

assumed model already at small degrees of superheat.

4. By the usual procedure [5] we find that the system becomes unstable in the temperature region

$$\theta_1 \leq \theta \leq \theta_2, \text{ where } \theta_1 = T[1 + 2(T/\Delta E)]; \quad \theta_2 = \Delta E/7.$$

The value of the critical intensity of the light (corresponding to  $\theta = \theta_1$ ) is

$$I_1 = \frac{\alpha}{k} \left( \frac{2}{\alpha r_{ph}} \right)^{3/2} \left( \frac{T}{\Delta E} \right)^{15/4} \exp(\Delta E/4T).$$

An estimate for silicon at  $T \approx 300^\circ\text{K}$ , illuminated with a neodymium laser, gives a value  $I_1 \approx 4 \times 10^{24}$  quanta/cm<sup>2</sup>sec. It is interesting to note that the obtained value is close to the threshold intensity at which destruction of the sample takes place [3].

An analysis of the system of equations (1) and (2) gives grounds for assuming that under certain conditions there is realized a regime in which the density  $n$  and the temperature  $\theta$  experience periodic oscillations.

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- [1] L.M. Blinov, E.A. Bobrova, V.S. Vavilov, and G.N. Galkin, Fiz. Tverd. Tela 9, 3221 (1967) [Sov. Phys.-Solid State 9, 2537 (1968)].
- [2] E.A. Bobrova, V.S. Vavilov, and G.N. Galkin, Fiz. Tekh. Poluprov. 3, 1232 (1967) [Sov. Phys.-Semicond. 3, 1032 (1967)].
- [3] A.A. Grinberg, R.F. Mekhtiev, S.M. Ryvkin, V.M. Salmanov, and N.D. Yaroshetskii, Fiz. Tverd. Tela 9, 1390 (1967) [Sov. Phys.-Solid State 9, 1085 (1967)].
- [4] J.S. Blakemore, Semiconductor Statistics, Pergamon, 1961.
- [5] A.A. Andronov, A.A. Vitt, and S.E. Khaikin, Teoriya kolebanii (Theory of Oscillations), Moscow, 1959.

#### DAMPING OF HIGH FREQUENCY SOUND IN A FERMI-BOSE LIQUID

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With decreasing temperature, solutions of He<sup>3</sup> in He<sup>4</sup> become laminated at an He<sup>3</sup> concentration in the solution  $x > 6 \times 10^{-2}$  ( $x = n_3/(n_3 + n_4)$ , where  $n_3$  and  $n_4$  are the numbers of He<sup>3</sup> and He<sup>4</sup> atoms per unit volume). At  $x < 6 \times 10^{-2}$  and  $T \ll T_F$ , where  $T_F$  is the degeneracy temperature of the Fermi component ( $T_F = 0.3^\circ\text{K}$  at  $x = 6 \times 10^{-2}$ ), the solution forms a mixture of Fermi and Bose liquids. We shall henceforth make use of the phenomenological theory of a Fermi-Bose liquid [1, 2]. Calculation shows that the contribution of the phonons to the damping of the sound at  $T \ll T_F$  can be neglected in a wide range of concentrations, and then the dissipation is determined by the collisions of the Fermi excitations with one another. In the collisionless region there can propagate in a Fermi-Bose liquid, in general, both first sound (density oscillations) and zero sound, but the available data on the Landau  $f$ -function [3] indicates that zero sound experiences a strong Landau damping in the solution.

We consider below the damping of first sound in a solution in the collisionless regime. The first to calculate the damping coefficient of first sound in a Fermi-Bose liquid in the limit  $x \ll 1$  was Baym [4]. Baym used the  $\tau$ -approximation, where the relaxation time of the Fermi excitations  $\tau$  plays the role of an arbitrary parameter. In essence, the  $\tau$ -approximation means that only the zeroth and the first harmonics of the expansion of the distribution function with respect to the angles is retained in the collision integral of the kinetic equation (see, e.g., [5]). With Eq. (3) as an example, this means that it would be necessary to retain in its right-hand side the first two terms of the sum. However, as will be shown below, the coupling of the kinetic equation for the Fermi excitations with the equations for continuity and superfluid motion makes the contribution of the second harmonic to the damping of the same order as of the first two even in the limit  $x \ll 1$ .

The problem of the damping of high-frequency sound in a solution, as will be seen from what follows, is analogous to the problem of damping of zero sound in a Fermi liquid [6]. Just as in [6], we find the damping coefficient as the first term of the series of perturbation theory in the small parameter  $1/\omega\tau$ , defined by the relation (the notation is the same as in [6])

$$\frac{1}{\omega\tau} = \frac{m^*3T^2}{12\omega\pi^4\hbar^6} \int \frac{W(\theta, \phi)}{\cos \theta/2} \frac{d\Omega}{2\pi}, \quad (1)$$

and we neglect corrections of relative order  $T^2/T_F^2$ .

The kinetic equation of interest to us for the Fermi excitation in a liquid differs from the kinetic equation for a pure Fermi liquid [6] in the presence of two terms proportional to  $\delta n_4$  and to the superfluid velocity  $v_s$ , and their form is determined by the form of the spectrum of the Fermi excitations in the solution [2].

$$\epsilon(\mathbf{p}) = \epsilon_0(n_3, n_4) + \frac{p^2}{2m^*} + \frac{\Delta m}{m^*} \left(1 + \frac{F_1}{3}\right) p v_s + \int f(\mathbf{p}, \mathbf{p}') \delta n(\mathbf{p}') d\mathbf{r}'. \quad (2)$$

Here

$$\Delta m = m^* \left(1 + \frac{F_1}{3}\right)^{-1} - m_3, \quad \frac{F_1}{3} = \frac{\rho_F m^*}{\pi^2 \hbar^3} \overline{f \cos \chi}.$$

After making exactly the same transformations as in [6] and eliminating the variables  $n_4$  and  $v_s$  with the aid of the continuity and superfluid-motion equations, identical to those written out in [2], the kinetic equation takes the form

$$\begin{aligned} & (\xi - \cos \chi) \sum_{n=0}^{\infty} \nu_n(\tau) P_n(\cos \chi) - \langle \nu_0 \rangle \cos \chi \left\{ F_0 - 3\lambda \frac{m_4}{m^*} \left[ a v^2 \left( a - \frac{s^2 m_3}{v^2 m_4} \right) \right. \right. \\ & \left. \left. + \frac{\Delta m}{m_4} \left( 1 + \frac{F_1}{3} \right) \left( a - \frac{m_3}{m_4} (1 + \beta x) s \cos \chi \right) \right] \right\} - \langle \nu_1 \rangle \cos \chi \left\{ \frac{F_1}{3} \cos \chi - \lambda \left[ s a \right. \right. \\ & \left. \left. + \frac{\Delta m}{m_4} \left( 1 + \frac{F_1}{3} \right) (1 + \beta x) \cos \chi \right] \right\} = \frac{3is}{4\omega\tau} \left( 1 + \frac{\tau^2}{\pi^2} \right) \sum_{n=0}^{\infty} \frac{\beta_n}{\beta_0} \nu_n P_n(\cos \chi), \quad (3) \end{aligned}$$

where  $v_n^0$  denotes the value of  $v_n$  without allowance for the damping,

$$\lambda = \frac{x}{1 + \beta x - s^2/u^2}, \quad u^2 = \frac{n_4}{m_4 v_F^2} \frac{\partial \mu_{04}}{\partial n_4} = \frac{c^2}{v_F^2},$$

$c$  is the velocity of first sound in  $\text{He}^4$ ,  $\mu_{04}$  is the chemical potential of pure  $\text{He}^4$ ,

$$\alpha = \frac{n_4}{m_4 c^2} \frac{\partial \epsilon_0}{\partial n_4}, \quad \beta = \frac{n_4^2}{m_4 c^2} \frac{\partial^2 \epsilon_0}{\partial n_4^2}, \quad \langle v_n \rangle = - \int_{-\infty}^{\infty} \frac{\partial r_0}{\partial t} v_n(t) dt, \quad s = \frac{\omega}{k v_F},$$

$$\xi = s \left( 1 + \frac{i}{\omega \tau} \right), \quad \beta_n = \int \frac{d\Omega}{2\pi} \frac{W(\theta, \phi)}{\cos \theta/2} \left[ P_n(\cos \theta) - P_n \left( \cos^2 \frac{\theta}{2} \right. \right. \\ \left. \left. + \sin^2 \frac{\theta}{2} \cos \phi \right) - P_n \left( \cos^2 \frac{\theta}{2} - \sin^2 \frac{\theta}{2} \cos \phi \right) \right], \quad t = \frac{\epsilon - \mu_3}{T},$$

$$k v = k v \cos \chi.$$

Equation (3) coincides in form with the corresponding equation in [6]. Therefore, proceeding exactly as in [6], we obtain the damping coefficient of first sound in the solution. We note that in the derivation of (3) we never used the smallness of  $x$ . An essential limitation in the derivation of (3) is the fact that in writing out the integral terms of the spectrum (2) we can take into account an arbitrary but finite number of Landau parameters. We retain in (3) only  $F_0$  and  $F_1$ . We shall not write out for the damping coefficient of the first sound the formula that is correct for arbitrary  $x$ , since it is too cumbersome, and proceed immediately to the case  $x \ll 1$ .

We note that when  $x \ll 1$  we have  $s \approx u$  (see [2]) and consequently  $\lambda \sim 1$ , and in addition  $u^2 \sim x^{-2/3}$ . We recall that  $F_n \sim p_F \sim x^{1/3}$ . We see thus that the first terms in the curly brackets following  $\langle v_0 \rangle$  and  $\langle v_1 \rangle$ , resulting from the integral term in (2), contain an extra small factor  $\sim x$  compared with the terms resulting from the coupling. Therefore terms of Fermi-liquid origin can be neglected in (3). Simple calculations show that  $\langle v_2 \rangle / \langle v_0 \rangle \sim 1$ . As a result,  $v_2$  makes a contribution to the damping of the same order as  $v_0$  and  $v_1$ . Retaining in the expression for the damping coefficient the term of lowest degree in  $x$ , we have

$$\gamma = \text{Im } k = \frac{x m_4 p_F^2 T^2}{90 \pi^2 \hbar^6 c^3} \left( \alpha + \frac{\Delta m}{m_4} \right)^2 \int \frac{d\Omega}{2\pi} \frac{W(\theta, \phi)}{\cos \theta/2} \left[ 1 + \frac{1}{R} \left( 1 - 3 \sin^4 \frac{\theta}{2} \sin^2 \phi \right) \right], \quad (4)$$

where

$$R = \frac{m_4}{m^*} \left( \alpha + \frac{\Delta m}{m_4} \right)^2 - \frac{\Delta m}{m_4}, \quad m^* \approx m^*(x=0), \quad \Delta m = m^* - m_3.$$

Recognizing that  $W(\theta, \phi) \sim x^{1/3}$ , we obtain  $\gamma \sim x^2$ . If we take (1) to be the definition of the relaxation time  $\tau$ , then (4) differs from Baym's result in the presence of the second term in the square brackets, which, as already mentioned, is the result of taking  $v_2$  into account.

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- [1] I. Bardeen, G. Baym, and D. Pines, Phys. Rev. 156, 207 (1967).
- [2] I.M. Khalatnikov, Zh. Eksp. Teor. Fiz. 55, 1919 (1968) [Sov. Phys.-JETP 28 1014 (1969)].
- [3] I.C. Wheatley, Quantum Fluids, Amsterdam, 1966.
- [4] G. Baym, Phys. Rev. Lett. 17, 952 (1966).
- [5] A.A. Abrikosov and I.M. Khalatnikov, Usp. Fiz. Nauk 66, 177 (1958) [Sov. Phys.-Usp. 1, 68 (1958)].
- [6] B.S. Luk'yanchuk and D.M. Semiz, Zh. Eksp. Teor. Fiz. 60, 1067 (1971) [Sov. Phys.-JETP 33, No. 3 (1971)].

NONEQUILIBRIUM DISSOCIATION OF A MOLECULAR GAS UNDER THE INFLUENCE OF RESONANT LASER RADIATION, WITH ALLOWANCE FOR VIBRATION-VIBRATION COLLISIONS

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1. With the appearance of powerful lasers, new possibilities were uncovered for acting on intermolecular processes, namely excitation and dissociation of molecules, initiation and control of chemical reactions, etc. Research on the excitation and dissociation of isolated molecules in a laser field is the subject of a number of papers [1 - 5]. In [1] is considered the interaction between a molecule and a field in the case when the acting force either coincides with the field frequency and resonant transitions are then possible through several levels, or in the case when the acting force is non-linear in the field, and different field harmonics take part in the excitation. The probabilities of two- and three-quantum dissociation of a molecule were calculated in [2, 5]. Reference [4] is devoted to excitation of a Morse oscillator. In all the cited papers, the threshold density of the laser radiation was  $10^{14} - 10^{16}$  W/cm<sup>2</sup>.

Dissociation and chemical reactions in molecular gases, induced by resonantly absorbed infrared laser radiation, were experimentally investigated in [6]. The threshold fluxes were of the order of  $10^2$  W/cm<sup>2</sup>. Physically this circumstance, unlike in [1 - 5], is due to selective interaction of the radiation with one of the vibrational degrees of freedom of the molecule and with vibration-vibration relaxation of the excitations. In this case the vibrational temperature can noticeably exceed the translational temperature of the medium, as was pointed out, for example, in [7].

In the present paper we solve the nonstationary self-consistent problem of dissociation of a molecular gas under the influence of resonant laser radiation under conditions when the duration of the pulse is limited only by the time  $t_m$  of the vibration-translation relaxation. It is assumed that the considered vibrational branch of the molecule is almost harmonic for vibrational exchange up to levels from which dissociation proceeds with a probability close to unity. On the other hand, the anharmonicity is always important for the absorbed laser radiation, so that absorption is possible only between the first two levels of the molecule.

2. We assume that the vibrational energy exceeds the quantum energy  $h\omega$  of the harmonic oscillator. In this case the time dependence of the molecule energy distribution function  $F(t, \epsilon)$  ( $\epsilon$  is the energy and  $t$  is the time) can be described by a kinetic equation in the diffusion approximation:

$$\frac{\partial \Phi}{\partial t} = - \frac{\partial J}{\partial \epsilon},$$

$$J = -\epsilon \left[ \Phi + (1 + \nu(\epsilon - \epsilon_*)) \frac{\partial \Phi}{\partial \epsilon} \right]. \quad (1)$$