserved for an ideal liquid. (In a superfluid liquid $\oint V_S dr = 2\pi hn/m$, where n is an integer which, of course, can change as the external conditions are continuously varied.) It turns out as a result that the angular momentum of the superfluid part changes in proportion to ρ_S. This, of course, does not contradict the law of angular-momentum conservation; the momentum is transferred to the normal part of the liquid, and from it to the stationary and secured vessel. The measurement of the momentum of the superfluid part after the change in temperature in the manner given above has confirmed the proportionality of L to ρ_S. It must be emphasized, however, that this result is perfectly obvious from the point of view of the theory, although the possibility of actual realization of such an experiment is of undisputed interest.

Measurement of the density of the superfluid helium has been the subject of a recently published paper by Andronikashvili and Tsakadze[19]. The authors used a sensitive pycnometer to measure changes in the helium density with accuracy $\Delta\rho/\rho \sim 10^{-5}$. The measurements have shown that the density of the liquid is changed by the rotation much more than if we were dealing with contripetal compression. The change in the density at an angular velocity 30 rad/sec and a temperature 1.74°K is $\Delta\rho/\rho \approx 3 \times 10^{-4}$, which is several dozen times larger than compression calculated from the usual values of compressibility and the well known formula for centrifugal pressure. An impression is gained that the rotating He II has a certain anomalous compressibility. This, of course, is very difficult to explain on the basis of the existing concepts. Since there is no such effect in helium above the λ point, the density is discontinuous at the λ point, so that the transition in the rotating helium is of first order. No noticeable shift of the transition point is observed in this case.

2. RESEARCH ON LIQUID He^3

The second of the quantum liquids known to us is liquid He^3. The excitations in this liquid obey the Fermi statistics, and its properties in the temperature intervals investigated so far have nothing in common with the properties of He^4. In particular, He^3 is not superfluid. To the contrary, its viscosity increases sharply with increasing temperatures. This, however, raises the following question, which is perhaps most interesting of all those concerning this liquid: Does He^3 become superfluid at lower temperatures? After the microscopic theory of superconductivity was developed it became obvious that in principle a transition of a Fermi system into the superfluid state is possible. It follows from this theory that a Fermi gas whose particles are attracted by arbitrarily small forces becomes superfluid at sufficiently low temperatures, in analogy with electrons in a superconducting metal. Consequently, immediately after the publication of the papers on superconductivity theory, the hypothesis was advanced that such a situation can be realized in He^3 [20-22]. The most important question, however, was whether the attraction between the elementary excitations necessary for the transition into the superfluid state can exist in He^3. (As is well known, in a superconducting metal such an attraction is effected by exchange of virtual phonons, something which in no case pertains to He^3.) Moreover, it is quite obvious that the elementary excitations in He^3 repel each other in principle. Otherwise, this liquid would be superfluid even at temperatures of the order of the Fermi temperature, since the interaction between the excitations in a Fermi liquid are far from weak.

A more detailed investigation has shown, however, that for a transition into the superfluid state, it is sufficient that the attraction between the excitations exist for at least one value of the relative angular momentum of the interacting excitations. It also turned out that attraction forces, similar to Van der Waals forces between remote atoms, exist at sufficiently large momenta, between the excitations in any uncharged Fermi liquid. The strong repulsion between excitations at small distances can change the magnitude but not the sign of such a long-range interaction[21]. It was proved as a result that He^3 will actually become superfluid at a sufficiently low temperature. An estimate of this temperature is a very difficult theoretical problem. If we assume that the interaction between the excitations is precisely the same as between isolated helium atoms, then the transition temperature is approximately 8×10^{-3} °K [21-23]. We can state, however, that the presence of neighboring atoms greatly influences the magnitude of the interaction. Thus, attraction between remote excitations decreases in comparison with attraction between free atoms by a factor

$$Z = \left(\frac{m^* c^2}{m c_0^2}\right)^2$$

where m* is the effective mass of the excitations in

He^3, m the mass of the He^3 atom, c^{-2} the compressibility of the liquid He^3, and c_0^{-2} the compressibility of an ideal Fermi gas with mass m and density equal to the density of liquid helium. The degree of attenuation of the interaction Z for liquid He^3 turned out to be unexpectedly large, $Z \approx 65$. If we assume that at a relative excitation momentum $l = 2$ the interaction is weakened by the same factor, then the temperature of the transition turns out to be much lower—of the order 2×10^{-4} °K. (It can be assumed that the attraction between the excitations first appears at exactly $l = 2$; this value of the momentum is responsible for the superfluidity.) The strong discrepancy between the foregoing estimates, and their approximate character, show that the temperature T_c of transition into the superfluid state can be obtained reliably only by experiment. Unfortunately, the experimental situation is far from clear at present. Two groups of investigators--Peshkov in the USSR [24] and Abel, Anderson, Black, and Wheatley in the USA [25] —have succeeded in carrying out experiments with He^3 at very low temperatures, down to 4×10^{-3} °K. The cooling was by adiabatic demagnetization of a paramagnetic salt—cerium magnesium nitrate ($CeMg(NO_3)_6$) with Curie temperature 3.2×10^{-3} °K. Peshkov, after measuring the specific heat of a mixture of the salt with liquid He^3 and subtracting the specific heat of the same installation without the helium, observed a jump-like anomaly of the curve of the specific heat of helium at a temperature 5.5×10^{-3} °K. It is natural to assume that this is the jump of specific heat at the point of phase transition into the superfluid phase. The authors of [25], however, observed no such jump. Their measurements of the coefficient of self-diffusion of helium atoms and of the nuclear magnetic susceptibility have likewise led to no observation of any anomalies down to 3.6×10^{-3} °K. The reason for such a discrepancy is not clear at present. The experimental apparatus used by the different workers, although employing similar ideas, differed in many essential details. Nor do we know whether the temperature scales employed by them coincide. It is clear that the question can be solved only by further research.

Considerable interest has been evinced in recent years by investigations of the properties of He^3 in a temperature region where it is certainly not superfluid. At such temperatures, He^3 is a Fermi liquid which can be described by Landau's theory of Fermi liquids (see the review [26]). In particular, its specific heat should decrease with decreasing temperature like $C = AT$. This law is satisfied experimentally, although with not very high accuracy. In our opinion this is probably connected with the insufficiently low temperatures of the experiments, and possibly with the insufficient accuracy of the temperature scales in this region. Another conclusion was deduced by Anderson. On the basis of analysis of the specific-heat data published in [27-30], he made the statement that the specific heat of He^3 actually satisfies the relation [31]

$$C = AT \ln T \tag{7}$$

and that consequently the Fermi-liquid theory is in need of revision. He considers a possible way towards such a revision to be allowance for the interaction between excitations by exchange of zero-sound quanta (for more details, see [26]). On the basis of this idea, Balian and Fredkin developed a theory, according to which the specific heat of He^3 actually obeys (7) and the velocity of the elementary excitations on the Fermi surface vanishes [32]. In our opinion, however, the assumptions based on the theory are intrinsically contradictory. The main assumption of the authors is that the amplitude of the scattering of the excitations by one another, which describes their interaction, has for small k and ω (ω—energy, k—momentum transferred by one excitation to another) the form

$$\Gamma = \frac{\alpha}{\omega^2 - k^2c^2},$$

where c is the velocity of zero sound. Yet it can be shown by using Landau's results [33] that for small k the functions ω and Γ depend only on the ratio ω/k. This means that near the pole Γ should actually be of the form

$$\Gamma = \frac{\alpha' k^2}{\omega^2 - k^2c^2}.$$

However, when the amplitude Γ has this form, allowance for the exchange of zero-sound quanta leads to no change whatever in the theory.

[1] L. P. Kapitza, DAN SSSR 18 (1), 29 (1948).
[2] L. D. Landau, JETP 11, 592 (1941).
[3] I. M. Khalatnikov, UFN 59, 69 and 60, 673 (1956).
[4] I. M. Khalatnikov, Vvedenie v teoriyu sverkhtekuchesti (Introduction to the Theory of Superfluidity), Nauka, 1965.
[5] H. Palevsky, K. Otnes, and K. E. Larson, Phys. Rev. 112, 11 (1959).
[6] J. Yarnell, G. Arnold, P. Bendt, and E. Kerr, Phys. Rev. 113, 1379 (1959).
[7] L. P. Pitaevskiĭ, JETP 39, 1168 (1959), Soviet Phys. JETP 9, 830 (1959).
[8] A. Woods, Preprint (1964).
[9] D. Henshaw and A. Woods, Phys. Rev. 121, 1266 (1961).
[10] H. Jackson and E. Feenberg, Rev. Mod. Phys. 34, 686 (1962).
[11] É. L. Andronikashvili, Yu. G. Mamaladze, S. P. Matinyan, and D. S. Tsakadze, UFN 73, 3 (1961), Soviet Phys. Uspekhi 4, 1 (1961).
[12] R. G. Arkhipov, UFN 88, 185 (1966), Soviet Phys. Uspekhi 9, 174 (1966).
[13] G. Rayfield and F. Reif, Phys. Rev. Letts. 11, 305 (1963).

[14] R. Douglass, Phys. Rev. Letts. **13**, 791 (1964).
[15] G. Careri, W. McCormick, and F. Scaramuzzi, Phys. Letts. **1**, 61 (1962).
[16] J. Reppy and D. Depatie, Phys. Rev. Letts. **12**, 187 (1964).
[17] H. Hall, Phil. Trans. Roy. Soc. (London) **A250**, 980 (1957).
[18] W. Vinen, Proc. Roy. Soc. **A260**, 218 (1961).
[19] É. L. Andronikashvili and D. S. Tsakadze, JETP Letters **2**, 278 (1965), transl. p. 177.
[20] K. Brueckner, T. Soda, P. Anderson, and P. Morel, Phys. Rev. **118**, 1442 (1960).
[21] L. P. Pitaevskiĭ, JETP **37**, 1794 (1959), Soviet Phys. JETP **10**, 1267 (1960).
[22] V. Emery and A. Sessler, Phys. Rev. **119**, 43 (1960).
[23] L. P. Gor'kov and L. P. Pitaevskiĭ, JETP **42**, 600 (1962), Soviet Phys. JETP **15**, 417 (1962).
[24] V. P. Peshkov, JETP **46**, 1510 (1964), Soviet Phys. JETP **19**, 1023 (1964).
[25] W. Abel, A. Anderson, W. Black, and J. Wheatley, Phys. Rev. Letts. **14**, 129 (1965).
[26] A. A. Abrikosov and I. M. Khalatnikov, UFN **66**, 177 (1958), Soviet Phys. Uspekhi **1**, 68 (1959).
[27] W. Abel, A. Anderson, W. Black, and J. Wheatley, Physics (1965).
[28] B. Abraham, D. W. Osborne, and B. Weinstok, Phys. Rev. **98**, 499 (1965).
[29] A. Anderson, W. Recse, and J. Wheatley, Phys. Rev. **130**, 495 (1963).
[30] M. Strongin, G. Zimmerman, and H. Fairbank, Phys. Rev. **128**, 1483 (1962).
[31] P. W. Anderson, Preprint (1965).
[32] R. Balian and D. Fredkin, Phys. Rev. Letts. **15**, 408 (1965).
[33] L. D. Landau, JETP **35**, 97 (1958), Soviet Phys. JETP **8**, 70 (1959).

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