Excitation and Control of the Rayleigh-Taylor Instability in a Plasma with a Curved Magnetic Field

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The Rayleigh-Taylor instability is generated in a plasma flowing along a curved magnetic field. The instability, which appears only in the bad-curvature region of the magnetic field, is controlled by changing the flow speed and Larmor radius of the ions. A drastic suppression of the instability is observed in the presence of weak magnetic shear.

It is of current interest to investigate plasma instabilities in a toroidal device.[1] The Rayleigh-Taylor instability is one of the fundamental instabilities in such a device, because an effective gravitation is provided by the curved magnetic field.[2] To our knowledge, however, no clear-cut experiment on this instability has been published, although some related experiments were performed in plasmas of different geometries,[3] i.e., a mirror-confined plasma and a θ-pinch plasma. D'Angelo *et al.* observed some instabilities in a Q-machine plasma with a variable magnetic field curvature.[4] But, it is difficult to relate their result to the Rayleigh-Taylor instability, although the result was discussed theoretically by Chen.[5] Here we report an experiment performed in a plasma with curved magnetic field and helical winding, designed to investigate the Rayleigh-Taylor instability. The steady-state plasma is produced by surface ionization in a single-ended Q machine.[6] The plasma is in equilibrium, i.e., the electric field induced by the curvature drift is canceled out by the short-circuit path through the hot plate.[7] Plasma parameters are easily changed to control the instability.

The experimental apparatus is shown in Fig. 1. The vacuum chamber (diameter $\simeq 4.8$ cm) consists of a straight part (60 cm) and a curved part (major radius $R = 60$ cm, angle of arc $= 140$ deg). The curved magnetic field $\vec{B}$ of up to 2.5 kG is produced along the chamber. The helical conductors with pitch of 31.4 cm and minor radius of 3.5 cm are wound around the chamber to form an $l=2$ helical winding, producing weak magnetic shear. The inverse of the shear length, $L_s^{-1} = (d\iota/dr)/(2\pi R/r)$ (ι is the rotational transform) is variable, for example, from 0 to 2.4×10^{-3} cm^{-1} (at the radial edge of the plasma) for $B = 2.25$ kG. The collisionless potassium plasma (diameter $\simeq 2.0$ cm, density $n_0 = 10^8$–10^9 cm^{-3}), produced under an electron-rich condition, is terminated at a movable 4.5-cm-diam end plate biased negatively. Thus, ions are accelerated by the electron sheath in front of the hot plate[8] and flow along $\vec{B}$ toward the end plate. The flow speed V_0 is controlled by changing the sheath which depends on the potassium feed and the hot-plate temperature $T(=2000$–$2500\,^\circ\mathrm{K})$. The electron temperature T_e is nearly equal to T. The ion temperature perpendicular to $\vec{B}$, $T_{i\perp}$, can also be approximated by T although the ion temperature parallel to $\vec{B}$ is smaller than T because of the sheath acceleration.[6,8] The background gas pressure is kept less than 3×10^{-6} Torr. Three Langmuir probes (1×1 mm^2) P_1, P_2, and P_3 are used to measure the plasma parameters and their fluctuations.

Initially, the end plate is situated at the entrance of the curved part in order to determine the properties of the plasma for the uniform, straight magnetic field. The ion-current fluctuations displayed on a spectrum analyzer show that a drift wave (DW) of 5.9 kHz is excited spontaneously.[9] Then, by moving the end plate along the curved part of the tube, the plasma is guided into the curved magnetic field with no shear. The fluctuations measured by P_1, P_2, and P_3 have

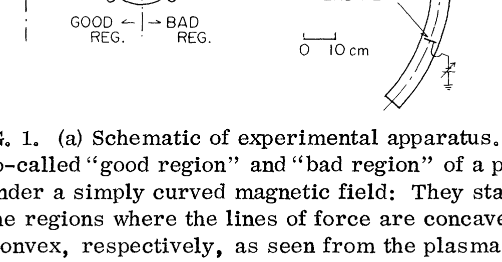

FIG. 1. (a) Schematic of experimental apparatus. (b) So-called "good region" and "bad region" of a plasma under a simply curved magnetic field: They stand for the regions where the lines of force are concave and convex, respectively, as seen from the plasma.

another peak around 7.4 kHz in addition to the DW mentioned above. As shown in Fig. 2, this instability appears only in the bad-curvature retion (see Fig. 1) of the magnetic field. The percentage fluctuation $\tilde{n}/n_0$ is about 10% at the position of the peak amplitude. There is a phase difference of about 180 deg between the fluctuations of density and potential. The instability propagates azimuthally, with the mode $m \simeq 4$, in the direction of electron diamagnetic drift which coincides with that of $\vec{E}_0 \times \vec{B}$ drift ($\vec{E}_0$ is the steady radial electric field) in our experiment. The axial phase change is not observed and thus the parallel wave number $k_z \simeq 0$ ($\ll 1/L$, L being the plasma length). Similar properties were reported previously by us on the instability observed at much higher frequency ($\simeq 50$ kHz).[10] They were well explained by effects of magnetic field curvature on a transverse Kelvin-Helmholtz (KH) instability. The instability described here, however, appears only when the magnetic field is curved, in contrast to the KH instability driven by the $\vec{E}_0 \times \vec{B}$ plasma rotation with velocity shear, which should appear even under a uniform, straight magnetic field. All features of this new instability are quite reasonable for the Rayleigh-Taylor (RT) instability. In this work, the KH instability is not expected to grow, because its instability condition is not satisfied for our values of E_0 and dE_0/dr (see Fig. 2) which are much smaller than those in Ref. 10.

Since the effective gravitation is due to the plasma (ion) flow along the curved magnetic field, it is important to measure the dependence of the instability on V_0. The result is shown in Fig. 3(a), where a charge-exchange method is adopted to measure V_0.[11] With an increase in V_0, the peak amplitude around 7.4 kHz increases, being consistent with the prediction for the RT instability. Figure 3(b) shows the dependence on B. The instability appears for $B \gtrsim 1.75$ kG. The amplitude increases with an increases with an increase in B. This dependence will be explained later by the effect of finite ion Larmor radius on the RT instability. The interesting result is the dependence of the instability on the shear field. As shown in Fig. 3(c), the amplitude of the RT in-

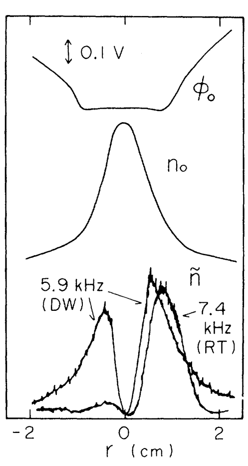

FIG. 2. Radial profiles of perturbed density $\tilde{n}$ of the RT and the DW, measured by the probe P_2, together with the plasma potential φ_0 and density n_0.

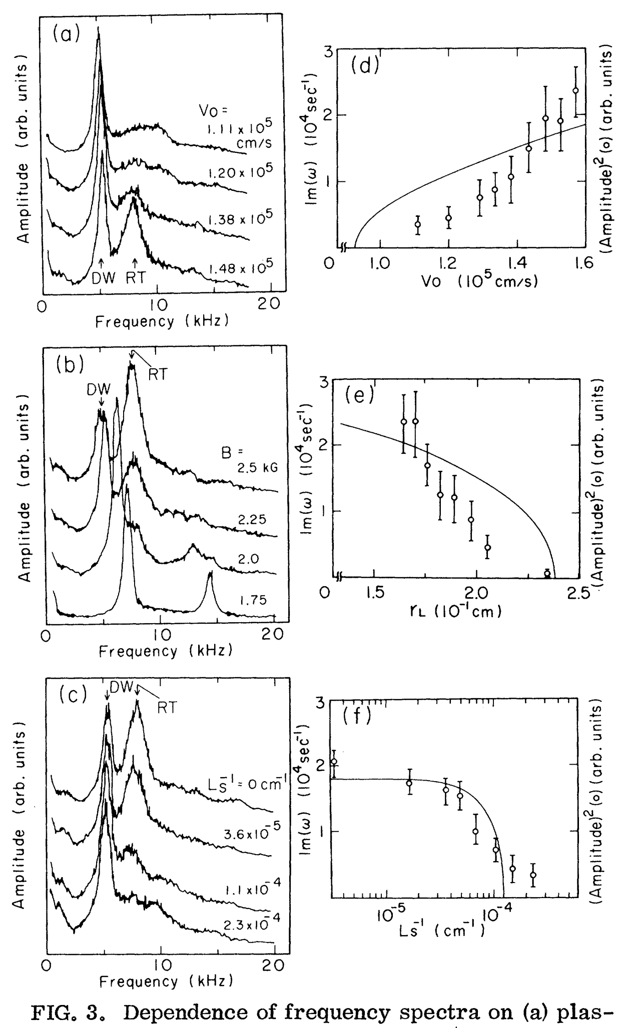

FIG. 3. Dependence of frequency spectra on (a) plasma flow speed V_0 at $B = 2.25$ kG and $L_s^{-1} = 0$, (b) magnetic field B at $V_0 = 1.57 \times 10^5$ cm/sec and $L_s^{-1} = 0$, and (c) magnetic shear length L_s at $V_0 = 1.57 \times 10^5$ cm/sec and $B = 2.25$ kG, measured by the probe P_2 in the bad region. Corresponding growth rates given by Eq. (5) are shown by the solid lines, together with the measured values of squared amplitudes (circles), in (d)–(f).

stability is sensitive to the magnetic shear even at small values of L_s^{-1}, and decreases as L_s^{-1} is increased. This suppression is expected: Magnetic shear results in $k_z \neq 0$, and electrons moving along the magnetic field over the shortened wavelength to cancel out the electric field fluctuations.[2] In our case, it is almost impossible to recognize the RT signal at $L_s^{-1} = 2.3\times 10^{-4}$, although the DW is not affected by such a small magnetic shear.[12]

The dispersion relation of the Rayleigh-Taylor instability is derived on the basis of the two-fluid model. For simplicity, a Cartesian coordinate system is used. We assume a low-β collisionless plasma with density gradient $dn_0/dx = -\kappa n_0$ under a uniform, straight magnetic field $\vec{B}$ in the z direction. Uniform gravity forces $m_e \vec{g}_e$ and $m_i \vec{g}_i$ are assumed for electrons and ions, respectively, in the x direction. Conservation equations are now solved by linearizing and looking for a solution of the form $\exp[i(-\omega t + k_y y + k_z z)]$, where k_z $(=k_y \Delta x/L_s$, Δx is the width of the instability) represents the shear effect mentioned above.[2] After simple algebra, we get two equations including perturbed density $\tilde{n}$ and potential $\tilde{\varphi}$:

$$\omega_e \frac{\tilde{n}}{n_0} + \left(-\frac{c_e^2 k_y \kappa}{\omega_{ce}} + \frac{c_e^2 k_z^2}{\omega_e}\right)\left(\frac{e\tilde{\varphi}}{KT_e} - \frac{n}{n_0}\right) = 0, \quad (1)$$

$$\left[\left(1 - \frac{c_i^2 k_y^2}{2\omega_{ci}^2}\right)^2 - \left(\frac{\omega_i}{\omega_{ci}} - \frac{c_i^2 k_y \kappa}{2\omega_{ci}^2}\right)^2\right]\omega_i \frac{\tilde{n}}{n_0} + \left(\frac{\omega_i c_i^2 k_y^2}{\omega_{ci}^2} - \frac{c_i^2 k_y \kappa}{\omega_{ci}}\right)\left(\frac{e\varphi}{KT_{i\perp}} + \frac{n}{n_0}\right) = 0, \quad (2)$$

where $\omega_e = \omega - k_y c_e^2 \kappa/\omega_{ce} - k_y g_e/\omega_{ce}$, $\omega_i = \omega + k_y c_i^2 \kappa/\omega_{ci} + k_y g_i/\omega_{ci}$, $c_e^2 = KT_e/m_e$, and $c_i^2 = KT_{i\perp}/m_i$. $V_0 k_z$ is small enough to be neglected. The derivation of Eqs. (1) and (2) is performed as in Chen's theory,[13] although no shear was considered in his case. For $\omega_i/\omega_{ci} \ll 1$, $|\kappa| r_L/\sqrt{2} \ll 1$, $T_e k_y r_L/\sqrt{2} T_{i\perp} \ll 1$, $|g| \ll |\kappa| c_i^2$ [r_L is the Larmor radius, $g = g_i + (m_e/m_i) g_e$], Eqs. (1) and (2) yield the dispersion relation of the Rayleigh-Taylor instability:

$$\omega_i^2 - \kappa k_y r_L^2 \omega_{ci}/2 + \kappa g - m_i \omega_{ci}^2 k_z^2 / m_e k_y^2 = 0 \quad (3)$$

which gives the frequency $\mathrm{Re}\,\omega$ and the growth rate $\mathrm{Im}\,\omega$:

$$\mathrm{Re}\,\omega = -\kappa k_y r_L^2 \omega_{ci}/4, \quad (4)$$

$$[\mathrm{Im}\,\omega]^2 = \kappa g - (\kappa k_y \omega_{ci}/4)^2 r_L^4 - m_i \omega_{ci}^2 k_z^2 / m_e k_y^2, \quad (5)$$

where $T_e \simeq T_{i\perp} \simeq T$ is used. The first term on the right-hand side of Eq. (5) represents gravitational excitation in the bad-curvature region ($g\kappa > 0$) of the magnetic field, and gives a stabilizing effect in the good-curvature region ($g\kappa < 0$). The second term is damping due to the finite Larmor radius of ions.[14] The effect of the shear field is given by the last term which works to suppress the instability. The growth rates given by Eq. (5) are shown by solid lines in Figs. 3(d), 3(e), and 3(f), where the measured values $\kappa \simeq 1.2$ cm^{-1}, $k_y \simeq 4.0$ cm^{-1}, $T \simeq 0.2$ eV, and $\Delta x \simeq 1.0$ cm are used. In these figures, the values of V_0, B, and L_s are varied, corresponding respectively to Figs. 3(a), 3(b), and 3(c). Since the linear growth rate should be proportional to the squared saturation amplitude of the instability, A^2,[15] the measured values of A^2 are also plotted. Good agreements are found between the calculated growth rates and A^2 vs V_0, B, and L_s. The frequency calculated from Eq. (4), 7.7 kHz, is consistent with the measured value, 7.4 kHz, when the frequency shift caused by the $\vec{E}_0 \times \vec{B}$ drift is included.

Our work demonstrates clear generation of the Rayleigh-Taylor instability in a plasma flowing along a curved magnetic field. The instability is controlled by changing the flow speed, ion Larmor radius, and magnetic shear. All features of the instability are well explained theoretically. Finally, this work provides important information for magnetohydrodynamic interchange modes closely related to the Rayleigh-Taylor instability, which were recently discussed for instabilities in a tokamak device.[1]

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[1] A. H. Glasser *et al.*, Phys. Fluids 19, 567 (1976); B. Coppi, Phys. Rev. Lett. 39, 939 (1977); D. Dobrott *et al.*, Phys. Rev. Lett. 39, 943 (1977).

[2] A. B. Mikhailovskii, *Theory of Plasma Instabilities* (Consultants Bureau, New York–London, 1974), Vol. 2, Chaps. 6 and 8.

[3] A. C. Kolb *et al.*, Phys. Fluids 8, 1005 (1965); E. Hintz and A. C. Kolb, Phys. Fluids 8, 1347 (1965); Yu. B. Gott *et al.*, Nucl. Fusion Suppl. Pt. 3, 1045 (1962).

[4] N. D'Angelo *et al.*, Phys. Rev. Lett. 11, 525 (1963).

[5] F. F. Chen, Plasma Phys. 7, 399 (1965).

[6] R. W. Morley, *Q Machines* (Academic, New York, 1975).

[7] A. Komori *et al.*, Plasma Phys. 19, 283 (1977).

[8] N. Sato *et al.*, Phys. Rev. Lett. 34, 931 (1975).

[9] P. F. Little and C. R. Middleton, Nucl. Fusion 9, 67 (1969).

[10]H. Sugai *et al.*, Phys. Fluids 20, 90 (1977).
[11]S. A. Andersen *et al.*, Phys. Fluids 14, 728 (1971).
[12]F. F. Chen and D. Mosher, Phys. Rev. Lett. 18, 639 (1967); P. E. Scott *et al.*, Phys. Rev. Lett. 25, 996 (1970).
[13]F. F. Chen, Phys. Fluids 8, 1323 (1965).
[14]M. N. Rosenbluth and A. Simon, Phys. Fluids 8, 1300 (1965).
[15]L. D. Landau and E. M. Lifshitz, *Fluid Mechanics* (Pergamon, London, 1960), p. 104.

Suppression of ω_{ci} Instability in a Mirror-Confined Plasma by Injection of an Electron Beam

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A hot ion plasma ($\bar{E}_{i\perp} \simeq 400$ eV, $n \simeq 10^{13}$ cm^{-3}) is trapped in a minimum-B mirror system. A strong instability at ω_{ci} is observed for neutral densities $< 10^{11}$ cm^{-3}. We have observed the suppression of this instability when an electron beam (≥ 42 kW) is injected into the plasma. We believe the suppression is due to a population of hot electrons produced and trapped in the mirror by the beam-plasma interaction.

Mirror-confined plasmas may exhibit strong microinstabilities driven by the inverted nature of the particle distribution.[1] Recently, it has been shown that the hole in velocity space can be filled by injection of a warm plasma stream. This has resulted in nearly classical loss times of hot, highly ionized plasmas.[2] A problem with this technique is that the warm plasma cools the electrons of the confined plasma. Ioffe *et al.*[3] describe a different method of populating the cold part of the velocity distribution. They use electron-cyclotron resonance heating to produce a small population of energetic electrons near the mirror midplane with $v_\perp \gg v_\parallel$. These are magnetically trapped and produce a local depression of the space potential. It is believed that this inhibits the plasma loss and helps to fill the velocity distribution. In their experiments they were able to suppress the drift cyclotron loss-cone instability.

We report here an alternative scheme for suppressing these instabilities, by creating a trapped hot electron population by the interaction of an electron beam and the main plasma. Earlier studies[4,5] of electron beam-plasma interactions have shown that a small population of very hot electrons ($\gtrsim 20$ keV) with $v_\perp \gg v_\parallel$ is produced when high-power electron beams are injected into mirror-confined plasmas.

In order to produce a loss-cone distribtuion that will drive instabilities, the ions must have a mean free path $\gtrsim 100$ mirror bounce times. This requires $\bar{E}_{i\perp} > 100$ eV and neutral density $n_0 < 10^{11}$ cm^{-3}. These conditions have been met in our apparatus, Fig. 1. The mirror is of the minimum-B type, with the hexapole field produced by ferrite permanent magnets. The axial field is produced by an array of fifteen coils that form the confining mirror and the guide field for the plasma injected by the Ti-washer plasma gun. A series of baffles and Ti getter pumps keep neutral gas from the gun from reaching the mirror, so that $n_0 < 10^{11}$ cm^{-3} during the injection and containment periods of the experiment. The plasma gun pulse length may be varied between 0.1 and 2 msec. To prevent residual plasma from the gun entering the mirror after "turnoff," a divertor coil may be energized that magnetically disconnects the mirror and the gun. This ensures that the plasma will not be stabil-

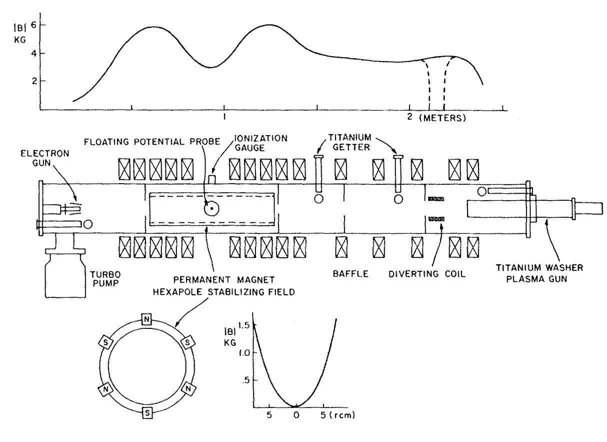

FIG. 1. Experimental apparatus for studying loss-cone microinstabilities.