Experimental observation of the drift-dissipative instability in afterglow plasmas

1970 J. Phys. A: Gen. Phys. 3 L21

(http://iopscience.iop.org/0022-3689/3/3/005)

which is precisely the result obtained by Zeldovich. Note that even if we had used the correct relativistic expression for the kinetic energy and included the rest-mass energy, the potential energy would have been dominant in the high-density limit.

If the many-body potentials are included in equation (4) a_j is positive if V_j is repulsive, and negative if V_j is attractive. Many-body potentials are extremely difficult to derive and in general may be repulsive or attractive. In this connection it is interesting to note that the static nuclear many-body potentials derived using old-fashioned perturbation theory (Drell and Huang 1953) were found to be repulsive for j odd and attractive for j even. Hence, as $n \to \infty$, the behaviour of E/Ω is not known, quite contrary to the results of Zeldovich and Harrison.

It should be noted that the true ground-state energy is lower than that given in the Hartree approximation because of the neglect of the Hartree–Fock and correlation energies. These corrections were also neglected by Zeldovich.

In conclusion, we find that the energy per unit volume in the Hartree approximation is given by a power series in the density. The terms in n^3, n^4, ... come from the contributions to E/Ω of the 3, 4, ... body potentials respectively and were not included by Zeldovich and Harrison. As $n \to \infty$ it is not known whether E/Ω tends to a limit or not so that we cannot say anything about the corresponding behaviour of the speed of sound. This is hardly surprising considering the complexity of ultra-dense matter. We suggest that, if all the many-body potentials are properly taken into account, the speed of sound will not exceed that of light.

The author is grateful to Professor S. S. Schweber for suggesting the importance of many-body forces in dense matter and for many discussions. He is also grateful to Professor P. Ledoux for his helpful comments.

Brandeis University,
Waltham,
Massachusetts, U.S.A.

G. Chanmugam†
12th January 1970

Drell, S. D., and Huang, K., 1953, *Phys. Rev.*, **91**, 1527–42.
Harrison, E. R., 1965, *Astrophys. J.*, **142**, 1643–5.
Kadanoff, L. P., and Baym, G., 1962, *Quantum Statistical Mechanics* (New York: Benjamin).
Primakoff, H., and Holstein, T., 1939, *Phys. Rev.*, **55**, 1218–34.
Tsuruta, S., and Cameron, A. G. W., 1966, *Can. J. Phys.*, **44**, 1895.
Zeldovich, Ya. B., 1962, *Sov. Phys.-JETP*, **14**, 1143–7.

Experimental observation of the drift-dissipative instability in afterglow plasmas

Abstract. This letter presents results which support the general shape of the ω against k_z dispersion curve for the drift-dissipative instability, predicted using the 'two-fluid' equations in slab geometry. The wave is observed to be self-excited in the k_z region where a positive growth rate is predicted.

In afterglow plasmas, the usual possible causes of instability such as axial current, non-isotropic velocity distributions, imposed electric fields, etc., are absent. Consequently, most experiments on afterglows have shown them to be stable; however,

† Now at Institut d'Astrophysique, Université de Liège, Cointe-Sclessin, Belgium.

recent papers (Pigache and Harding 1969, Dodo 1969) have reported instabilities in this type of plasma. This paper reports further measurements on a low-frequency self-excited oscillation present in an inhomogeneous afterglow, in which there was a radial density gradient perpendicular to an axial homogeneous magnetic field B_0 variable between 100 G and 3 kG.

The plasma was produced either by using a short pulse of electrons emitted from a hot cathode to ionize a neutral gas (helium or hydrogen), or by 'pulsing-off' an rf discharge. Typical initial densities $n_0 \sim 10^{10}\ \mathrm{cm}^{-3}$ were obtained, and an average electron temperature $T_e \sim 1000$ K was measured with a single Langmuir probe. The experiments were performed in both helium and hydrogen plasmas, in two different tubes of 2·5 and 5·0 cm diameter each 200 cm long. This varied the inverse density scale length

$$\kappa = \frac{1}{n_0}\frac{\mathrm{d}n_0}{\mathrm{d}r}$$

and resulted in values 2·3 cm^{-1} and 1·3 cm^{-1} respectively, for the two tubes. The mode number of the instability was determined by amplitude and phase measurements from four single probes spaced around the azimuth. It was found to be

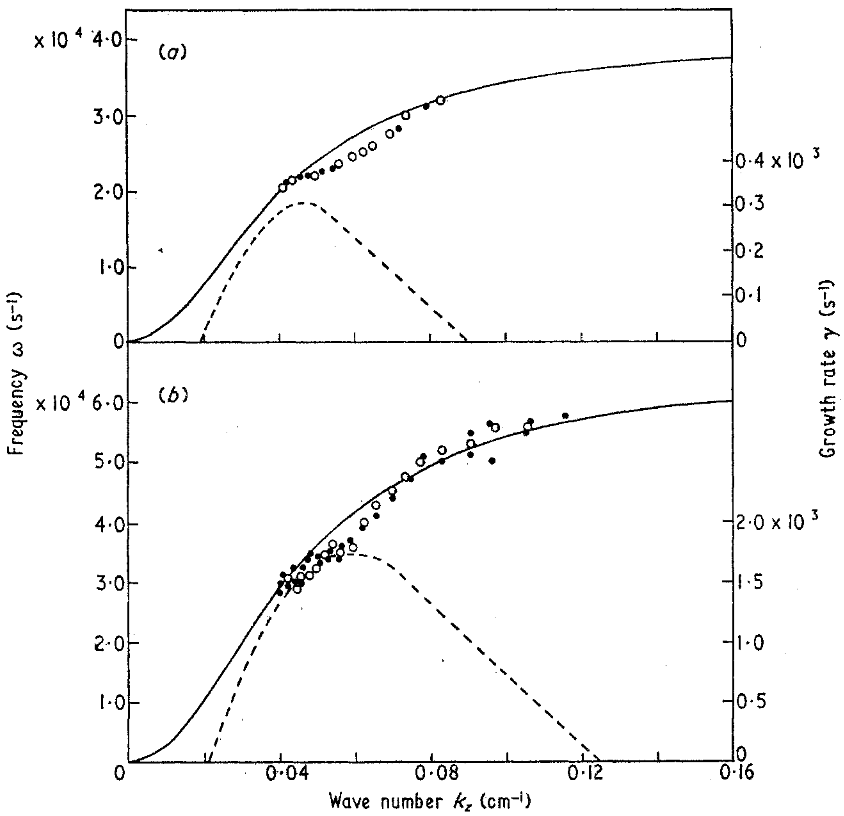

Figure 1. Dispersion diagram for a hydrogen plasma; $m = +1$. (*a*) $H = 500$ G, small tube, and (*b*) $H = 350$ G. Points marked are experimental points for the Re(ω). The full curve shows the theoretically computed curve for Re(ω) and refers to the left-hand scale. The broken curve shows the theoretically computed growth rate γ and refers to the right-hand scale. ● rf plasma. ○ Hot cathode plasma.

mainly an $m = +1$ propagating mode, although in the larger tube some results were obtained for an $m = +2$ mode. Longitudinal amplitude and phase measurements were made using an axially moving single probe, and these showed that the oscillation was a standing wave in the axial direction with a wavelength λ_z approximately equal to the plasma column length. By employing a movable anode the plasma column length was varied, thus changing the axial wavelength. In this way, the oscillation frequency ω was measured as a function of axial wavenumber $k_z(= 2\pi/\lambda_z)$, and an ω–k_z dispersion relationship was obtained. This experiment was repeated under differing conditions by varying the axial magnetic field and the neutral pressure for both the helium and hydrogen plasmas contained in the two different tubes. It was found that the oscillation frequency ω tended to be independent of time (or density) in the afterglow decay, and that ω reduced in value as $k_z \to 0$, and tended towards the drift frequency ω^* as k_z assumed larger values. Typical results are shown in figures 1(*a*) and 1(*b*), for a hydrogen plasma, under the conditions shown marked on the diagram.

A theory of the drift-dissipative instability of a plasma in a magnetic field, with no longitudinal current, has been developed by Timofeev (1964 *a,b*). The 'two fluid' equations of motion were considered in 'slab' geometry, in which both an electron–neutral collision frequency ν_e and an ion–neutral frequency ν_i were taken into account. These equations, in the electrostatic approximation, were linearized for the ions and electrons separately, and density n' and potential perturbations ϕ' of the form $\exp\{-i(\omega t - \bar{k} . \bar{r})\}$ were assumed. By eliminating these quantities n' and ϕ', using the equations of continuity, the following complex quadratic dispersion equation was obtained:

$$\omega^2 + i\omega(\omega_s + D_e k_z{}^2 + \nu_i) - i\omega_s\omega^* - D_e k_z{}^2 \nu_i = 0 \qquad (1)$$

where

$$\omega^* = -\frac{k_y c T_e \kappa}{eB_0}$$

$$\omega_s = \frac{k_z{}^2 \Omega_e \Omega_i}{k_\perp{}^2 \nu_e}$$

and $D_e = T_e/m\nu_e$. Here $k_\perp{}^2 = k_x{}^2 + k_y{}^2$, and $\Omega_e\,(= eB/mc)$ and $\Omega_i\,(= eB/M_i c)$ are the electron and ion cyclotron frequencies respectively. This dispersion equation has been obtained in the approximations: (i) $\omega \ll \Omega_i$, (ii) $\nu_e \ll \Omega_e$, and (iii) $\nu_i \ll \Omega_i$, and these assumptions are reasonably well satisfied experimentally. Equation (1) can be solved numerically, under the appropriate conditions, for $\mathrm{Re}(\omega)$ and the growth rate $\gamma = \mathrm{Im}(\omega)$, as a function of k_z.

In order to compare theory with experiment, it is necessary to know the values of the collision frequencies ν_e and ν_i. The electron–neutral collision frequency ν_e was obtained by using a microwave cavity technique (Buchsbaum and Brown 1957). They showed that

$$\nu_e = \frac{\pi f_0{}^2}{\Delta f}\left(\frac{1}{Q} - \frac{1}{Q_0}\right)$$

where Q_0 is the 'Q' value of the empty cavity and Q its value loaded with the plasma, f_0 is the resonant frequency and Δf the change in resonant frequency when plasma is present. A cylindrical cavity ($f_0 \sim 1350$ MHz, $Q_0 = 5000$) oscillating in the TM_{010}

mode was used on the outside of the plasma tube, and with a constant Δf value chosen (corresponding to a constant density value in the afterglow) the change in Q value was measured as the neutral pressure p in the tube was changed. A linear relationship between p and $1/Q$ resulted, thus indicating that the electron–neutral collision frequency was the measured quantity. Values of

$$\nu_e/p = 0{\cdot}70 \pm 0{\cdot}15 \times 10^9 \, \mathrm{s}^{-1} \, \mathrm{torr}^{-1}$$

for the hydrogen plasma and

$$\nu_e/p = 0{\cdot}80 \pm 0{\cdot}15 \times 10^9 \, \mathrm{s}^{-1} \, \mathrm{torr}^{-1}$$

for the helium plasma, were computed. The ion–neutral collision frequency ν_i proved to be a more difficult parameter to measure, and simple kinetic theory considerations on the hard-sphere model predict $\nu_i/\nu_e = v_i/4v_e$ where v_i and v_e are the ion and electron thermal velocities respectively. Thus, values of

$$\nu_i/p = 1{\cdot}0 \pm 0{\cdot}15 \times 10^6 \, \mathrm{s}^{-1} \, \mathrm{torr}^{-1}$$

for the helium plasma and

$$\nu_i/p = 1{\cdot}50 \pm 0{\cdot}15 \times 10^6 \, \mathrm{s}^{-1} \, \mathrm{torr}^{-1}$$

for the hydrogen plasma, were adopted.

Using these values, $\mathrm{Re}(\omega)$ and the growth rate $\gamma = \mathrm{Im}(\omega)$ were computed as a function of k_z from equation (1) for the particular conditions prevailing in each experiment. Typical theoretical curves are shown in figures 1(*a*) and 1(*b*). The full curves indicate the $\mathrm{Re}(\omega)$ and refer to the left-hand scale, and the broken curves show the predicted growth rate γ and refer to the right-hand scale. It is seen that in this case good agreement is obtained, and similarly good agreement is obtained for other $m = +1$ and $m = +2$ results, in spite of the approximations in the theory and the possible errors in the experiment. Therefore, it is concluded that it is the drift-dissipative instability which is observed in these plasmas.

UKAEA Research Group,
Culham Laboratory,
Abingdon, Berks.

M. W. ALCOCK
B. E. KEEN
12th February 1970

BUCHSBAUM, S. J., and BROWN, S. C., 1957, *Phys. Rev.*, **106**, 196.
DODO, T., 1969, *Plasma Phys.*, **11**, 881.
PIGACHE, D. R., and HARDING, G. N., 1969, *Plasma Phys.*, **11**, 805.
TIMOFEEV, 1964 a, *Sov. Phys.–Tech. Phys.*, **8**, 682.
—— 1964 b, *Sov. Phys., Dokl.*, **8**, 890.

Screening of the photon propagator in many-body optics

Abstract. An internally consistent treatment of the interaction of light with a molecular fluid in terms of a screened photon propagator is reported. Screening simplifies the description of multiple scattering in terms of Ursell functions and the treatment of surface effects. In a translationally invariant approximate theory the screened photon propagator and the screened radiation reaction are expressed in terms of the refractive index.