The experimental setup is shown in Fig. 1. Here 1 - mirror with reflection coefficient 99% at λ = 6943 Å, 2 - ruby crystal 120 mm long, 3 - lens with f = 15 cm, focusing the laser radiation in the center of a cell 4 (40 cm long) with the investigated liquid, 5 - slit spectrograph.

We used in the investigation benzene, cyclohexane, and carbon tetrachloride, which differ greatly from each other both with respect to excitation of SMBS and with respect to excitation of SRS.

Initially the laser resonator was made up of mirror 1 and the end of the ruby rod. The emission oscillograms (Fig. 2) show that the giant pulse was formed by one of the first free-generation "spikes." The energy of the single pulse incident on the medium was ~ 1.5 and ~ 1.2 J for C_6H_6 and C_6H_{12} respectively. The pulse duration at half-height was ~ 25 nsec for C_6H_6 and C_6H_{12} and ~ 20 nsec for CCl_4. The high pulse energy is due to the relatively high reflection coefficient of backward SMBS, which greatly decreases the threshold level of the inverted population in the ruby crystal and increases the radiation energy density in the resonator. Spectral investigations performed with the aid of a Fabry-Perot interferometer have demonstrated the presence of several (4 - 5) SMBS components in the radiated pulse. The width of the individual components was ~ 10^{-2} cm^{-1}, which is much less than the radiation line width in the free-generation mode (0.35 cm^{-1}).

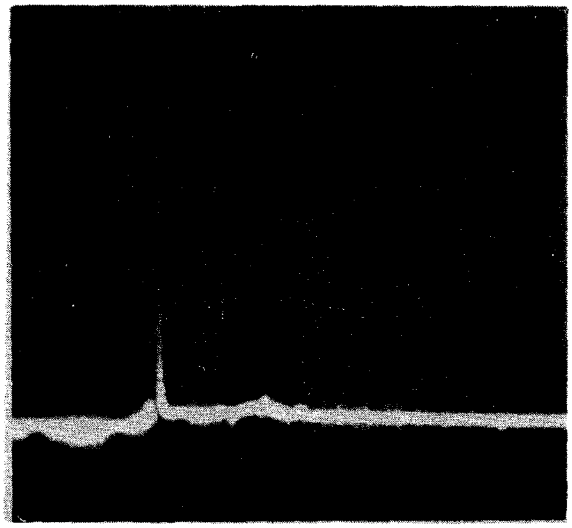

Fig. 2. Oscillogram of pulse incident on medium. Sweep duration 500 500 μsec.

In benzene, four Stokes components were excited with a shift of 992 cm^{-1}. The power of the strongest, second Stokes component exceeded 15 MW. In C_6H_2 there were excited two Stokes components pertaining to the 2852 cm^{-1} molecular vibration, two Stokes components of 801 cm^{-1} vibration, and also two "combination" Stokes frequencies with shifts (2852 + 801) and (2 × 2852 + 801) cm^{-1}. The power of the most intense component with the 2852 cm^{-1} shift reached 20 MW. In CCl_4, four Stokes components with a shift of 459 cm^{-1} were excited, with the second Stokes component having the largest power (~ 10 MW).

It was established that the simple method used in the present investigation to excite SRS in liquids ensures a much higher power of the excited pulse and a more effective excitation of the SRS than a setup with a saturable absorber for laser Q-switching, operating at higher ruby-rod pumping levels.

[1] D. Pohl, Phys. Lett. 24A, 239 (1967).
[2] A.Z. Grasyuk, V.V. Ragul'skii, and F.S. Faizulov, ZhETF Pis. Red 9, 10 (1969) [JETP Lett. 9, 6 (1969)].

ANOMALOUS RESISTANCE OF A PLASMA IN THE CASE OF ION-ACOUSTIC TURBULENCE

G.E. Bekshtein and R.Z. Sagdeev
Submitted 2 February 1970
ZhETF Pis. Red. 11, No. 6, 297 - 300 (20 March 1970)

The so-called anomalous resistance of a rarefied plasma, which has been observed in a number of experiments [1, 2], is attributed to the scattering of the electrons by the random fields of oscillations that become unstable when the current flows. It is assumed that the principal role is played here by the ion-acoustic instability [3, 4]. In the general case, however, the picture is greatly complicated by the "runaway" of some of the electrons. In this sense, the simplest case is that of a current perpendicular to the magnetic field. This is precisely the situation realized in experiments with collisionless shock waves. The anomalous resistance of the plasma, corresponding to the measured values of the thickness of the shock wave front, is lower by approximately one order of magnitude than that obtained in the theory of weak turbulence for the ion-acoustic instability [3, 4]. This means that the nonlinear Landau damping on ions (the main stabilizing effect in [3, 4]) is small and a more effective damping mechanism is needed.

A different model was considered in [5], but there the current velocity u turns out to be too large and does not increase with time, thus contradicting the experimental data.

We propose below a different mechanism of development of ion-acoustic turbulence for the case of a current flowing across the magnetic field[1]. If the electrons are magnetized ($\nu_e/\omega_{He} \ll 1$, where $\nu_e \sim \omega_{pe} W/mT_e$ is the electron scattering frequency and W the oscillation energy density), their velocity distribution is almost isotropic in the reference frame connected to the current. Then the electronic increment is determined by two quantities (u and the effective temperature T_e of the electrons)

$$\gamma_e \sim \frac{M}{m} \frac{\omega_k^3}{k^2} \frac{u}{(T_e/m)^{3/2}} \cos\theta' \tag{1}$$

(θ' - angle between the wave vector and the direction of the current). The current excites a broad spectrum of ion-acoustic oscillations. The contribution to the ion damping is made by a small fraction of the ions located in the "tail" of the ion distribution function $f_i(v, \theta)$ obeying the quasilinear equation. This equation and the condition $\gamma = \gamma_e + \gamma_i = 0$ (closeness to the instability threshold for all waves with $\cos\theta' < 0$ and $\gamma < 0$ at $\cos\theta' < 0$) admit of a self-similar formulation, such that the ratios of the electron current velocity to the ion thermal energy and to the electron energy, and the fraction of the resonant ions ($\Delta n/n$), remain constant as the plasma becomes heated. The ion velocity distributions and the oscillation spectrum then also remain constant. The latest measurements of the current an- energy of ions agree precisely with such a steady state [6].

Let us express the spectral energy density of the oscillations in the form $w_k(t) = A(t)w(z, \theta')$ ($z = kr_0$ and r_0 - Debye radius), and let us introduce in place of t a new variable

$$\tau = \frac{2\pi e^2}{M^2 \omega_{pi}} \int_0^t W(t)dt.$$

Then the self-similar velocity variable has the form $\xi = v/\tau^{1/2}$. The quasilinear equation and the condition $\gamma = 0$ have in terms of the self-similar variables rather complicated final forms and are not presented here. So far, the parameters $u_0 = u/(T_e/m)^{1/2}$ remains free. The obtained equations are so complicated, that it is impossible to obtain their exact solution, i.e., the functions $f(\xi, \theta)$ and $w(z, \theta')$. However, the values of u of interest to us and the

[1]Usually $\omega_{He} \ll \omega_{pe}$ and the magnetic field does not influence the oscillations.

ion and electron energies can be estimated without knowing the exact solution. To this end, we find the first and second moments of the equation for $f_i(v, \theta)$ - the momentum and energy conservation laws.

We introduce the effective temperature T_i of the resonant ions[2)]. The energy conservation law yields the well-known expression [4]

$$\frac{T_e}{T_i} \sim \frac{ux}{c_s}; \quad \left(c_s^2 = \frac{T_e}{M}; \quad x = \frac{\Delta n}{n}\right). \tag{2}$$

estimating the ion decrement to be

$$\gamma_i \sim \frac{\omega_k^3}{k^2} \frac{xc_s}{(T_i/M)^{3/2}}$$

and comparing it with the electron decrement (1), we obtain

$$x \sim \left(\frac{m}{M}\right)^{1/4}\left(\frac{T_i}{T_e}\right)^{1/4}; \quad u \sim c_s\left(\frac{M}{m}\right)^{1/4}\left(\frac{T_e}{T_i}\right)^{5/4}. \tag{3}$$

we consider further the momentum of the resonant ions. The momentum lost by the electrons in scattering goes over to the ions, $P_i = nmu\nu_e$. Since $T_i \gtrsim T_e = Mc_s^2$, the ion distribution function can be represented in the form $f_i(v, \theta) = f_0(v) + f_1(v, \theta)$, where the anisotropy part is

$$f_1 \sim \frac{c_s}{(T_i/M)^{1/2}} f_0 \lesssim f_0 .$$

Thus,

$$P_i \sim \int M\, v f_1 d^3v \sim nxMc_s .$$

Finally we obtain

$$u \sim c_s\left(\frac{M}{m}\right)^{1/4}; \quad T_i \sim T_e; \quad x \sim \left(\frac{m}{M}\right)^{1/4}. \tag{4}$$

Relations (4) contain, of course, factors on the order of unity, which can be determined without knowledge of the exact solution only by comparison with the experimental data. Let us calculate the width of the shock wave δ, by determining ν_e with the aid of the formulas given above

$$\delta \sim \frac{c}{\omega_{pe}}\left(\frac{M}{m}\right)^{1/4}\frac{v_A}{c_s} \sim \frac{c}{\omega_{pe}}\left(\frac{M}{m}\right)^{1/4} \tag{5}$$

($v_A/c_s \sim 1$ for not very weak shock waves).

In a hydrogen plasma, experiment yields $\delta \sim 10\ c/\omega_{pe}$, meaning that the numerical factor in (5) is close to unity. The "isotropic effect" measured for the cases He, Ar, and Xe agrees with (8) within the limits of experimental errors [2].

[1] Plasma Physics and Controlled Nuclear Fusion Research, II, IAEA, Vienna, 1966.

[2)] $T_i \gtrsim T_e$ by definition, in order to make the ions resonant.

[2] Plasma Physics and Controlled Nuclear Fusion Research, I, IAEA, Vienna, 1969.
[3] B.B. Kadomtsev, in: Voprosy teorii plazmy (Problems of Plasma Theory), Vol. 4, 1964.
[4] R.Z. Sagdeev, Proc. Symposia in Appl. Mathematics 18, 281 (1965).
[5] L.P. Rudakov and L.V. Korablev, Zh. Eksp. Teor. Fiz. 50, 220 (1966) [Sov. Phys.-JETP 23, 145 (1966)]; L.M. Kovrizhnykh, ibid. 51, 1795 (1966) [24, 1210 (1967)].
[6] R.Kh. Kurtmullaev, Proc. Conf. on Collisionless Shock Waves, Frascati, 11 - 18 June 1969 (in press).

"MIXMASTER UNIVERSE" AND THE COLD VARIANT

Ya.B. Zel'dovich
Institute of Applied Mathematics, USSR Academy of Sciences
Submitted 5 February 1970
ZhETP Pis. Red. 11, No. 6, 300 - 303 (20 March 1970)

Recently Misner [1, 2] raised the question of a cosmological model in which the observed isotropy of the universe (in particular, the isotropy of the relict radiation [3] might have a natural explanation. In other words, one seeks in the initial stage at high density a more general anisotropic and inhomogeneous solution that approaches asymptotically in the course of time Friedmann's homogeneous and isotropic solution observed at the present time.

Misner proposed that such properties is possessed by the solution of Belinskii, Lifshitz, and Khalatnikov [4, 5], which they obtained during the course of an investigation of the singularity of the equations of general relativity theory.

The initial solution is homogeneous but anisotropic, and describes a closed world. The solution differs in that it has "no horizon," i.e., the light signal has time to run around the entire world many times during the anisotropic stage.

During the early stage, matter (quanta, particles, antiparticles) does not influence the evolution of the solution. The gravitation of matter can be neglected.

If the matter has a homogeneous distribution, a transition to Friedmann's isotropic closed solution takes place in the course of time when the gravitation of the matter becomes appreciable.

The inhomogeneous distribution of matter can be regarded as a density perturbation superimposed on the homogeneous distribution. In the stage during which the gravitation of matter does not influence the general expansion, it should likewise not influence the perturbations. This means that there is no gravitational instability, and the perturbations attenuate, during the course of the adiabatic expansion, under the influence of the viscosity (particularly neutron viscosity [6]) and as the result of the formation of shock waves [7].

The initial inhomogeneity of the metric, which does not depend on the presence of matter ($R_{ik} = 0$), can be regarded as gravitational waves superimposed on the homogeneous solution. The gravitational waves also attenuate during the course of the expansion.

The "no horizon" condition is obviously necessary to permit equalization of the inhomogeneity and damping out of the perturbations. The closedness of the world limits the maximum length of the perturbation wave; this limitation is necessary for adiabatic damping to be applicable. It is the very possibility of equalization of the homogeneities which is a new property of the mixmaster