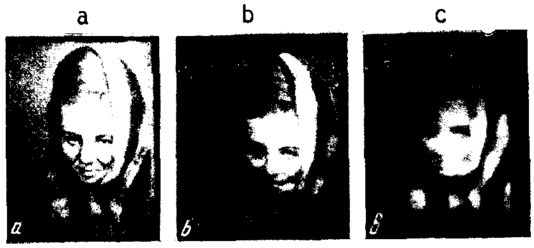


DRSh-1000 lamp using two light filters (ZhZS-9 and PS-7), and (c) in the light from an incandescent lamp through a glass light filter (KS-13). The dimensions of the hologram correspond to a 24 x 36 mm frame of a miniature camera.

This raises the question whether it is also possible, by foregoing the redundant information in the hologram, to use a light source of equally broad spectral composition to obtain a hologram on an area corresponding to the value of D in formula (1).

For each point of the object this is indeed so. In this case the problem is fully equivalent to the problem of determining the number of interference fringes seen in the light from a given source. On going over to the entire aggregate of points of the object, it is necessary to take into account the fact that zero phase differences between the reference beam and the beam from a given point of the object do not occur in identical locations on the hologram plane. Therefore broadening of the spectrum leads not to a narrowing of the region where the interference pattern (hologram) is sharp, but to a decrease in its contrast, which eventually vanishes at a certain spectral width. The permissible spectral width depends on the dimensions of the object and on the concrete optical scheme used to obtain the hologram.

Thus, a light source which is perfectly adequate for the reconstruction of an image of satisfactory quality may turn out to be utterly unsuitable for the production of a hologram. At the same time, there may exist a large number of problems and technical solutions in which the loss of information contained in the hologram is offset by the simplicity of reconstruction of the hologram in ordinary light sources.

INSTABILITY OF FERMI SYSTEMS AND SPECIFIC HEAT OF LIQUID He^3

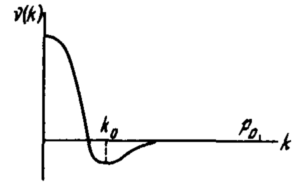
D. A. Kirzhnits and Yu. A. Nepomnyashchii
P. N. Lebedev Physics Institute, USSR Academy of Sciences
Submitted 23 May 1966
ZhETF Pis'ma 4, No. 3, 86-90, 1 August 1966

1. The disparity between low-temperature data on the specific heat of He^3 [1] and the predictions of the Fermi-liquid theory [2,4] has been attracting interest of late. We can attempt to explain this disparity [1)] by assuming that at some still-unattained temperature T_c the system experiences a second-order phase transition, as a result of which the specific heat has a peak of width ΔT near T_c. It follows from experimental data that $T_c < 0.01°$, $\Delta T \sim 1°$, and at any rate $\Delta T/T_c \gg 1$. There can therefore be no talk of a transition to the superfluid state that might be produced as a result of long-range attraction forces [5], since for such a transition [6,7] $\Delta T/T_c \sim (mT_c/p_0^2)^4 \ll 1$ (p_0 = limiting momentum, m = particle mass).

We propose in this note a possible explanation of the anomaly of the specific heat of

He^3, based on the fact that the long-range attraction forces are capable also of leading to a phase transition of an essentially different nature, with a value of $\Delta T/T_c$ which is assuredly larger. We refer to a rearrangement of the system not in the "particle-particle" channel, as in the case of superfluidity, but in the "particle-hole" channel; the system goes in this case into a unique spatially-inhomogeneous state. Thus the anomaly under discussion is, so to speak, a certain "percursor" of such a transition.

2. Let us consider the simplest model (see the figure). $\nu(k)$ is the Fourier transform of the pair interaction potential. Attraction takes place when $k \ll p_0$, with $\nu(0) > 0$ and $\int d^3k\nu(k) > 0$, which prevents the system from collapsing. It is assumed that the system is compressed, $\eta = p_0/k_0 \gg 1$, and that the ratio of the particle-pair interaction to the kinetic energy, $\alpha = (m|\nu(k_0)|k_0^3)/\pi^2 p_0^2$, is small; however, the ratio of the total interaction to the kinetic energy is $\alpha\eta^3 \equiv 1 + \gamma \gtrsim 1$.



We choose the Hartree approximation as that of zero order (the exchange term is small for compressed systems); the corresponding Green's function is determined by the equation

$$\left(i\frac{\partial}{\partial t} - \frac{\hat{p}^2}{2m} + i\int dx''V(x - x'')G_0(x, x'')\right)G_0(x, x') = \delta(x - x'). \tag{1}$$

When the foregoing conditions are satisfied, we can confine ourselves to polarization diagrams only [8]. Taking the translation-invariant solution of (1), we have for the vertex part

$$\Gamma(\vec{k}, \omega) = \left(\frac{1}{\nu(k)} - \Pi(\vec{k}, \omega)\right)^{-1}, \tag{2}$$

where $\Pi(\vec{k}, \omega) = -2i\int d^4pG_0(p)G_0(p + k)$. When $\gamma > 0$ and $T < T_c$ Eq. (2) has a pole at imaginary frequency, thus demonstrating the instability of the system in the "particle-hole" channel. T_c itself is determined, when $\gamma \ll 1$ and $k = k_0$, by the expression $(p_0k_0/2m)(\gamma/\ln 2 \ln 1/\gamma)$.

3. The heart of the matter is the instability (against small density variations) of the translation-invariant solution of (1), which now corresponds not to the minimum energy, but to a stationary point [9,10]. The modified Green's function which does not lead to an imaginary pole in (2) is none other than the stable solution of (1) corresponding to an energy minimum and describing an inhomogeneous density distribution.

We can arrive at the same results by starting from the analogy with the theory of superconductivity and by describing the condensation of correlated "particle-hole" pairs. The corresponding equation of Gor'kov [11], containing an extra "particle-hole" pair in the intermediate state, coincides exactly with (1).

4. When $\gamma \ll 1$ (or when $|T - T_c| \ll T_c$) a stable state of the system corresponds to weakly-modulated periodic density distribution. As in the case of electron motion in the lattice field, the excitation spectrum has a gap of a dielectric type; in our case, however, many Brillouin zones are contained inside the Fermi sphere.

We note that in the model under consideration the usual superfluidity is strongly suppressed and the modified vertex function has no imaginary pole at all in the "particle-particle" channel (at not too small values of γ).

5. In conclusion, let us estimate the value of $\Delta T/T_c$ (see Sec. 1). The expression for the thermodynamic potential is

$$\Phi - \Phi_0 = \frac{mp_0}{3\pi^2}\int dx\eta(x)\hat{O}\eta(x) + \dots, \tag{3}$$

where $\hat{O} = 1 - \nu(\hat{p})\sum_{\omega}\Pi(\hat{p}, \omega)$ and $\eta(x)$ is the deviation of the density from homogeneous. Estimating the role of the fluctuations by Ginzburg's method [6], we get

$$\Delta T/T_c \sim (k_0^2/mT_c)^2. \tag{4}$$

If this ratio is small, for which very small values of k_0 are required, then the foregoing calculation, which was carried out within the framework of the modified Hartree approximation, can be regarded as valid everywhere except in a small vicinity about T_c. A case of greater interest is when (4) is not small: this is precisely what should occur if the proposed explanation is correct [2)]. It is necessary here to go beyond the framework of the zeroth approximation and to take into account polarization diagrams describing the density fluctuations. Such a calculation is now under way. Even now, however, there is assurance of the correctness of the deduction that $\Delta T/T_c \gg 1$ is possible. The specific-heat anomaly at $T > T_c$ arises formally because the corresponding pole in γ has a residue that increases without limit when γ approaches zero or when T approaches T_c. We note, on the other hand, that the fluctuations connected with the "particle-hole" pair are in general much more significant than "superfluid" fluctuations; the energy of such a pair is relatively low (it does not contain the large term characteristic of the "particle-particle" pair and describing the kinetic energy of the pair as a whole).

A detailed exposition of the problems touched upon here will be published elsewhere.

The authors are deeply grateful to V. L. Ginzburg and the participants of the seminar under his direction for numerous useful discussions.

[1] W. Abel, A. Anderson, W. Black, and J. Wheatley, Physics 1, 337 (1965); Phys. Rev. Lett. 15, 875 (1965).

[2] P. W. Anderson, Physics 1 (1965).

[3] R. Balian and D. Fredkin, Phys. Rev. Lett. 15, 408 (1965).

[4] L. P. Pitaevskii, UFN 88, 490 (1966), Soviet Phys. Uspekhi 9, in press.

[5] L. P. Pitaevskii, JETP 37, 1794 (1959), Soviet Phys. JETP 10, 1267 (1960).

[6] V. L. Ginzburg, FTT 2, 2031 (1960), Soviet Phys. Solid State 2, 1824 (1961).

[7] L. D. Landau and E. M. Lifshitz, Statisticheskaya fizika (Statistical Physics), Sec. 80, Nauka, 1964.

[8] D. A. Kirzhnits, Polevye metody teorii mnogikh chastits (Field Methods of Many-particle Theory), Gosatomizdat, 1963.

[9] A. Overhauser, Phys. Rev. Lett. 4, 415 (1960).

[10] D. J. Thouless, The Quantum Mechanics of Many-body Systems, Academic Press, 1961.

[11] A. A. Abrikosov, L. P. Gor'kov, and I. E. Dzyaloshinskii, Metody kvantovoi teorii polya

v statisticheskoi fizike (Methods of Quantum Field Theory in Statistical Physics), 1962.

1) The explanation proposed in [3] has been refuted in [4].

2) By substituting the experimental data in (4) (see Sec. 1), we obtain the reasonable estimate $k_0/p_0 \sim 0.1 - 0.2$.

ON THE USE OF AN ELECTRON SYNCHROTRON AS A MASER

A. A. Sokolov and I. M. Ternov
Physics Department, Moscow State University
Submitted 23 May 1966
ZhETF Pis'ma 4, No. 3, 90-92, 1 August 1966

The question of obtaining stimulated emission from electrons moving in a magnetic field [1,2] or crossed (magnetic and electrostatic) fields [3] has been discussed many times recently. The papers cited, however, consider the nonrelativistic or weakly-relativistic case, involving emission at the fundamental harmonic, which can be used to obtain radiation at a wavelength of the order of the orbit radius. We wish to show that in the relativistic case it is possible, with an electron moving in a magnetic field only, to make the stimulated emission prevail over absorption in a definite band of high harmonics corresponding to a certain resonance.

If the incident electromagnetic wave of frequency ω is linearly polarized (the vector of electric intensity $\mathcal{E}$ lies in the plane of the orbit) and propagates perpendicular to a constant magnetic field $\vec{H}$, causing cyclic motion of the electron, then we obtain for the total power of the stimulated emission and absorption of a harmonic ν the following formula [4]:

$$W_\nu = -\frac{4e^2\mathcal{E}^2\tau}{m_0}\frac{m_0c^2}{E}\nu J_\nu'^2(\nu\beta)\left(\frac{1-\beta^2}{\beta}\frac{J_\nu(\nu\beta)}{J_\nu'(\nu\beta)} - \frac{3}{2\nu}\beta^2\right), \quad (1)$$

where e is the electron charge, τ the average lifetime, E the energy, m_0c^2 the rest energy, $c\beta$ the velocity of motion, and J_ν a Bessel function of order ν.

Formula (1) has been derived for resonance, when

$$\omega = \nu\Omega = \nu\frac{eHc}{E}. \quad (2)$$

In the case when

$$\frac{1-\beta^2}{\beta}\frac{J_\nu(\nu\beta)}{J_\nu'(\nu\beta)} > \frac{3}{2\nu}\beta^2 \quad (3)$$

we obtain ultimately stimulated absorption ($W_\nu < 0$), which takes place, for example, in the nonrelativistic case ($\beta \ll 1$). As shown by Schneider [1], emission in the nonrelativistic approximation is possible for motion in a magnetic field only if resonance is violated.

In the case when