

PLASMA HEATING AND CONFINEMENT
WITH
MODE-CONTROLLED ION-CYCLOTRON RF FIELDS

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APRIL 1984

DEPARTMENT OF ELECTRONICS
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PREFACE

The electromagnetic interactions of radio frequency (RF) waves with plasmas have been basic and important problems in plasma physics. The interactions are also extensively investigated in conjunction with thermonuclear fusion research.

The purpose of the present thesis is to study experimentally the plasma heating and confinement with the application of RF waves. The main technique we employ is to use mode-controlled RF fields for the excitation of the RF waves. By properly choosing the mode of the applied RF fields, we can achieve the radial confinement of a plasma in a single axisymmetric mirror through the stabilization of flute instabilities, the improvement of the heating efficiency in ion cyclotron heating through the elimination of nonlinear convective plasma loss, and the improvement of the axial confinement in RF-plugged open ended systems.

The experimental results and techniques developed in this thesis can be readily applicable to larger-scale experiments, and in fact, have been successfully put into use in several present-day devices.

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CHAPTER 1

INTRODUCTION

1. ROLES OF ICRF WAVES IN MAGNETIC CONFINEMENT SYSTEMS

Radio frequency (RF) wave applications in magnetic confinement systems have held an important position in this decade in researches of controlled thermonuclear fusion. RF waves can directly couple with charged particles providing energy deposition and/or momentum transfer in a desired position in the phase-space of particles. This enables us to utilize the RF waves as an efficient method for plasma heating and confinement. Among several frequency ranges, it may be the ion cyclotron range of frequencies (ICRF) that is used most extensively for additional heating and controls of plasmas in tokamaks and open ended systems. A MW-level RF power in ICRF was successfully coupled into a plasma to achieve an efficient heating up to a few keV in tokamaks [1]. Reduction of high-Z impurities through a control of temperature profile by ICRF heating was investigated [2]. Studies of toroidal current generation by waves in ICRF as well as by lower hybrid and electron cyclotron waves are also performed [3].

For open ended systems, RF waves seem to have much more potentially important uses in viewpoints of RF heating and confinement of plasmas. Since the confinement of a plasma with fusion-oriented parameters was demonstrated in the mirror

machine 2XIIB [4], investigations of open ended systems have become active, and recently, large mirror machines based on a new concept of ambipolar potential confinement were constructed [5-7]. Although such a mirror machine, the tandem mirror, greatly improves the plasma confinement, and hence, the Q value (energy amplification factor) over a simple mirror, there are still problems such as resonant radial diffusion of particles due to asymmetric configuration, large gas load and technological restriction associated with high-energy neutral beam injections (NBI) at plug cells, needs for an efficient heating of a central cell plasma, and so on. Use of RF waves may resolve these problems, and in fact, is being investigated in detail. For example, MHD-stable radial confinement of a tandem mirror plasma may be achieved by RF stabilization [8] instead of non-axisymmetric MHD anchors, and the high-energy NBI can be replaced by ICRF heating and sustainment of plug plasmas [9]. It is also predicted that the pumping of cold ions in a thermal barrier or sloshing ion formation can be performed by RF waves [10].

Besides the tandem mirror, the RF wave applications have played an important role in Radio Frequency Containment (RFC) experiments [11-13]. Large amplitude ICRF fields applied at open ends produce the ponderomotive potential to confine particles axially just as the ambipolar potential in the tandem mirror. This configuration, the RF plugging, requires an efficient method for creating the ponderomotive potential.

The roles of ICRF waves described above are characterized as energy and momentum transports from waves to particles

through particle acceleration and deceleration by Lorentz force and, for large amplitude waves, ponderomotive force. For a cylindrical plasma surrounded by a conducting wall, which is the most general situation in real experiments, the wave characteristics, e.g., field distribution and polarization are strongly dependent on geometrical parameters as well as plasma parameters. The magnitude, the direction, and their spatial distribution of the Lorentz force and the ponderomotive force are dependent on such a mode structure of the RF field. Hence, in order to realize efficient RF heating and confinement, it is important to select RF mode so that desired energy and/or momentum transports take place in a desired position of phase-space. Stix [14] was the first who used the antenna which selectively excites a specific mode. He chose $m = 0$ mode for ion cyclotron heating, where m is the azimuthal mode of RF fields in a cylindrical coordinate. Several experimental reports by Russian authors are available concerning the effect of RF mode selection in heating and stabilizing plasmas. Diky et al. [15] showed that the $m = \pm 1$ mode (linear combination of $m = +1$ and $m = -1$ modes) was better than the $m = 0$ mode for heating the core of a plasma in a circular torus. Demirkhanov et al. [16] used a phased helical antenna to excite Alfvén wave of $m = +2$ mode selectively. The plus sign of m means that the RF field rotates macroscopically in left-hand sense with respect to the axial static magnetic field. It was also reported that a plasma in a linear solenoid was formed into a column when the RF field (not a wave) of $m = -1$ mode with a frequency much lower than the ion cyclotron frequency was

applied [17]. Hipp et al.[18] calculated the loading resistance of a phased helical antenna when launching $m = 0, +1$, or -1 wave in ICRF. Knox et al.[19] demonstrated the selective excitation of $m = +1$ slow wave and $m = -1$ fast wave in ICRF using the phased helical antenna. Although the importance of RF mode selection has been recognized in several experiments mentioned above, it seems that little attention has been paid so far to systematic selection of RF mode in viewpoints of optimization of Lorentz force and/or ponderomotive force.

In the present thesis, we present the experimental studies on plasma heating and confinement with mode-controlled RF fields in ICRF. The experiments described here are RF stabilization of MHD instabilities for radial confinement in an axisymmetric mirror, observation of a new type plasma loss associated with the ICRF heating and its elimination for an improvement of the heating efficiency, and efficient RF plugging for axial confinement in open ended systems. The main idea we employ is to control the azimuthal mode of applied RF waves as well as the selection of wave branches for the purpose of efficient RF controls of plasma transports. In the following Sections, historical backgrounds and critical issues of these subjects are described.

2. ALTERNATIVE METHOD FOR MHD ANCHORS IN A TANDEM MIRROR

Mirror trap is one of the most common configurations in open ended systems. In 1975, it was demonstrated that

a plasma with fusion-oriented parameters was created in the mirror machine 2XIIB with high-energy NBI and warm plasma stabilization of micro instabilities [4]. It is well known that the mirror trap has several advantages over a tokamak such as a simpler configuration with good accessibility, steady state operation, confinement time independent of machine size, high- β operation, possible use of a direct energy converter, and so on. However, the most serious problem that prevented the mirror from being a fusion reactor was the low Q value due to losses through open ends even if a direct converter with very much high efficiency was incorporated. A new concept "ambipolar potential confinement" was proposed by Dimov et al. [20] and Fowler and Logan [21] as a method for enhancing the Q value considerably. Based upon the concept, tandem mirrors were constructed in Tsukuba Univ. and later in LLNL and Univ. of Wisconsin. The tandem mirror consists of a long central cell and minimum-B end plug cells in both ends. Ions in the central cell is confined by an ambipolar potential created in plug cells with high-energy NBI. This configuration was experimentally confirmed to work well in GAMMA6 [5], TMX [6], and Phaedrus [7]. It was further shown that the ambipolar potential could be created with less plasma density in plug cells and less NBI power provided that the thermal contact between electrons of the central and plug cells was broken down by a thermal barrier [22]. The experiment of the thermal barrier is now being carried out in TMX-U [23]. Rapid progress of the tandem mirror research enables us to ensure enough longitudinal confinement of a plasma. It has

been, however, pointed out that the resonant transverse diffusion due to non-axisymmetric configuration degrades the radial confinement significantly. In fact, although the original tandem mirror had minimum-B end plugs, they were axisymmetrized and only anchor regions remain asymmetric in present-day and future machines [24-26]. The non-axisymmetric anchor regions are necessary to maintain magnetohydro-dynamic (MHD) stability of a whole plasma. Although sophisticated designs of the anchors have been made to minimize the resonant diffusion, it is impossible to completely eliminate the enhanced diffusion in this configuration. It is, therefore, important to investigate other methods for maintaining MHD stability in order to realize a purely axisymmetric tandem mirror.

The most deleterious MHD instability in an axisymmetric mirror is the flute mode [27]. This instability is driven by a radial force due to "bad" curvature of the magnetic field line. Ions and electrons drift oppositely with each other in azimuthal direction by the radial force so called the effective gravity. If a small perturbation occurs at a plasma-vacuum interface, a charge separation is created and a resultant electric field acts as to enhance the perturbation through the $E \times B$ drift. The plasma is then lost radially.

If we apply an RF wave (or field) in such a way that the ponderomotive force of the RF field acts on ions or electrons oppositely to the effective gravity, the relative azimuthal drift of ions and electrons is eliminated and the plasma becomes stable against the flute perturbation [28-30]. The experiment of stabilizing a precessional drift wave using

this method was done by Yamamoto et al.[31]. In the succeeding paper [32], they showed flute stabilization by an RF field which was generated electrostatically by rod electrodes aligned to cover almost the entire plasma surface. In their experiment, however, the observed instability was not clearly identified. Furthermore, the density was restricted to 10^{12}cm^{-3} , so the availability of this stabilization method was not confirmed for higher densities where the electrostatic RF field excitation would become inefficient. Yasaka et al.[33] showed the experimental identification of the curvature-driven flute instability in a density range higher than 10^{13}cm^{-3} . They also demonstrated the stabilization of the flute instability by the RF field which was electromagnetically coupled and applied only in the bad-curvature region in a single axisymmetric mirror [34]. The stabilization method was confirmed to be effective in a wide range of density up to 10^{14}cm^{-3} .

To clarify the utility of the RF stabilization in a fusion plasma, it is necessary to obtain a scaling relation of required RF field strength to plasma parameters. An empirical scaling of the RF stabilization in a single axisymmetric mirror was obtained [35]. The results of these experiments suggest that the RF stabilization could be an alternative for MHD anchors of a tandem mirror.

3. ICRF HEATING BY ROTATING RF FIELDS

ICRF heating is divided into two categories; one is ion cyclotron heating (ICH) in which ion cyclotron (slow) waves are absorbed by ions through fundamental ion cyclotron damping, and the other is fast wave heating in which magnetoacoustic (fast) waves are absorbed through second harmonic damping in one ion species plasma or are absorbed via mode conversion near the ion-ion hybrid resonance layer in two ion species plasma. Ion cyclotron heating was first proposed by Stix [14] and was investigated in B-65, B-66, and C-stellarators in PPPL [36,37]. An azimuthally symmetric slow wave was excited by a Stix coil at a uniform magnetic field section, propagated toward weaker field side, and damped at a magnetic beach. In the local magnetic beach, ion temperature reached 3 keV. Although the wave was efficiently coupled and absorbed in the linear section of the racetrack torus, it could not propagate around whole of the torus.

In the Uragan stellarator, an azimuthally asymmetric slow wave of $m = \pm 1$ mode was launched by a Kharkov coil and propagated around the torus to bring a uniform heating of a whole plasma [15,38]. The asymmetric wave had one more advantage over the symmetric wave that the perpendicular component of the RF field was maximum on minor axis, hence the wave could heat a core of the plasma.

In circular machines such as ST tokamak [39] or Heliotron [40,41], the slow wave was propagated perpendicularly to the toroidal magnetic field with a fixed axial wave length. After-

wards, the density of tokamak plasmas became so high that the slow wave could not propagate into the core unless the axial wave length of the wave was chosen to be unreasonably short. Hence, the fast wave came into use for RF heating in tokamaks.

The ICRF heating was investigated also in open ended systems. The azimuthally symmetric slow wave was used in THM-2 mirror machine and plasma heating at first and second harmonic ion cyclotron frequencies was observed [42,43] . In TPD-III [44], plasma heating by the $m = +1$ slow wave was compared with that by the $m = 0$ slow wave. A detailed experiment concerning the mode effect on antenna loading was performed by Knox et al.[19] using $m = +1$ and $m = -1$ RF fields. They indicated that a proper selection of the direction of rotation might increase the efficiency of wave excitation.

It is of course desirable not only to achieve higher heating efficiency but to suppress enhanced plasma loss during ICRF heating, which was observed, for example, in ATC [45], Proto-Cleo [46], and Macrotron [47]. Since the excursion velocity greatly exceeds the thermal velocity in RF heating, nonlinear effects due to, for example, a ponderomotive force will become important. Although nonlinear processes in lower hybrid heating have been investigated in detail [48,49], few experimental reports are available for ICRF heating.

Yasaka and Itatani showed that the conventional slow wave of $m = +1$ mode gave rise to a new type of plasma transport leading to a reduction of plasma confinement in a linear device [50]. The transport is produced through a radial convective motion of ions caused by the nonlinear interaction of $m = +1$

and $m = -1$ RF fields. It was found that the anomalous transport could be eliminated if the slow wave was of $m = +1$ rotating mode. The use of the $m = +1$ mode improves the overall heating efficiency by a factor of 1.7 [50]. The improvement of heating efficiency by the use of the slow wave of $m = +1$ rotating mode was also confirmed in toroidal [51] and mirror plasmas. These experiments have first established the superiority of using the rotating RF field in slow wave heating of both linear and toroidal plasmas.

Adati et al.[52] performed ICH experiments in the double cusp machine RFC-XX following in the wake of our experiments.

They changed the mode of the slow wave, from non-rotating to rotating mode continuously, and found that the best heating occurred in the case of the $m = +1$ rotating mode. This result also supports our conclusion.

4. IMPROVEMENT OF RF PLUGGING

Low Q value of open ended systems requires some efficient methods for reducing end losses. The principal method is the ambipolar potential confinement, which is explained in Section 2. The other method, which has longer history of investigation, is the RF plugging of end losses. If an RF field of appropriate mode and frequency is applied at an open end, charged particles are repelled along the static magnetic field by a ponderomotive potential created by the applied RF field. The ponderomotive potential confinement has an

outstanding advantage over the ambipolar potential confinement that the ions with a specific charge to mass ratio can be plugged selectively. This leads to a possibility that, in a reactor, only fuel ions are confined selectively and ash or impurity ions are automatically exhausted from the confinement region.

The RF plugging can be applied to linear solenoids, mirror machines, and cusp machines. In the course of researches in IPP, Nagoya, the effectiveness of plugging the cusp plasma has been established [53-56]. The loss flux through a line cusp was reduced by a ponderomotive potential of electrostatic ion cyclotron RF fields launched by parallel plate electrodes. The loss flux could be reduced to 5 % of the initial value with 800 V/cm RF electric field in a plasma of density of $1.3 \times 10^{12} \text{ cm}^{-3}$ [57]. The ability of plugging a specific ion species was also demonstrated [58].

For a point cusp or a mirror end, an electromagnetic coupling is used to create a ponderomotive potential in a larger-cross-section, higher-density plasma compared to the line cusp plasma. In the experiment performed by Watari et al. [59], it was shown that the azimuthally asymmetric RF fields produced by the type-III coil was much more effective for end loss reduction than the symmetric RF fields in the sense that the required RF electric field strength remained relatively small for higher densities. They also found that the plugging effect was due to the axial ponderomotive potential for ions, and electrons were confined by a resultant ambipolar potential. Since the perpendicular acceleration of ions by RF fields is

responsible for RF plugging, it is expected that a use of left-hand circularly polarized RF fields of $m = +1$ mode will enhance the end loss reduction compared to the linearly polarized RF fields of $m = \pm 1$ mode used in TPD-III.

Yasaka et al.[60] showed that the plugging effect was greatly improved when the $m = +1$ RF fields were employed instead of the $m = \pm 1$ RF fields. The particle and energy confinement times in an RF-plugged simple mirror were compared for different RF modes. The best confinement occurred when the $m = +1$ RF fields were used.

In the simple mirror, the flute instability was observed resulting in significant reduction of the confinement time. In this case, both the RF stabilization and the RF plugging are necessary for confinement. The simultaneous application of the RF stabilization and the RF plugging was reported [61].

The experiments of RF plugging in a different frequency range [62] or in a different magnetic field configuration [63,64] were also performed successfully although the density and the temperature used were relatively low.

In order to extend the results in TPD-III and to obtain prospect for using the RF plugging in a reactor, Nagoya group constructed the double cusp machine RFC-XX, which consisted of a central solenoid section and two cusp regions in both ends. By applying the RF voltage to the parallel electrodes located at two line cusps, the end loss flux through the line cusps was reduced to less than several % of the initial value. The energy confinement time of the central-section plasma is improved by a factor of 2 through the line cusp plugging

[60]. The end loss flux through the point cusps was about one-tenth of that through the line cusps, and hence had minor effects on total plasma confinement. However, once almost complete end loss reduction is achieved at the line cusps, an efficient RF plugging of the point cusps is necessary for further improvement of the energy confinement time. Yasaka et al.[65] applied the asymmetric, rotating RF fields to the point cusp plugging, according to the results in the simple mirror, and found that this type of RF field had better plugging effect than the non-rotating RF field.

The RF plugging of ICRF-heated plasma was also performed [65]. The streaming plasma with ion temperatures up to 250 eV could be plugged by the rotating RF field of $m = +1$ mode with a voltage less than 20 kV. A scaling formula of the RF plugging of the ICRF-heated plasma was obtained experimentally for the first time.

5. FRAMEWORK OF THE PRESENT THESIS

The present thesis is organized as follows:

In Chapter 2, basic properties of RF waves in ICRF in a cylindrical plasma are explained. It is described how the effects of Lorentz force and the ponderomotive force on energy and momentum transports depend upon the azimuthal mode of the RF fields. Forced excitation of RF fields and a method for controlling the azimuthal mode of RF fields are also discussed.

RF stabilization and associated radial confinement in a single axisymmetric mirror is presented in Chapters 3 and 4. In Chapter 3, low frequency instability observed in the mirror is identified as a curvature-driven flute instability for a wide range of densities up to 10^{14}cm^{-3} . It is demonstrated that the flute instability is stabilized by the RF field which is electromagnetically coupled and applied only in the bad-curvature region. Radial loss of the plasma is reduced significantly through the stabilization. This method of stabilization is revealed to be effective in densities from 10^{10}cm^{-3} to 10^{14}cm^{-3} .

In Chapter 4, the RF stabilization of the flute instability is investigated by varying the plasma density, the radius of curvature of the magnetic field line, and the ion temperature. It is described that a scaling formula of the RF stabilization which relates the required RF electric field strength to plasma parameters is obtained, predicting a reasonable value of 1.1 kV/cm for the stabilization of a fusion-grade plasma. The results of Chapters 3 and 4 suggest that the RF stabilization could be an alternative for MHD anchors of a tandem mirror.

Chapters 5 and 6 contain the studies on improvement of the heating efficiency of ICRF heating. In Chapter 5, we present that a conventional slow wave of $m = \pm 1$ mode gives rise to a new type of plasma transport leading to a reduction of radial confinement. The mechanism of the anomalous transport is investigated both theoretically and experimentally. It is found that the anomalous transport can be eliminated

if the slow wave is of rotating mode. The use of $m = +1$ rotating mode improves the overall heating efficiency by a factor of 1.7.

In Chapter 6, the slow wave of rotating mode is applied to ICRF heating of toroidal and mirror plasmas. It is also confirmed that the $m = +1$ rotating mode can reduce the anomalous loss leading to an improvement of the heating efficiency. These experiments have first established the superiority of the $m = +1$ RF fields in slow wave heating of linear and toroidal plasmas.

Chapter 7 contains the results on improvement of RF plugging in open ended systems. The particle and energy confinement times in an RF-plugged axisymmetric mirror are compared for different RF modes. The best confinement is achieved when the $m = +1$ RF field is applied for plugging.

The simultaneous application of the RF stabilization and RF plugging are also described. It is observed a combined effect that the plugging efficiency is enhanced with less total RF power as compared with the case of RF plugging only.

In the latter half of Chapter 7, the RF plugging of ICRF-heated plasma is performed at a point cusp. The streaming plasma with ion temperatures up to 250 eV can be plugged by the $m = +1$ rotating RF field. A scaling formula of the RF plugging of the ICRF-heated plasma is obtained for the first time.

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CHAPTER 2

MODE CONTROLLED ELECTROMAGNETIC FIELD IN ION CYCLOTRON RANGE OF FREQUENCIES

1. INTRODUCTION

Radio frequency waves in Ion Cyclotron Range of Frequencies (ICRF) have been used with growing importance in researches of additional heating and dynamic control of plasmas. There have been reported a number of theoretical studies concerning the wave coupling, propagation, and absorption in infinite or slab plasmas, while such investigations in a cylindrical, bounded plasma, which is the most general situation in real experiments, seem to be not enough. For a cylindrical plasma surrounded by a conducting wall, the wave characteristics, e.g., field distribution and polarization, are strongly dependent on geometrical parameters including antenna structures as well as plasma parameters. For example, it was shown that the cutoff frequency of a compressional Alfvén wave (or simply termed as a fast wave) significantly changes when the thickness of vacuum layer between the cylindrical plasma and the wall is changed [1].

The cylindrical plasma can support a variety of eigenmodes, each of which has a different characteristics of propagation. Stix [2] proposed a "Stix" coil to generate an ion cyclotron wave of azimuthally symmetric mode and to heat

a plasma via ion cyclotron damping of the wave. The use of the Stix coil was successful in the ion cyclotron heating experiments performed in the straight section of C-stellarator [3]. But the wave did not propagate through toroidal sections, and hence, did not bring a heating of the whole plasma. On the other hand, the same wave but of azimuthally asymmetric mode was experimentally revealed to propagate freely in a toroidal plasma [4]. These extremely different characteristics come from only the difference in the azimuthal mode of ion cyclotron waves.

Knox et al.[5], used a phased herical antenna to excite azimuthally asymmetric Alfvén waves in a linear machine. They showed that the antenna can selectively couple to either the slow wave or the fast wave by controlling the phase difference of RF currents in the antenna, i.e., the azimuthal mode of RF current distribution.

As described above, the propagation of waves in ICRF in a cylindrical plasma is strongly dependent on the azimuthal mode. It is, therefore, easily predicted that plasma transport phenomena are also strongly affected by the azimuthal mode of the wave, while few experimental reports in this standpoint are available so far. Before describing the experimental results which show the very evidence of the prediction, we here present in Section 2, a brief theoretical consideration of the dispersion and mode structure of waves in ICRF.

When a plasma is subjected to the large amplitude RF fields, nonlinear effects become important for plasma transports. As an example of positive utilization of these effects,

we can point out a series of studies on RF plugging of a cusp plasma [6-8]. Although the mechanism of plugging in relation to the axial ponderomotive force was investigated in detail by many authors, the effect of other components of the ponderomotive force has been looked over. In Section 3, the ponderomotive force is derived and effects of each component on plasma transports are discussed.

Forced excitation of RF fields in ICRF is accomplished by an RF antenna outside a plasma column. Energy and momentum transports due to the wave fields are dependent on mode structure. So, to optimize a desired transport, the RF antenna must be designed so that it generates coherent mode spectrum. The antenna design is discussed in Section 4.

2. EIGENMODE IN A CYLINDRICAL PLASMA

The behavior of waves in a plasma is governed by the wave equation given by

$$\vec{\nabla} \times \vec{\nabla} \times \vec{E} + \frac{\omega^2}{c^2} \overleftrightarrow{K} \vec{E} = 0, \quad (2-1)$$

where $\vec{E}$ is the electric field which varies as $\exp[j(-k_{\parallel}z + \omega t)]$ with $k_{\parallel}$ and ω being the parallel wave number and the angular frequency, respectively, c is the velocity of light, and $\overleftrightarrow{K}$ is the dielectric tensor. As far as the propagation of waves is concerned, it is sufficient to use the cold plasma approximation for ions. But, for electrons, we consider

both cases of hot and cold electrons. For a collisionless, homogeneous plasma immersed in a uniform static magnetic field in z-direction, the dielectric tensor [9] takes the form,

$$\vec{\epsilon} = \begin{pmatrix} S & jD & 0 \\ -jD & S & 0 \\ 0 & 0 & P \end{pmatrix}, \quad (2-2)$$

where

$$S = \frac{1}{2} (R + L), \quad D = \frac{1}{2} (R - L),$$

$$R = 1 - \sum_{\sigma} \frac{\omega_{p\sigma}^2}{\omega^2} \left(\frac{\omega}{\omega + \omega_{c\sigma}} \right), \quad (2-3)$$

$$L = 1 - \sum_{\sigma} \frac{\omega_{p\sigma}^2}{\omega^2} \left(\frac{\omega}{\omega - \omega_{c\sigma}} \right),$$

$$P = 1 - \sum_{\sigma} \frac{\omega_{p\sigma}^2}{\omega^2}, \quad (\text{cold electrons})$$

$$1 - \frac{\omega_{pi}^2}{\omega^2} + \frac{1}{k_{\parallel}^2 \lambda_D^2}, \quad (\text{hot electrons})$$

$\omega_{p\sigma}$ and $\omega_{c\sigma}$ denote the plasma and the cyclotron frequencies of the σ -th species, respectively, and λ_D is the Debye length. The explicit form of the wave equation for E_z is written as

$$[\nabla_T^4 + (a + c) \nabla_T^2 + (ac - bd)] E_z = 0, \quad (2-4)$$

where

$$a = (-k_{\parallel}^2 + k_0^2 S)P/S, \quad b = j\omega\mu_0 k_{\parallel} D/S,$$

$$c = -k_{\parallel}^2 + k_0^2 R L / S, \quad d = -j \omega \epsilon_0 k_{\parallel} D P / S,$$

$k_0 = \omega/c$, and μ_0 and ϵ_0 are the permeability and the permittivity in vacuum, respectively [10].

If we further assume that $\vec{E}$ is proportional to $\exp(-j\vec{k}_{\perp}\vec{r})$, Eq.(2-4) reduces to second order differential equation with two eigen values $k_{\perp 1}$ and $k_{\perp 2}$. The eigen values must satisfy the dispersion relation:

$$k_{\perp}^4 + (a + c) k_{\perp}^2 + (ac - bd) = 0. \quad (2-5)$$

In infinite plasmas, the propagation of waves is well explained by Eq.(2-5). Figure 2-1 shows the dispersion relation on $k_{\perp}$ versus ω_{pi}^2/ω^2 plane for a frequency of $\omega/\omega_{ci} = 0.9$. The other parameters are $k_{\parallel} = 21 \text{ m}^{-1}$ and $\omega/2\pi = 3.5 \text{ MHz}$. (Hereafter, if otherwise noted, the plasma contains single ion species of hydrogen.) It is seen that for the slow wave, the propagation is limited up to a density where $(k_{\parallel}c/\omega)^2 = S$ in the case of cold electrons, while for hot electrons with the temperature of $T_e = 30 \text{ eV}$, the wave can propagate into higher densities. Fast wave appears beyond a density where $(k_{\parallel}c/\omega)^2 = R$. The same plot is given in Fig.2-2 for $\omega/\omega_{ci} = 2.4$, $k_{\parallel} = 10.5 \text{ m}^{-1}$, and $\omega/2\pi = 10 \text{ MHz}$. Slow wave becomes evanescent for a cold electron plasma. There is no essential change in fast wave propagation as compared with Fig.2-1.

We then consider the case where a cylindrical homogeneous plasma with radius p is surrounded by a perfectly conducting cylindrical wall with radius w as shown in Fig.2-3. An RF

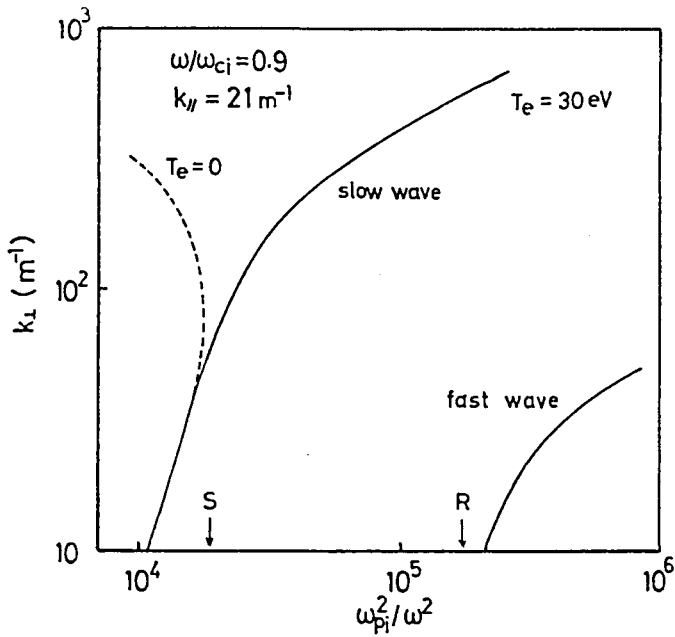


Fig.2-1 Dispersion relation of fast wave, slow wave with $T_e=0$, and slow wave with $T_e=30 \text{ eV}$ for $f = 3.5 \text{ MHz}$, $f/f_{ci} = 0.9$, and $k_{||} = 21 \text{ m}^{-1}$. The value of $(k_{||}c/\omega)^2$ is equal to $S(R)$ at the density labeled $S(R)$.

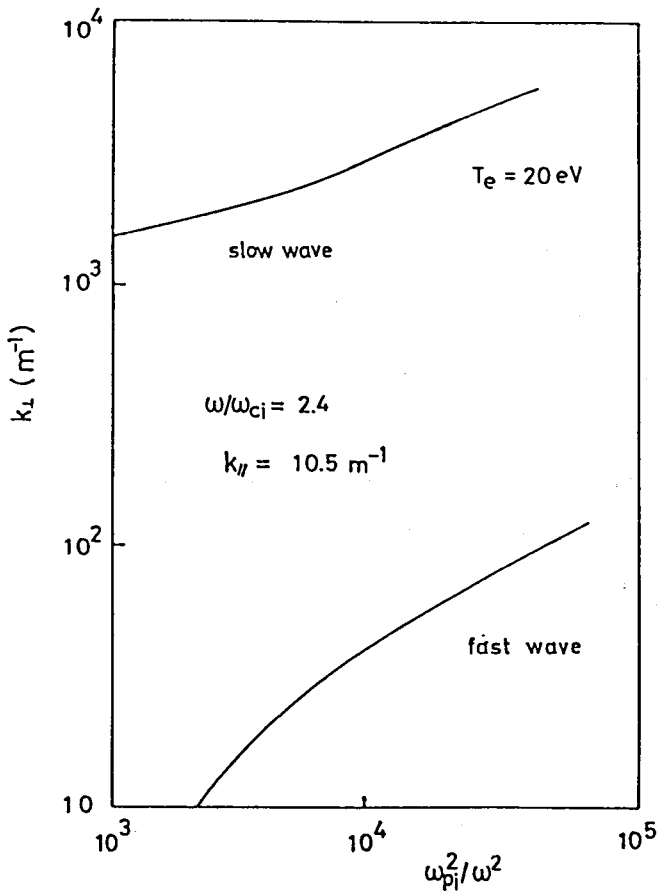


Fig.2-2 Dispersion relation of fast and slow waves for $f = 10 \text{ MHz}$, $f/f_{ci} = 2.4$, and $k_{||} = 10.5 \text{ m}^{-1}$. Slow wave is evanescent for cold electrons.

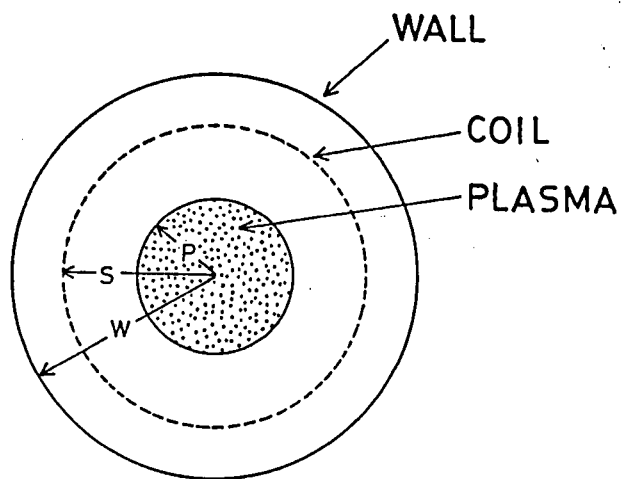


Fig.2-3 Cross section of a bounded cylindrical plasma.

antenna is approximated by a current sheet at $r = s$, and the static magnetic field is again in z -direction. The solution of Eq.(2-4) is given by

$$E_z = \sum_{\ell=1}^2 A_{\ell} J_m(k_{\perp \ell} r) \exp[j(m\theta - k_{\parallel} z + \omega t)], \quad (2-6)$$

where m is the azimuthal mode number, and $k_{\perp \ell}$ must satisfy Eq.(2-5). The other components of the electric fields and the magnetic fields are obtained from Maxwell's equations.

Restarting from Eq.(2-4) with $\omega_{ps} \rightarrow 0$, we can also obtain the fields outside the plasma in the same manner. The results are summarized in Appendix A. These fields must satisfy boundary conditions at plasma-vacuum and vacuum-wall interfaces, which are given by

$$\overleftrightarrow{T} \vec{A} = \vec{J}, \quad (2-7)$$

where $\vec{A}$ is a constant vector of field amplitudes, $\vec{J}$ is the RF current density in the antenna, and matrix $\overleftrightarrow{T}$ is defined in Appendix A. For an eigenmode (finite $\vec{A}$ with infinitesimal $\vec{J}$), a non-trivial solution requires

$$\det \overleftrightarrow{T} = 0. \quad (2-8)$$

The dispersion in the cylindrical plasma is given as a simultaneous solution of both Eqs.(2-5) and (2-8). The ω - $k_{\parallel}$ diagram of fast and slow waves in ICRF is shown in Fig.2-4 for $m = -1$ and $+1$ modes. The plasma has following parameters; density

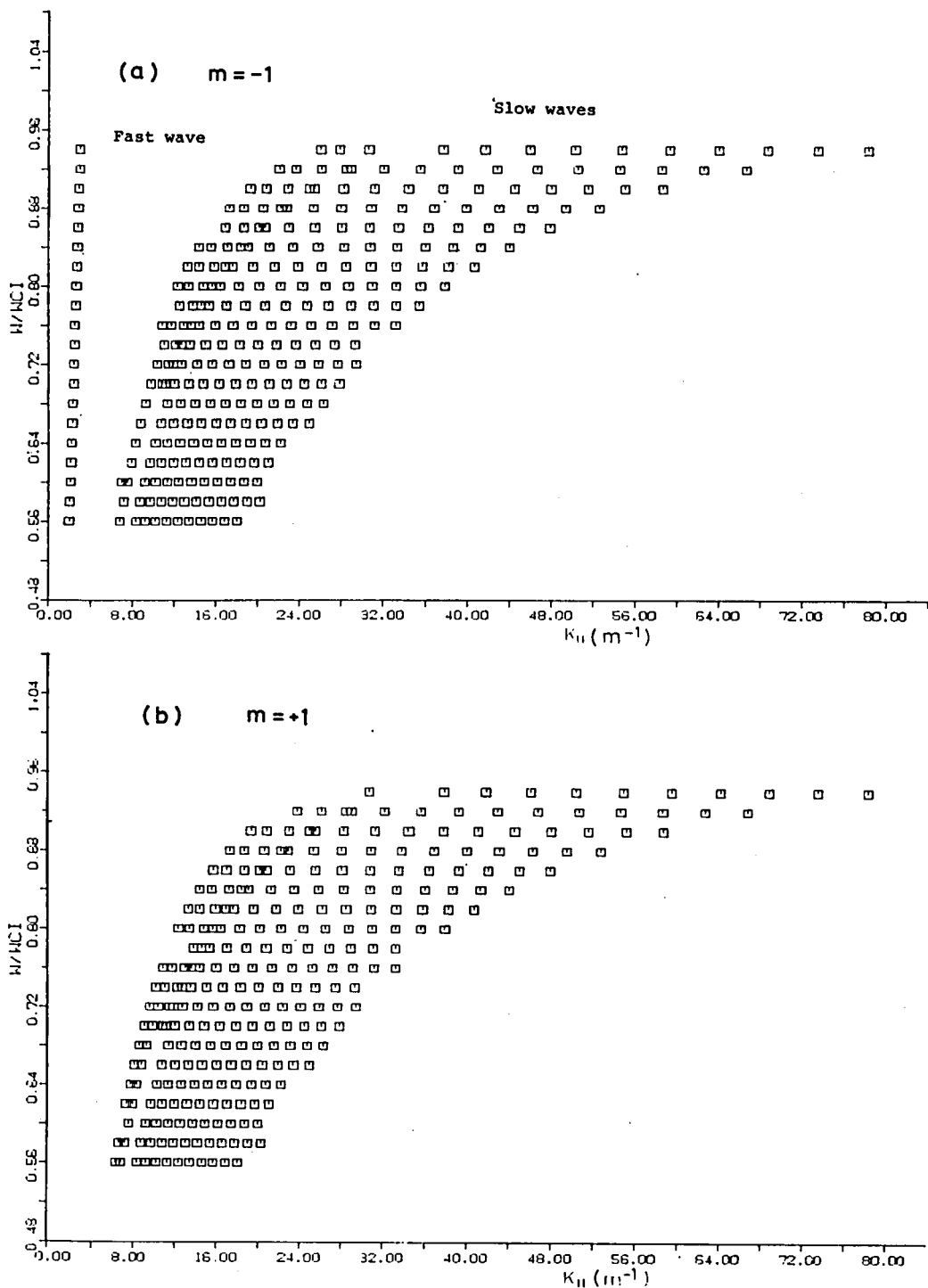


Fig.2-4 ω/ω_{ci} versus $k_{||}$ plot of (a) fast and slow waves of $m = -1$ mode, and (b) slow waves of $m = +1$ mode. The plasma parameters are $n_0 = 4 \times 10^{12} \text{ cm}^{-3}$, $B_0 = 0.3 \text{ T}$, $T_e = 0$, $p = 4 \text{ cm}$, and $w = 5 \text{ cm}$.

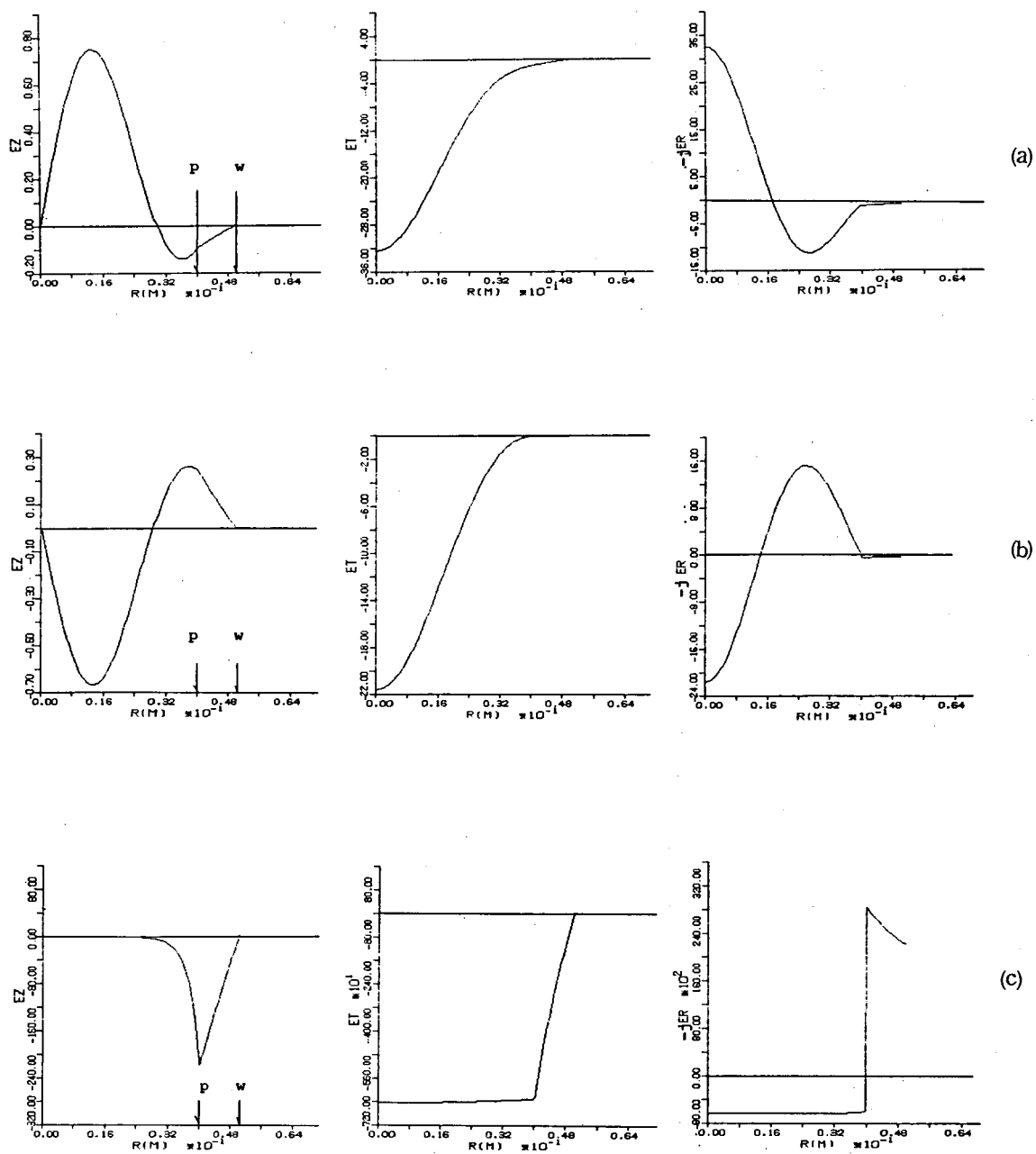


Fig.2-5 (Continued)

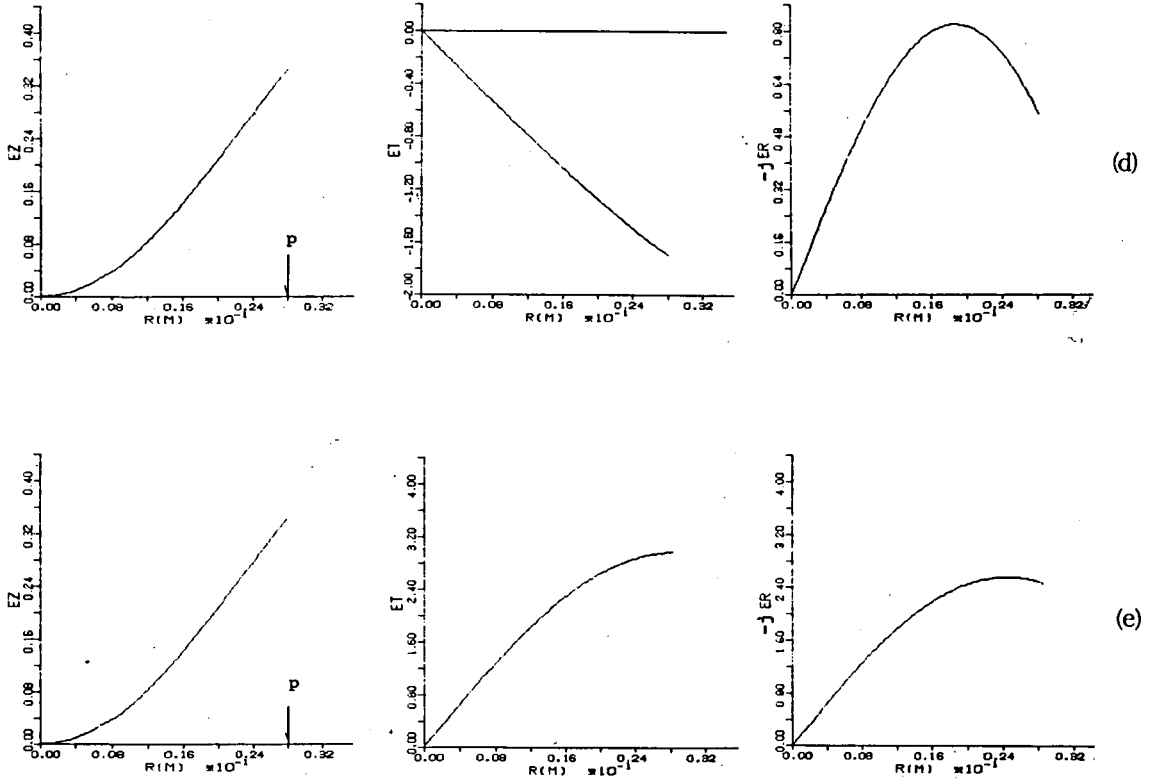


Fig.2-5 Radial distribution of electric fields for (a) $m = +1$ slow wave, (b) $m = -1$ slow wave, (c) $m = -1$ fast wave, (d) $m = +2$ fast wave, and (e) $m = -2$ fast wave. For (a), (b), and (c), $\omega/\omega_{ci} = 0.96$ and $k_{||} = 21 \text{ m}^{-1}$ with other plasma parameters being the same as in Fig.2-4. For (d) and (e), $\omega/\omega_{ci} = 2.4$, $k_{||} = 10.5 \text{ m}^{-1}$, $n_0 = 5 \times 10^{13} \text{ cm}^{-3}$, $B_0 = 0.31 \text{ T}$, $T_e = 20 \text{ eV}$, $p = 3 \text{ cm}$, and $w = 7.5 \text{ cm}$.

$n_0 = 4 \times 10^{12} \text{ cm}^{-3}$, $B_0 = 0.3 \text{ T}$, $T_e = 0$, plasma radius $p = 4 \text{ cm}$, and wall radius $w = 5 \text{ cm}$. The many branches for the slow waves correspond to discrete sets of $k_{\perp\ell}$'s which satisfy the boundary condition. The axial wave number of each branch of the slow waves in (a) or (b) becomes larger as the wave frequency approaches the ion cyclotron frequency. The fast wave appears in very small $k_{\parallel}$ range for $m = -1$ mode, while it becomes cutoff for $m = +1$ mode [1].

In Fig.2-5, the radial profiles of the electric field are shown for (a) $m = +1$ slow wave, (b) $m = -1$ slow wave, (c) $m = -1$ fast wave, (d) $m = +2$ fast wave, and (e) $m = -2$ fast wave. We plot $-jE_r$ instead of E_r . Hence one can easily find that E_r leads E_θ or E_z by 90° if $-jE_r$ is positive, and vice versa. For (a), (b), and (c), the frequency and $k_{\parallel}$ of the wave are chosen to be $\omega/\omega_{ci} = 0.96$ and 21 m^{-1} , respectively, and the plasma parameters are the same as in Fig.2-4. The $m = +1$ slow wave is left-hand circularly polarized on axis and has a maximum amplitude there. Both the $m = -1$ slow wave and $m = -1$ fast wave are right-hand circularly polarized near axis. Although the polarization reverses near the mid-radius for the $m = -1$ slow wave, the $m = -1$ fast wave is almost purely right-hand circularly polarized over all plasma radii.

In (d) and (e), we plot the $m = +2$ and -2 fast wave field with $\omega/\omega_{ci} = 2.4$ and $k_{\parallel} = 10.5 \text{ m}^{-1}$ in a plasma of $n_0 = 5 \times 10^{13} \text{ cm}^{-3}$, $B_0 = 0.31 \text{ T}$, $T_e = 20 \text{ eV}$, $p = 3 \text{ cm}$, and $w = 7.5 \text{ cm}$. The field pattern at a radius larger than 3.0 cm is not shown here. The perpendicular field for these modes is zero on axis and increases with r , in contrast to those of $m = +1$ and -1 modes.

3. EFFECT OF PONDEROMOTIVE FORCE

The equation of motion of particles in an electromagnetic field is given by

$$\frac{\partial \vec{v}_\alpha}{\partial t} + \vec{v}_\alpha \cdot \nabla \vec{v}_\alpha = \frac{q_\alpha}{m_\alpha} (\vec{E} + \vec{v}_\alpha \times \vec{B}), \quad (2-9)$$

where $\alpha = i(\text{ion})$ or $e(\text{electron})$, $\vec{v}_\alpha$ is the particle velocity, $\vec{B}$ is the magnetic field, q_α and m_α are the charge and the mass of the particle, respectively. Expanding physical quantities in power series of field amplitude and assuming $\vec{E}^{(0)} = 0$, $\vec{B}^{(0)} = B_0 \vec{z}$, and $\vec{v}_\alpha^{(0)} = 0$, we obtain the second order part of Eq.(2-9) [11]:

$$\frac{\partial \vec{v}_\alpha^{(2)}}{\partial t} = - \vec{v}_\alpha^{(1)} \cdot \nabla \vec{v}_\alpha^{(1)} + \frac{q_\alpha}{m_\alpha} (\vec{E}^{(2)} + \vec{v}_\alpha^{(1)} \times \vec{B}^{(1)} + \vec{v}_\alpha^{(2)} \times \vec{B}_0). \quad (2-10)$$

Since we are interested in the quasi-steady part, we take a short-time average over a few periods of the fundamental frequency. The slowly varying part of the solution of Eq.(2-10) is written as

$$\vec{v}_{\alpha\perp}^{(2)} = [(\vec{E}^{(2)} + \vec{S}_\alpha) \times \vec{B}_0 + \frac{m_\alpha}{q_\alpha} \frac{\partial \vec{E}_\perp^{(2)}}{\partial t}] / B_0^2, \quad (2-11)$$

$$\frac{\partial v_{\alpha\parallel}^{(2)}}{\partial t} = \frac{q_\alpha}{m_\alpha} (E_\parallel^{(2)} + S_{\alpha\parallel}),$$

where

$$\vec{S}_\alpha = \langle \vec{v}_\alpha^{(1)} \times \vec{B}^{(1)} \rangle - \frac{m_\alpha}{q_\alpha} \langle \vec{v}_\alpha^{(1)} \cdot \nabla \vec{v}_\alpha^{(1)} \rangle, \quad (2-12)$$

which represents so called ponderomotive force if multiplied by q_α . The steady electric field $\vec{E}^{(2)}$ arises due to a charge separation caused by the difference between $\vec{v}_i^{(2)}$ and $\vec{v}_e^{(2)}$.

The axial component of the ponderomotive force has been intensively investigated in researches of RF plugging [6-8]. It should be noted that the perpendicular components of $\vec{S}_\alpha$ take more important roles giving rise to an $\vec{S}_\alpha \times \vec{B}_0$ drift. The radial component of $\vec{S}_\alpha$ makes particles drift azimuthally. We can expect a cancellation or reversal of particle drift due to "bad" curvature of a magnetic field line leading to the stabilization of flute-type instabilities. This effect will be investigated in Chapters 3 and 4.

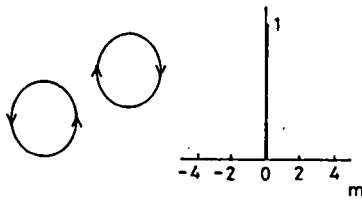
We can also see that a radial drift is caused by the azimuthal ponderomotive force. The radial drift of particles may produce a degradation of the particle confinement. On the other hand, this effect may be useful for the RF pumping in a thermal barrier. In Chapter 5, we will present the effects of the azimuthal ponderomotive force.

4. FORCED EXCITATION OF EIGENMODES

Several types of RF antennae have been used to launch waves in ICRF. Schematic views of the RF antennae are shown in Figs. 2-6 to 2-8 together with the Fourier spectrum a_m of the azimuthal antenna current I_θ , namely,

$$a_m = \frac{1}{2\pi} \int_0^{2\pi} I_\theta(\theta) \exp(-jm\theta) d\theta. \quad (2-13)$$

(a) STIX



(b) HALF TURN

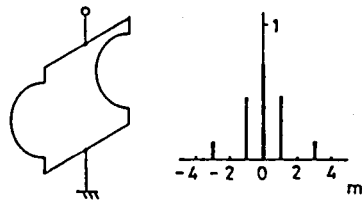
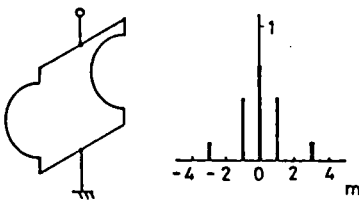
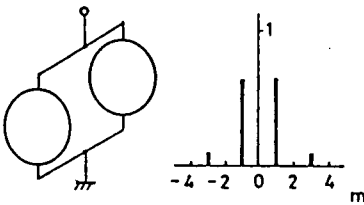


Fig.2-6 RF antennae for generating $m = 0$ RF fields.

(a) HALF TURN



(b) KHARKOV



(c) ROTATING

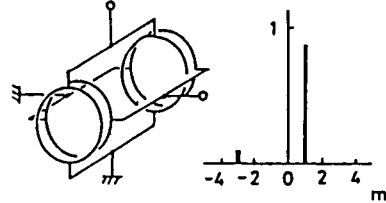
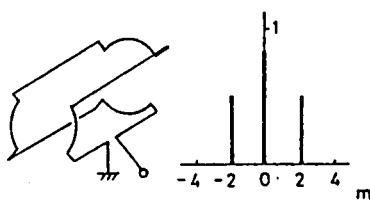


Fig.2-7 RF antennae for generating $m = \pm 1$, $m = +1$, and $m = -1$ RF fields.

(a)



(b) ROTATING

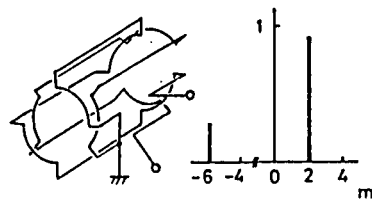


Fig.2-8 RF antennae for generating $m = \pm 2$, $m = +2$, and $m = -2$ RF fields.

The RF power P_w which is coupled to a wave from an antenna is represented by

$$P_w = s \operatorname{Re} < - \int_0^{2\pi} E_\theta|_{r=s} I_\theta^* d\theta > \quad (2-14)$$

where s is the radius of the antenna, and Re and $*$ denote the real part and the complex conjugate, respectively. So, if one wants to excite the wave of azimuthal mode m , one should use the RF antenna whose spectrum has a peak at the m .

Figure 2-6 shows the RF antennae which excite mainly $m = 0$ mode. The Stix coil was used in earlier studies of ICRF heating [2,3], but was replaced by the half turn antenna in experiments in ST tokamak [12] and TFR [13]. For the purpose of ion heating with the slow wave, it is desirable that the wave is left-hand circularly polarized at the plasma core as discussed in Section 2. Referring Fig.2-5, we find that the $m = +1$ mode is the most convenient one. The antennae in Figs.2-7(a) and (b) have $m = +1$ Fourier component, which is, however, not large as compared with other Fourier components. The Kharkov antenna [14] has the $m = +1$ and -1 components of the same amplitude and generates linearly polarized field in vacuum. Therefore, we call this field as the $m = \pm 1$ field. The pair of two half turn antennae used in TFR [13] and the type-III coil used in RFC-XX [15] are essentially the same as the Kharkov antenna in the sense that they all generate the $m = \pm 1$ RF field.

On the contrary, the antenna in Fig.2-7(c) can generate either the $m = +1$ or $m = -1$ component selectively, provided

that the phase of RF currents in two coil units is properly adjusted. Since the fields of these modes are rotating in azimuthal direction, they are termed as the rotating RF field. By using this antenna, it is possible to generate the RF field of purely coherent azimuthal mode.

As will be described in Chapter 3, the cancellation of particles' drift due to a bad curvature of the magnetic field line requires radially inward ponderomotive force at the plasma periphery. The RF field of $m = +2$ or $m = -2$ mode meets this requirement in the case of $\omega > \omega_{ci}$ because $E_{\perp}$ of the field peaks near the periphery and decreases toward the center [Figs.2-5(d) and (e)]. We illustrate the antennae for the excitation of the $m = +2$ or -2 RF field in Fig.2-8. Although the antenna (a) has the multi-mode structure, it will be shown in Chapter 3 that this antenna is sufficient for the excitation of the $m = \pm 2$ mode.

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CHAPTER 3

RF STABILIZATION OF HIGH-DENSITY PLASMA IN AN AXISYMMETRIC MIRROR

1. INTRODUCTION

Recently, a tandem mirror with minimum-B field has received considerable interest as an effective confinement system [1,2]. The plasma in a long central cell is confined by the ambipolar potential created in plug cells located at both ends. It has been, however, pointed out that the resonant radial diffusion due to non-axisymmetric configuration degrades the confinement. In fact, although the original tandem mirror [3-5] had minimum-B end plugs, they were axisymmetrized and only the MHD anchor regions remain asymmetric in present-day and future machines [6-8]. It is, therefore, important to investigate other methods for maintaining MHD stability in order to realize a purely axisymmetric tandem mirror.

Following the experiment performed by Arsenin et al.[9], a great deal of reports concerning the feed back stabilization of flute instabilities were published. This method, which was confirmed to be effective in low-density plasmas, may not be applicable for a fusion-oriented plasma because some electrodes must be placed close to a high-density, hot plasma. This unfavorable situation is the same for the line-

tying stabilization [10] scheme in which conducting plates contact with a plasma. Even when a surface line-tying stabilization [11] is employed, energy loss through the end plates would be significant.

Sidorof [12] discussed the possibility of stabilizing the flute mode by the application of radio frequency field. The experiment of stabilizing a precessional drift wave using this method was done by Yamamoto et al.[13]. In the succeeding paper [14], they showed the flute stabilization by the RF field which was generated electrostatically by the rod electrodes aligned to cover almost whole of the plasma surface. In their experiment, however, the observed instability was not clearly identified and remained to be experimentally clarified. Furthermore, the density was restricted within an order of 10^{12} cm^{-3} , so the availability of this stabilization method was not confirmed for higher densities where the electrostatic RF field excitation would become inefficient.

We here present the experimental identification of the MHD instability in a density range higher than 10^{13} cm^{-3} . The observed instability in the axisymmetric mirror machine HIEI is identified as the curvature-driven flute instability through the experimental determination of the dispersion relation and its comparison with the theoretical prediction. We further demonstrate the stabilization of the flute instability by the RF field which is electromagnetically coupled and applied only in the bad-curvature region. This method of stabilization is revealed to be effective in a wide range

of density, from 10^{10}cm^{-3} up to 10^{14}cm^{-3} . The results of this work suggest that the RF stabilization could be an alternative for MHD anchors of a tandem mirror.

2. DESCRIPTION OF THE EXPERIMENT

The experiment was carried out in the axisymmetric mirror HIEI as shown in Figs.3-1 and 3-2. The hydrogen plasma of density ranging from 10^{10} to $5 \times 10^{11}\text{cm}^{-3}$ (low density regime) or from 10^{13} to 10^{14}cm^{-3} (high density regime) is produced by TPD or MPD plasma guns [16], respectively, in a static magnetic field up to 1.1 T at the throat. The length between two mirror points is 1.2 m and the mirror ratio is variable from 1.4 to 4. The plasma diameter at the throat is 4 cm (FWHM), which is fixed by the limiter made of pyrex glass. Typical electron temperatures are $T_e = (15-20)$ eV, and ion temperatures without heating are $T_i = (10-15)$ eV. The gun injected plasma is terminated by an end plate of insulator located outside the mirror point.

The RF stabilization of MHD instabilities is accomplished by applying the RF field only in the bad-curvature region. The stabilizing antenna shown in Fig.3-3 consists of a copper strap of 2 cm width and 0.3 cm thickness formed just like a quadrupole winding of 7 cm in diameter and 30 cm in length. This antenna is placed at the midplane and produces the RF field of $m = \pm 2$ mode in frequency range $1.5 < f/f_{ci0} < 3$, where f is the applied RF frequency and f_{ci0} is the local ion

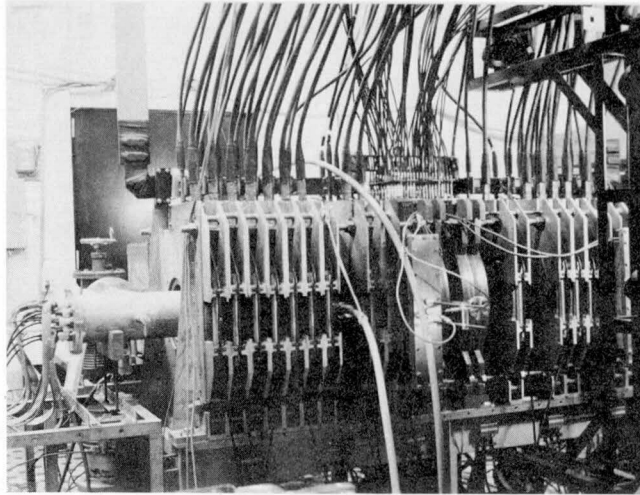


Fig.3-1 The photograph of HIEI.

HIEI

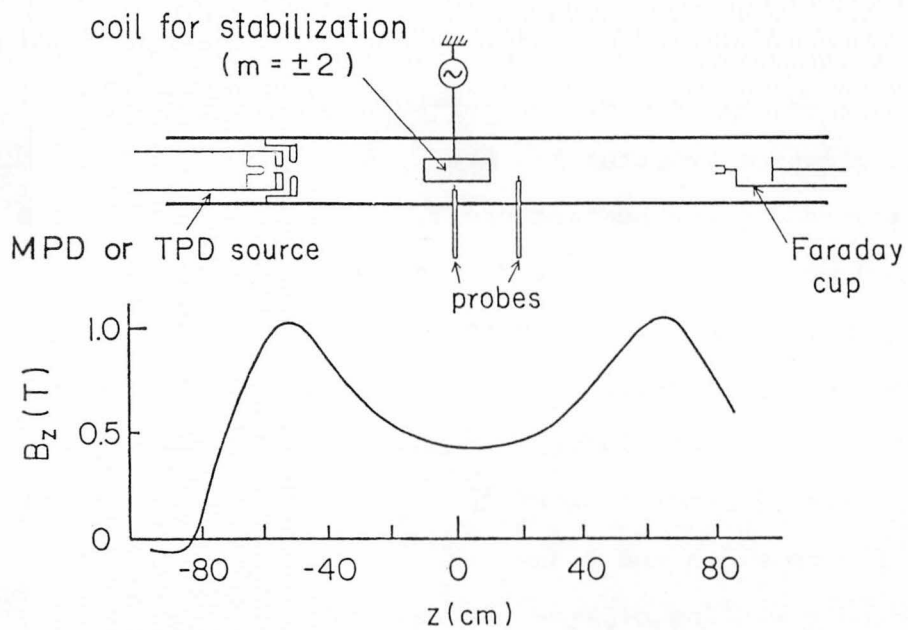


Fig.3-2 Experimental set up.

cyclotron frequency at the midplane. This antenna is driven by an RF self-oscillator which can deliver pulsed RF power up to 200 kW at $f = 10$ MHz.

The maximum RF current in the antenna was $50 A_{\text{rms}}$ in this experiment. The modulation of the static magnetic field due to the RF is less than 20 G at the plasma surface. This is an order of magnitude smaller than the difference between the confining magnetic field strengths at the plasma surface and on axis for typical operations. Hence, the modification of the mirror field by the RF magnetic field is negligible.

Plasma parameters and instability signals are measured with a charge exchange analyzer, diamagnetic loops, end loss analyzers, and electrostatic probes. The raw data from these equipments are A/D-converted, stored on magnetic tapes, and analyzed in real time mode with a digital data acquisition system using multiple microprocessors. The block diagram of this system is shown in Fig.3-4(a), where three blocks are connected to a main bus line. Each block including one CPU has a peculiar function such as controls of A/D converters, display and management of experimental parameters, controls of a graphic display, and so on. These functions are operated by a request from the bus master (CPU 1). An example of the data acquisition is shown in Fig.3-4(b), where raw data of voltage and current from a swept double probe (graph 1A and 1B) are converted into plot of V-I characteristics. We have also developed an FFT program, which is used for an analysis of instabilities described in next Section.

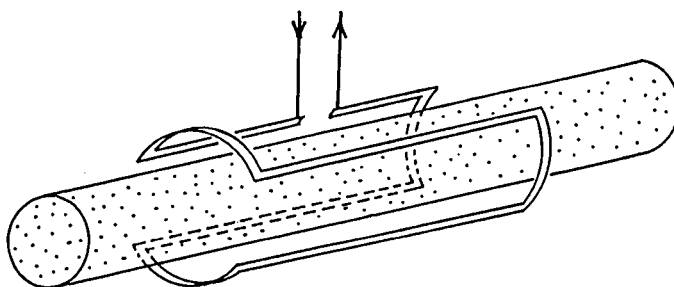
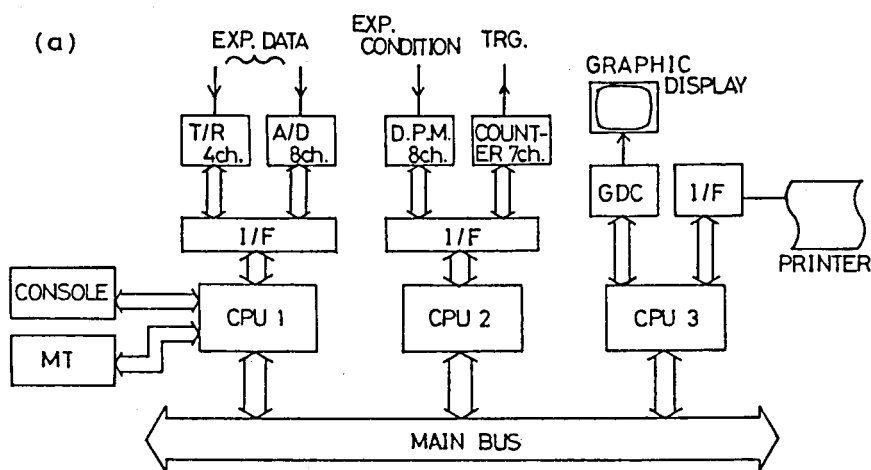


Fig.3-3
RF antenna for
stabilization.



DATA ACQUISITION SYSTEM

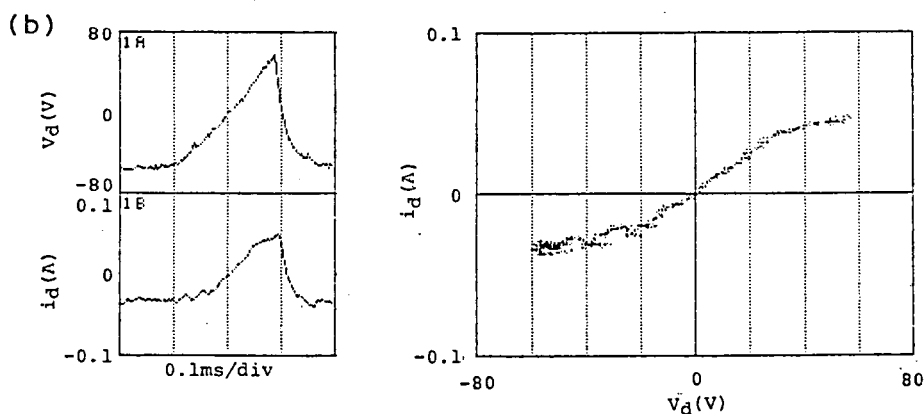


Fig.3-4 (a) Blockdiagram of data acquisition system, which consists of transient recorders(T/R), A/D converters(A/D), digital pannel meters(DPM), graphic display controller(GDC), and so on. (b) Raw data from a double probe (1A and 1B) and V_d-i_d plot of them.

3. IDENTIFICATION OF THE INSTABILITY

A. Measurement of the Mode

Fairly monochromatic oscillations of plasma density and potential are observed near the plasma periphery in both high- and low-density regimes, with a fluctuation level $\tilde{n}/n_0$ larger than 80%. In the following, we will determine the mode of the instability in both density regimes separately.

(1) Low density regime

Figure 3-5 shows typical oscillogram of the density fluctuation near plasma periphery ($r/r_L = 0.6$ with r_L being the local radius of the plasma which maps on the limiter radius at the throat.) and the ion saturation current at $r = 0$. As the instability grows up, the build-up of the central density slows down and then the density begins to decrease until the amplitude of the instability gets saturated. The density at quasi-steady state was about 60% of that without the instability. The decay time of the density after the gun turn-off was rapid as compared with the stable case.

The radial variation of the fluctuation level, plasma density, space potential, and T_i are plotted in Fig.3-6 for magnetic field strength at the throat $B_M = 0.38$ T, mirror ratio $R = 2.4$, and density $n_0 = 1 \times 10^{11} \text{ cm}^{-3}$. For this mirror ratio, the radius of curvature of the magnetic field line, R_c , is 4.3 m. The electron temperature on axis was 20 eV. The fluctuation level has a maximum around $r/r_L = 0.7$ at the midplane, where the density gradient is large, and has

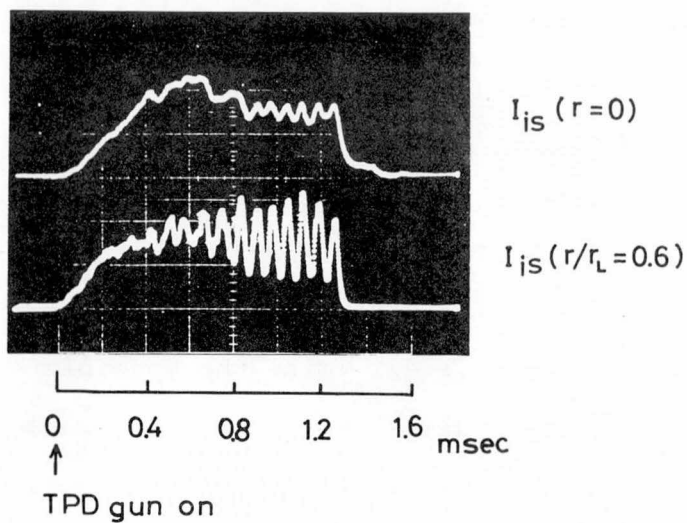


Fig.3-5 Typical oscillogram of fluctuation for low-density regime.

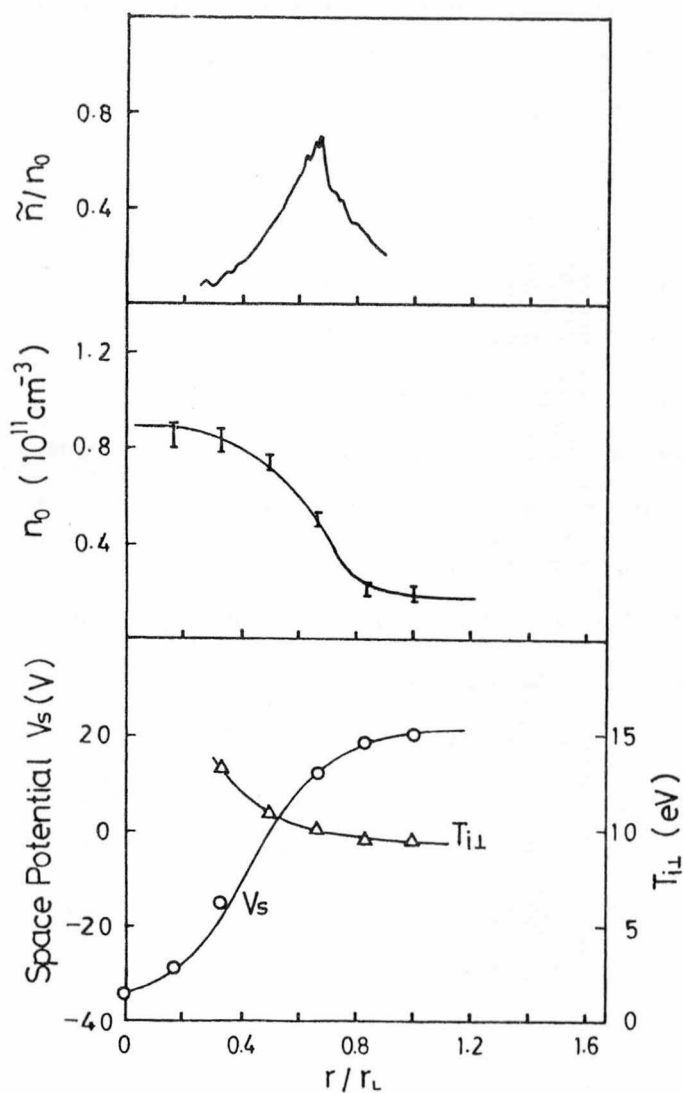
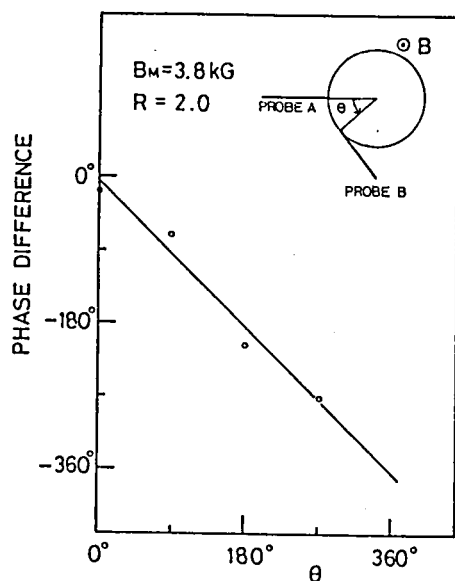
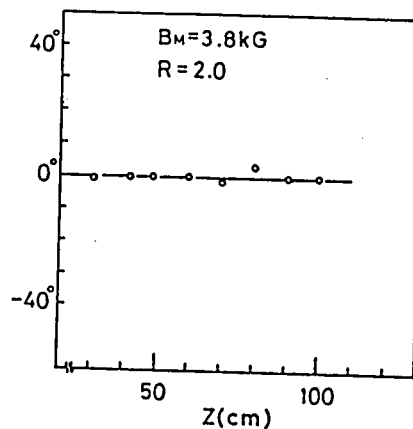


Fig.3-6 Radial variation of plasma parameters for low-density regime. $n_0 = 1.0 \times 10^{11} \text{ cm}^{-3}$, $B_M = 0.38 \text{ T}$, and $R = 2.4$.



(a)



(b)

Fig.3-7 Phase difference of the instability as a function of (a) azimuthal probe angle and (b) axial length.

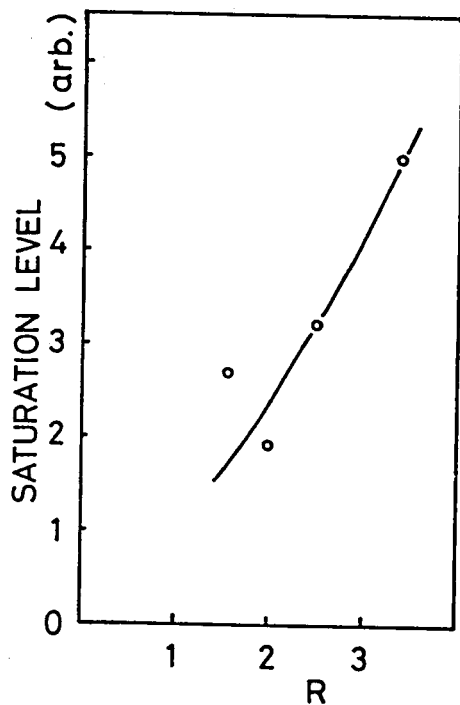


Fig.3-8 Fluctuation level versus the mirror ratio.

very small amplitude at the center. The frequency spectrum of the fluctuation was sharply peaked at about 15 kHz. The wave number of the fluctuation was measured by two single probes using cross-correlation technique. Setting one probe at a radius where the fluctuation level is a maximum, we move the other probe along the circumference with the same radius. The measured phase difference is plotted in Fig.3-7(a) as a function of azimuthal offset angle of two probes. We can see that the observed fluctuation is of $M = 1$ azimuthal mode. The direction of propagation coincides with that of electron diamagnetic drift. We also measured the phase difference along the magnetic field line to find that the axial wave number is nearly zero within the plasma column as shown in Fig.3-7(b).

As will be shown later, the fluctuation is stabilized by applying an RF field. We can measure the growth rate of the instability by measuring the time evolution of the fluctuation amplitude after switching off the RF pulse. The value of the growth rate Ω_i thus obtained is $0.2 < \Omega_i/\Omega < 0.4$ with Ω being the real angular frequency. The fluctuation level observed at the midplane increased with the mirror ratio R as shown in Fig.3-8, roughly indicating the increase of the growth rate with R .

(2) High-density regime

Figures 3-9 and 3-10 show, respectively, typical density fluctuation and radial variations of the fluctuation level, the density, and the space potential for $B_M = 0.51$ T, $R = 2.3$,

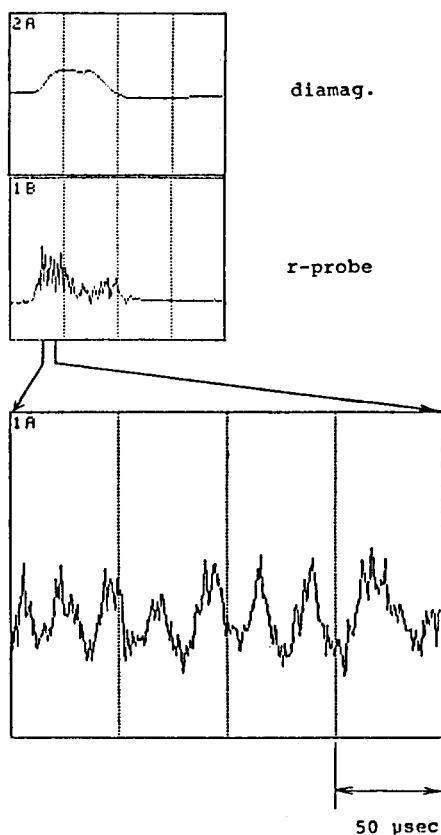


Fig.3-9 Typical oscillogram of the fluctuation in high-density regime.

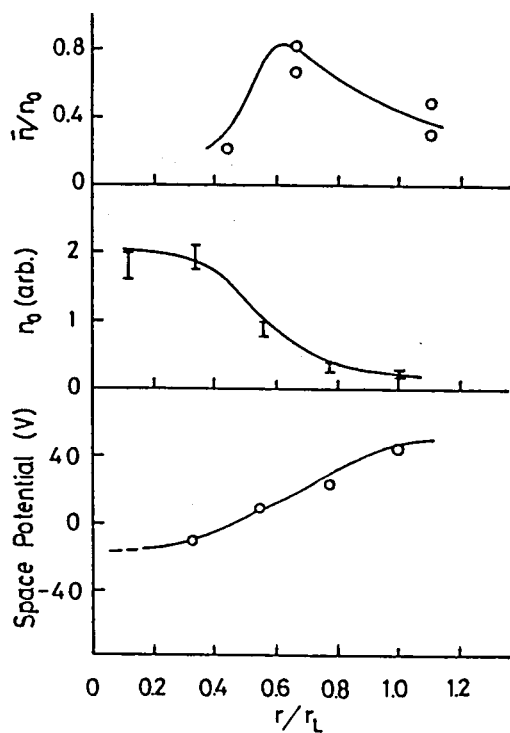


Fig.3-10 Radial distribution of plasma parameters in high-density regime. $n_0 = 2 \times 10^{13} \text{ cm}^{-3}$, $B_M = 0.51 \text{ T}$, $R = 2.3$, and $T_e = 11 \text{ eV}$.

$T_e = 11$ eV, and $n_0 = 2 \times 10^{13} \text{ cm}^{-3}$. The fluctuation level has a maximum around $r/r_L \approx 0.6$ [$r = (2.5-3)$ cm] at the midplane, and has very small amplitude on axis. The frequency spectrum of the fluctuation was sharply peaked at about 30 kHz and had a minor peak around 15 kHz. The result of the cross-correlation measurement is shown in Fig.3-11; (a) and (b) are auto power spectra of two probe signals, (c) and (d) are cross amplitude spectrum and phase spectrum, respectively. The phase difference of two probe signals at a frequency of 30 kHz, which is given by Fig.3-11(d), is plotted in Fig.3-11(e) as a function of the azimuthal offset angle of the probes. Three lines in the figure represent ideal cases of $M = 1$, 2, and 3 azimuthal modes. It is seen that the observed fluctuation is of $M = 2$ mode. We also measured the phase difference along the magnetic field line to find again that the axial wave number is zero within the plasma column. The fluctuation level versus R is given in Fig.3-12 for $B_M = 0.42$ T and $n_0 = 2 \times 10^{13} \text{ cm}^{-3}$. In this case the fluctuation level was measured at a radius of maximum instability amplitude for each mirror ratio. It should be noted that the fluctuation level becomes greater than several tens of % immediately after R exceeds unity. The normalized growth rate was $0.4 < \Omega_i/\Omega < 0.6$.

We finally summarize the results of the observation in low- and high-density regimes in Table 3-1.

B. Comparison with Theory

There have been many theoretical works concerning the

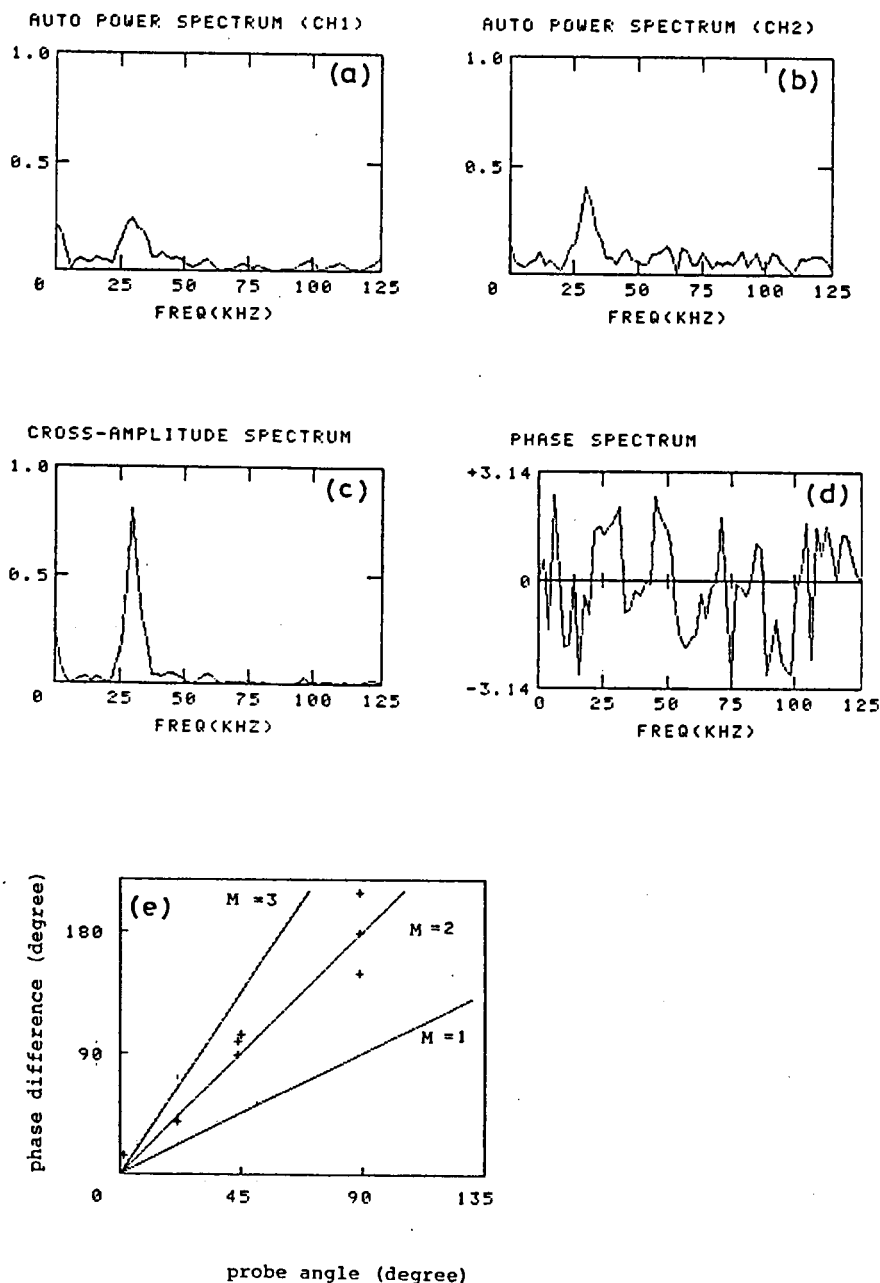


Fig.3-11 (a) Auto-power spectrum of one probe signal.
 (b) Auto-power spectrum of the other probe signal.
 (c) and (d) are cross-amplitude spectrum and phase spectrum of the two probe signals, respectively.
 (e) Phase difference obtained from (d) as a function of the offset angle of the two probes.

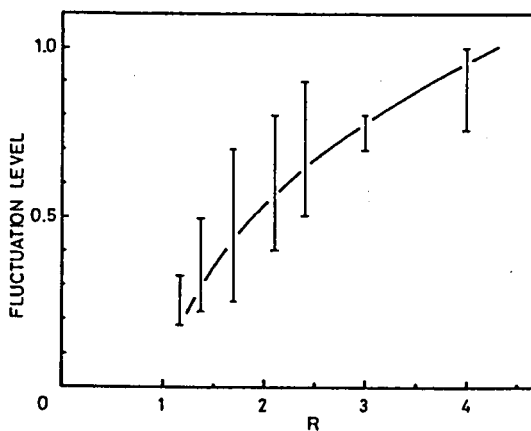


Fig.3-12 Fluc-
tuation level
versus the mirror
ratio for $n_0 =$
 $2 \times 10^{13} \text{ cm}^{-3}$ and
 $B_M = 0.42 \text{ T}$.

Table 3-1

regime	mode M	region r/r_L	frequency $\Omega/2\pi(\text{kHz})$	growth rate Ω_i/Ω
low density	1	0.6-0.7	12-20	0.2-0.4
high density	2	0.6-0.7	25-35	0.4-0.6

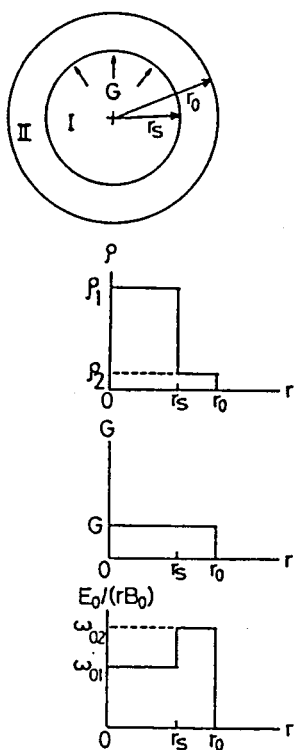


Fig.3-13 Theoretical model
of the plasma cross section.
 ρ is the normalized density,
 G is the effective gravity,
and ω_0 is the rotation
frequency.

dispersion relation of flute mode. In the present problem, the perpendicular wave length of the observed instability is of the order of plasma radius. Hence it is not appropriate to use a local approximation in deriving the dispersion relation. We therefore employ the finite Larmor radius equation derived by Rosenbluth and Simon[17]. This equation describes the low frequency flute perturbation of the form $\exp[j(M\theta + \Omega t)]$ propagating in a cylindrical plasma with arbitrary density, radial electric field, and gravity profiles, and is given by

$$\frac{1}{r} \frac{d}{dr} (T \frac{d\psi}{dr}) + \left[\frac{(1 - M^2)}{r^3} T + (M^2 G + \Omega^2 r) m_i \frac{dn_0}{dr} \right] \psi = 0, \quad (3-1)$$

where $T = \bar{\Omega}^2 m_i n_0 S r^3$, $\psi = E_\theta / \bar{\Omega}$, $\bar{\Omega} = \Omega - M\omega_0$, and $S = 1 - M / (r m_i n_0 \omega_{ci0} \bar{\Omega}) d(n_0 e T_i) / dr$. Here, E_θ is the azimuthal component of electric field of the flute fluctuation, $\omega_0 [= E_0(r) / r B_0]$ is the rotation frequency of the plasma drift due to the steady radial electric field E_0 and the external magnetic field B_0 (evaluated at the midplane), m_i is the ion mass, and T_i is in the unit of eV. For simplicity, we assume the radial variation of the plasma parameters such that n_0 and ω_0 have a discontinuity at $r = r_s$, but are constant in both sides $0 < r < r_s$ (region 1) and $r_s < r < r_0$ (region 2) [see Fig.3-13]. In region i, the normalized density $\rho_i = n_{0i} / (n_{01} + n_{02})$ and $\omega_{0i} = E_{0i}(r) / (r B_0)$ ($i = 1, 2$), and there is a perfectly conducting wall at $r = r_0$. The effective gravity G due to the curvature of the magnetic field line is given by $e T_i / R_c$, which is assumed to be constant in both regions. Following the procedure used by Kent et al.[18], we obtain the dispersion

relation from Eq.(3-1):

$$\begin{aligned}
& [\rho_1(\zeta-1)+\rho_2(\zeta+1)]\Omega^2 - 2[(M-1)\{\rho_1(\zeta-1)+\rho_2(\zeta+1)\} \\
& + 2\rho_2\zeta]\omega_{0s}\Omega + M\omega_{0s}^2[(M-1)\{\rho_1(\zeta-1)+\rho_2(\zeta+1)\} \\
& + 2\rho_2\zeta] - 2M(M^2-1)\omega_{0s}\gamma(\rho_1-\rho_2)(\zeta-1) \\
& + \frac{M}{r_{sm_i}} \frac{G}{\rho_1-\rho_2}(\zeta-1) = 0,
\end{aligned} \tag{3-2}$$

where $\zeta = (r_0/r_s)^{2M}$, $\omega_{0s} = (\omega_{01} + \omega_{02})/2$, and $\gamma = eT_i/(2r_s^2 eB_0)$. As we have replaced M in Eq.(3-1) by $-M$, Eq.(3-2) describes the perturbation of the form $\exp[j(M\theta - \Omega t)]$. From Fig.3-10, we can determine the parameters in Eq.(3-2) to be $\rho_1 = 0.77$, $2 < r_s(\text{cm}) < 3$, $4 < E_{01}(\text{V/cm}) < 5$, $20 < E_{02} < 25$ for the high-density regime. We have chosen the value of ρ_1 so that $(\rho_1 - \rho_2)/r_s$, which corresponds to the density gradient, is equal to the measured density gradient. The other parameters necessary for calculation are $B_0 = 0.22 \text{ T}$, $R_c = 4.3 \text{ m}$, and $T_i = 20 \text{ eV}$. As the measured values, especially, r_s and E_{0i} have some ambiguity, we calculate Eq.(3-2) for several combinations of r_s and E_{0i} and plot the frequency $\Omega/2\pi$ and the square of the growth rate Ω_i^2 in Fig.3-14 as a function of M . Although the values of $\Omega/2\pi$ and Ω_i have some spread, all curves predict that $M = 2$ mode is most unstable in a frequency range from 20 to 40 kHz.

The circles in Fig.3-14 denote experimental points for the high-density regime (see Table 3-1). The agreement between the theory and the experiment is very good in spite of the simplification of the theoretical model. It is also predicted

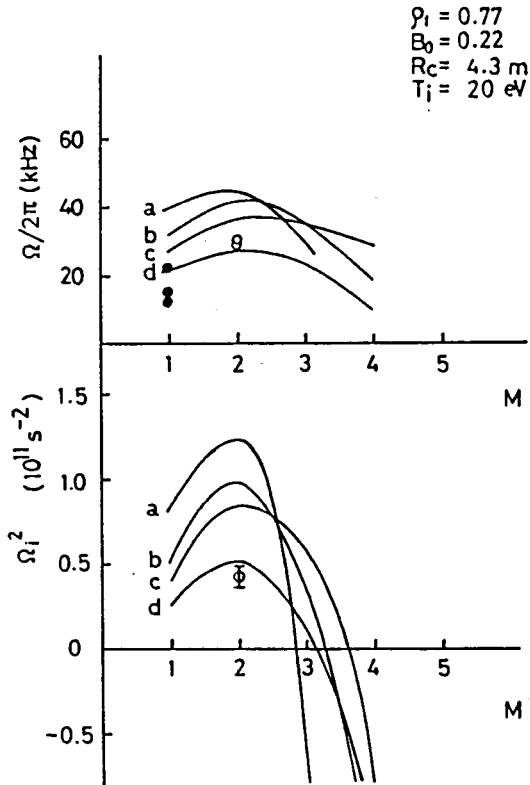


Fig.3-14 Dispersion relation of the instability. Solid curves are the solution of Eq.(3-2): curve a; $E_{01} = 5 \text{ V/cm}$, $E_{02} = 25 \text{ V/cm}$, $r_s = 2.0 \text{ cm}$, curve b; $E_{01} = 5 \text{ V/cm}$, $E_{02} = 25 \text{ V/cm}$, $r_s = 2.5 \text{ cm}$, curve c; $E_{01} = 5 \text{ V/cm}$, $E_{02} = 25 \text{ V/cm}$, $r_s = 3.0 \text{ cm}$, and curve d; $E_{01} = 4 \text{ V/cm}$, $E_{02} = 20 \text{ V/cm}$, $r_s = 3.0 \text{ cm}$. Open circles and closed circles are the experimental points for $M = 2$ and $M = 1$ modes, respectively.

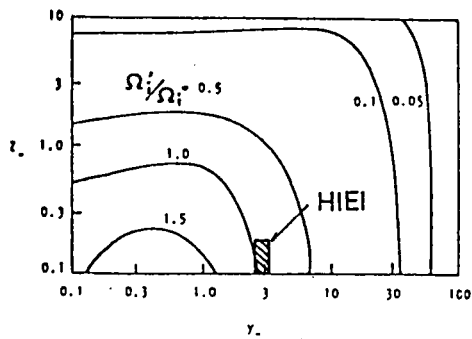


Fig.3-15 Modified growth rate of the flute instability due to the line-tying effect on y - z plane. (Prater, 1974)

from Eq.(3-2) that the growth rate increases with decreasing R_c , in accordance with the experimental result in Fig.3-12. Thus, we conclude that the observed fluctuation in the high-density regime is the flute instability.

For the low-density regime, we also calculate Eq.(3-2) to find that $10 < \Omega/2\pi$ (kHz) < 20 and $0.4 < \Omega_i/\Omega < 0.6$ for $M = 1$ mode, and that $\Omega_i^2 < 0$ for $M \geq 2$. These values are in quantitative agreement with the measured ones shown in Table 3-1. The difference of the most unstable mode number from that in the high-density regime is mainly due to larger $|E_0|$ and larger $d|E_0|/dr$ in the TPD plasma. We can say that the instability arising in the low-density regime is also identified as the flute instability.

C. Effect of Line-Tying

It has been said that the flute instability hardly occurs in moderate-scale experiments because a mirror plasma contacts directly, or through an ambient plasma, with end walls. In our experiment, however, it seemed that the end effects were negligible on flute destabilization. As shown in Fig.3-12, the fluctuation level of the flute mode exceeds 0.1-0.2 (typical for drift mode) and increases monotonously immediately after the mirror ratio exceeds unity. If there exists some stabilizing effect, the fluctuation level should remain small and constant until the curvature effect overcomes the stabilizing effect. So, in the following, we consider the line-tying effect in our experiment in more detail.

The MPD gun used in the experiment consists of a rod

cathode and a hollow ring anode separated each other by a Macor insulator. The magnetic flux tube on which the instability has a maximum amplitude at the midplane maps on the insulator surface, not on metal electrodes. The other end of the plasma is terminated on the end target made of pyrex glass. Thus there is no contact of the plasma with conducting walls at least in a radial region where the instability amplitude is large.

Prater [19] reported that the longitudinal current flow to conducting end plates can decrease or even increase the growth rate of the flute instability depending on the admittance of the sheath at the contact region. He showed that if the sheath parameters, y_- and z_- , are small, the growth rate of the flute instability is enhanced and vice versa. Here, $y_- = \omega_{ci}^2 r_s^2 / (2c_s^2 L\Omega)$ and $z_- = v_{Ti} / (2L\Omega)$ with c_s , L , and v_{Ti} being, respectively, the sound velocity, the length of the ambient plasma, and the ion thermal velocity. Figure 3-15 shows the value of Ω_i' / Ω_i on $y_- - z_-$ plane where Ω_i' and Ω_i are the growth rate of the flute mode with and without the wall effect. Calculation of Eq.(14) in Ref.19 under the present experimental condition gives the sheath parameters shown as a hatched region in Fig.3-15. Since $0.8 < \Omega_i' / \Omega_i < 1.1$, it is seen that, according to the Prater's theory, the stabilizing effect is very weak, if any. This prediction is in agreement with our observation. Thus, the second reason is the relatively low sheath admittance which prevents electrons from moving freely to cancel the charge separation due to flute perturbation. We conclude that the effect of

line-tying on the flute instability is negligible in our experimental condition for these two reasons.

4. RF STABILIZATION

A. Stabilization Measurements

The flute instability described in the previous Section is investigated in the presence of RF fields. During a quasi-steady state of the plasma, which lasts about 1.5 msec, the RF field of $m = \pm 2$ mode is applied with the RF antenna shown in Fig.3-3 for a period of about 0.5 msec. The density fluctuation of the flute instability is sampled for this period except for each 50 μ sec of beginning and end of the RF pulse. Figure 3-16(a) shows the flute fluctuation for $B_M = 0.39$ T, $R = 1.8$, and $n_0 = 5 \times 10^{10} \text{ cm}^{-3}$ (low-density regime). If we apply the RF field ($\omega/\omega_{ci0} = 2.9$ at the midplane) for a duration of the gate pulse shown in the figure, the amplitude of the flute fluctuation decreases with increasing the RF antenna current I_{RF} [2.5 A_{rms} in (b) and 3.5 A_{rms} in (c)]. The power spectra of the fluctuation during the RF pulse is displayed in Fig.3-16(d) for $I_{RF} = 0, 1.1, 2.5$, and 3.5 A_{rms} . For $I_{RF} = 3.5 A_{rms}$, the spectral component of the flute fluctuation around 20 kHz completely disappeared. Figure 3-17 also shows the power spectra of the flute instability in a plasma with $B_M = 0.62$ T, $R = 2$, and $n_0 = (3-5) \times 10^{13} \text{ cm}^{-3}$ (high-density regime). The arrow in each spectrum denotes the D.C. level of the ion saturation current. When

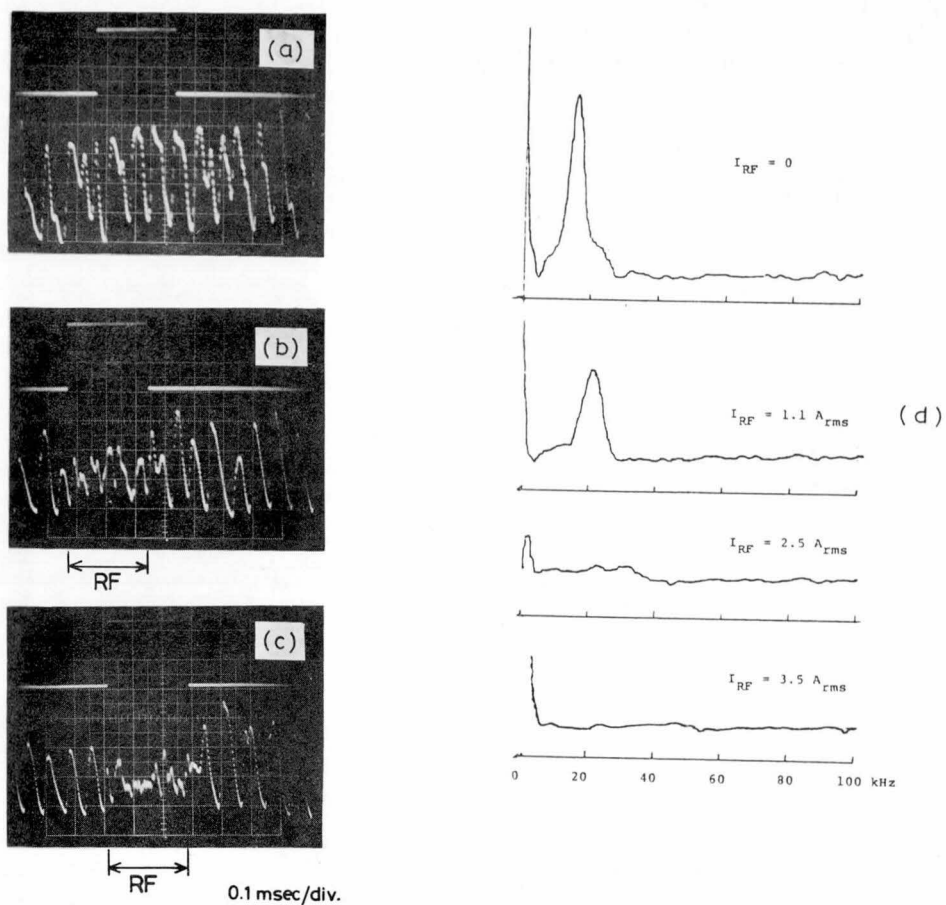


Fig.3-16 RF stabilization of the flute mode in the low-density regime. (a), (b), and (c) are the fluctuation signals at the plasma periphery. RF is applied during the gate pulse with (a) $I_{RF} = 0$, (b) $I_{RF} = 2.5 A_{rms}$, and (c) $I_{RF} = 3.5 A_{rms}$. (d) Power spectra of the flute mode for several I_{RF} .

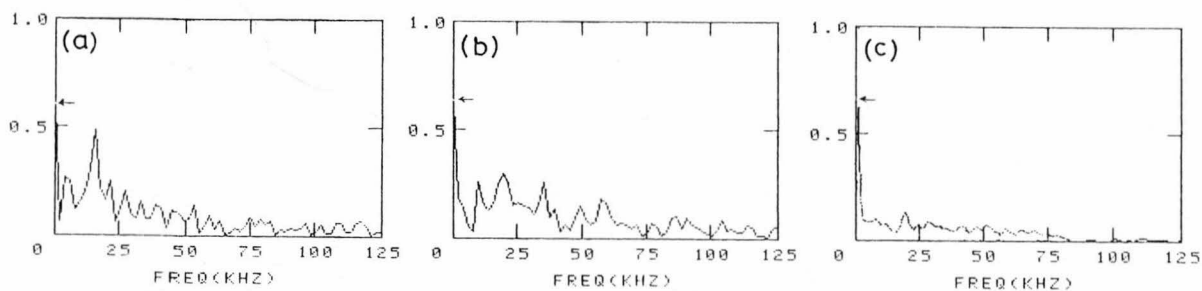


Fig.3-17 RF stabilization of the flute mode in the high-density regime. The power spectra of the flute mode are shown for (a) $I_{RF} = 0$, (b) $I_{RF} = 20 A_{rms}$, (c) $I_{RF} = 25 A_{rms}$.

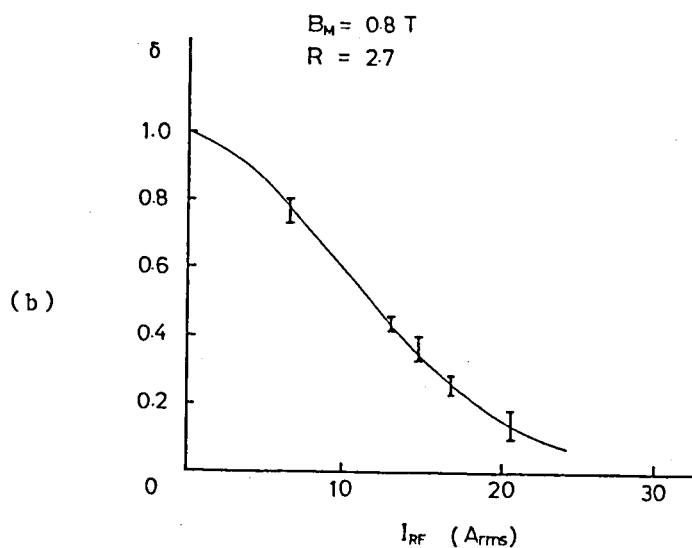
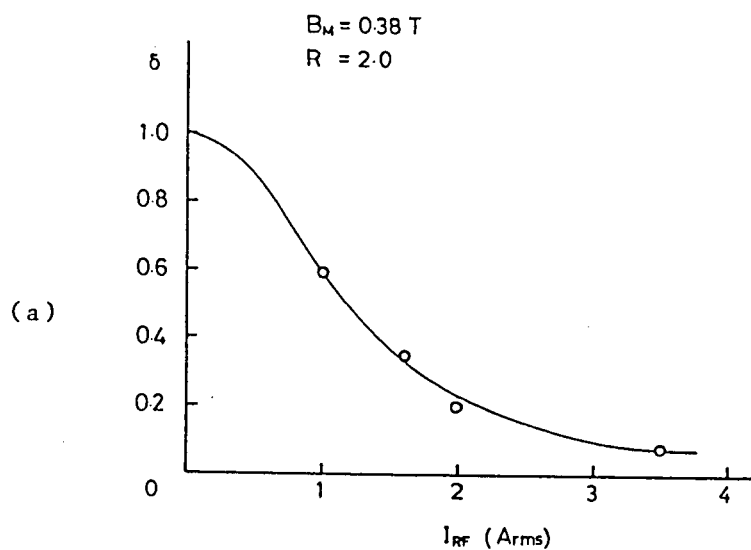


Fig.3-18 Suppression ratio versus I_{RF} for (a) $n_0 = 6 \times 10^{10} \text{ cm}^{-3}$ and (b) $n_0 = 1 \times 10^{13} \text{ cm}^{-3}$. The value of ω/ω_{ci} is 3.4 for (a) and 2.2 for (b).

the $m = \pm 2$ RF field ($\omega/\omega_{ci0} = 2.4$) is applied, the spectral amplitude at 20 kHz decreases [(b)], and further increase of I_{RF} results in almost complete suppression of the flute instability [(c)]. As the vertical gain is doubled in the spectrum (c) so that the D.C. level is approximately equal to those in (a) and (b), one can readily compare the fluctuation level each other. In practice, the ion saturation current for (c) at the plasma periphery decreased to half of the initial value showing that the radial diffusion is reduced through the flute stabilization. The flute instability disappeared everywhere in the plasma, even though the RF stabilization was applied only for 30 cm length around the midplane.

When I_{RF} was increased, it was observed a little upward shift of the instability frequency [see Fig.3-16(d)], which was hardly notable in the high-density regime. This may be attributed to a change in the plasma density and/or its profile caused by the RF field.

In Figs.3-18(a) and (b), the ratio of the spectral amplitude of the instability with stabilizing RF to that without the RF, δ , is plotted versus I_{RF} for the low- and high-density regimes, respectively. In the low-density regime, several amps of I_{RF} was sufficient for the stabilization with $\delta < 0.2$, while much larger current was required for the high-density regime.

B. Estimation of the Ponderomotive Force

When the large amplitude RF field is applied to the plasma,

charged particles feel the ponderomotive force as described in Section 3 of Chap.2. If the radial component of the ponderomotive force is such that the resultant azimuthal drift of ions or electrons tends to cancel the charge separation of the flute instability, the charge accumulation is reduced and the stabilization of the flute instability will be achieved.

The RF electric field in the plasma may be represented by

$$\begin{aligned} E_r^{(1)} &= E_r(r) \cos \phi \\ E_\theta^{(1)} &= E_\theta(r) \sin \phi \\ E_z^{(1)} &= E_z(r) \sin \phi \end{aligned} \quad (3-3)$$

where $\phi = m\theta - k_z z + \omega t$. The radial component of the ponderomotive force [see Eq.(2-12)] is written as

$$\begin{aligned} F_{\alpha r} &= q_\alpha S_{\alpha r} \\ &= q_\alpha \langle v_{\alpha\theta}^{(1)} B_z^{(1)} - v_{\alpha z}^{(1)} B_\theta^{(1)} \rangle - m_\alpha \langle v_{\alpha r}^{(1)} \frac{\partial v_{\alpha r}^{(1)}}{\partial r} \\ &\quad + v_{\alpha\theta}^{(1)} \frac{1}{r} \frac{\partial v_{\alpha r}^{(1)}}{\partial \theta} + v_{\alpha z}^{(1)} \frac{\partial v_{\alpha r}^{(1)}}{\partial z} - \frac{1}{r} v_{\alpha\theta}^{(1)2} \rangle, \end{aligned} \quad (3-4)$$

where angular brackets denote high-frequency average. The values of $\vec{B}^{(1)}$ and $\vec{v}^{(1)}$ can be obtained using Eq.(3-3), Maxwell's equations, and the equation of motion. Substituting these values into Eq.(3-4), we obtain

$$\begin{aligned}
F_{ir} = & \frac{e^2}{2m_i} \frac{1}{\omega^2 - \omega_{ci}^2} \left[\frac{1}{\omega} (\omega_{ci} E_r - \omega E_\theta) (E_\theta' + \frac{E_\theta}{r} + \frac{m}{r} E_r) \right. \\
& - \frac{1}{\omega^2 - \omega_{ci}^2} \{ (\omega E_r - \omega_{ci} E_\theta) (\omega E_r' - \omega_{ci} E_\theta') + \frac{1}{r} (\omega_{ci} E_r - \omega E_\theta)^2 \\
& \left. - \frac{m}{r} (\omega_{ci} E_r - \omega E_\theta) (\omega E_r - \omega_{ci} E_\theta) \} \right] \quad \text{for ions} \quad (3-5)
\end{aligned}$$

$$F_{er} = \frac{e^2}{2m_e} \frac{1}{\omega^2} E_z E_z' \quad \text{for electrons} \quad (3-6)$$

where the prime denotes the derivative with respect to r . In above equations, we have retained only major terms. It should be noted that Eq.(3-5) differs from the conventional expression of the radial ponderomotive force

$$F_{ir} = \frac{-e^2}{4m_i} \frac{1}{\omega^2 - \omega_{ci}^2} \nabla_r E^2, \quad (3-7)$$

which has been employed for the explanation of the RF stabilization [13,14]. In the present experimental condition, the azimuthal scale length of the RF field is comparable to the radial scale length. Furthermore, the polarization of the RF field varies with r . Since these effects are not included in Eq.(3-7), we must use Eq.(3-5) for the present problem.

In order to calculate Eq.(3-5) and (3-6), we must determine the RF field distribution first. Figures 2-5(d) and (e) in Chap.2 show the calculated RF electric field distribution for the $m = +2$ and -2 modes, respectively, in a

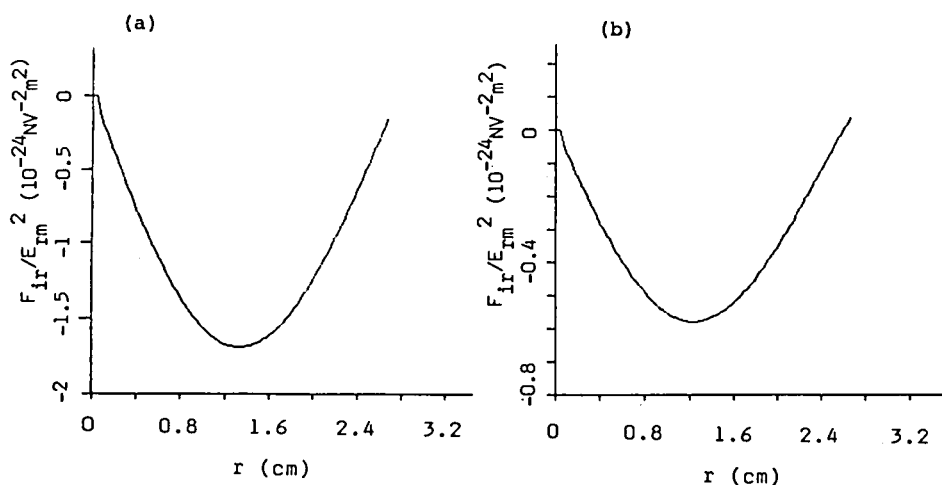


Fig.3-19 Calculated value of the radial ponderomotive force for ions, F_{ir} , as a function of the plasma radius for (a) $m = -2$ mode and (b) $m = +2$ mode. The assumed electric fields are shown in Fig.2-5 (d) and (e) in Chap.2.

cylindrical, sharp boundary, cold plasma. The plasma parameters are the same as in Fig.3-17 except for the fact that the density profile in the calculation is rectangular rather than parabolic as in the experiment. In general, $E_\theta(r)$ increases with r , while $E_r(r)$ becomes a maximum at a certain radius within a plasma radius. We define the maximum value of $|E_r(r)|$ as E_{rm} . Substitution of $\vec{E}(r)$ in Fig.2-5 into Eq.(3-5) gives the radial ponderomotive force for ions. The value of F_{ir} normalized by E_{rm}^2 is shown in Fig.3-19 as a function of r for (a) $m = -2$ and (b) $m = +2$ modes. For both modes, F_{ir} is negative in almost the entire plasma column producing the ion drift in $+\theta$ direction, which cancels the charge separation resulting from the curvature drift. The RF field generated by the antenna shown in Fig.3-3 is of $m = \pm 2$ mode. Since $\vec{B}^{(1)}$ of $m = \pm 2$ mode is a linear combination of those of $m = +2$ and -2 modes, the value of F_{ir} for $m = \pm 2$ mode will be in the range between the above two values. [Note that the cross-term such as $v_{i\theta} \cos(2\theta - k_z z + \omega t) \times B_z \cos(-2\theta - k_z z + \omega t)$ vanishes if averaged over θ and t .]

We can calculate F_{er} in Eq.(3-6) in the same manner. The ponderomotive force for electrons is produced mainly by the product of B_θ and v_{ez} . But, in the present experiment, electron-neutral and electron-ion collision frequencies are much larger than ω , hence the phase correlation between B_θ and v_{ez} is easily destroyed by the collisions. Therefore, we can expect no ponderomotive force for electrons. For a case of weak collisions, F_{er} would become greater than F_{ir} .

The necessary condition for the RF stabilization of the

flute instability may be written as

$$- F_{ir} > G . \quad (3-8)$$

The value of the RF magnetic field is measured to be $|B_\theta|(r/r_L = 0.8) \sim 6 G_{pp}$ (Measured B_θ is maximum at this position.) under the same condition as in Fig.3-17(c). Since the calculated radial distribution of $\vec{E}$ and $\vec{B}$ determines the relative magnitude of the RF field, we can estimate the value of E_{rm} by fitting the calculated B_θ to the measured one. This procedure gives $E_{rm} \sim 16$ V/cm. The absolute value of the radial ponderomotive force is estimated from Figs. 3-19(a) and (b) to be

$$- F_{ir} = (1.6-4.3) \times 10^{-18} \text{ Newton/particle.}$$

On the other hand, the effective gravity is given by

$$\begin{aligned} G &= eT_i/R_c \sim (1.6 \times 10^{-19} \times 20)/4.3 \\ &= 7.4 \times 10^{-19} \text{ Newton/particle.} \end{aligned}$$

Although the value of F_{ir} is only an order of magnitude estimation, it can be seen that the stability condition of Eq.(3-8) could be satisfied in the case of Fig.3-17(c).

C. Improved Confinement

If the flute instability is suppressed, radial particle diffusions should be reduced and, consequently, the plasma

confinement would be improved. We here show a typical example demonstrating an improvement of the plasma confinement. Figure 3-20 shows the oscillogram of (a) the density perturbation of the flute instability measured at $r/r_L = 0.7$, (b) the radial loss flux at the midplane, (c) the radial loss flux near the throat, and (d) diamagnetic signal for $B_M = 0.26$ T, $R = 2.3$, and $n_0 = 5 \times 10^{10} \text{ cm}^{-3}$. The radial loss collector consists of negatively biased plane electrode and grounded mesh grid in front of the electrode. For one axial position, four collectors are set with 90° interval along the azimuthal direction at a radius larger than that of the limiter. The signal shown in (b) or (c) is the sum of ion currents detected by these four collectors. As shown in (a), the flute instability arises just after the gun injection, and by the application of the $m = +2$ RF field the instability amplitude eventually decreases. As a consequence, the radial loss flux reduces to about (20-30)% of the initial value at the midplane [(b)] and to (60-80)% even at the throat where no RF is applied [(c)].

The energy rate equation is written as

$$\frac{dW}{dt} = -P_{\parallel} - P_{\perp} + P_{in} - \frac{W}{\tau} , \quad (3-9)$$

where W is the total plasma energy, P_{in} is the power input from the gun, $P_{\parallel}$ and $P_{\perp}$ are the parallel and perpendicular power conduction losses, respectively, and τ represents the loss time due to charge exchange, collision, and radiation. If we assume that the input power from the gun

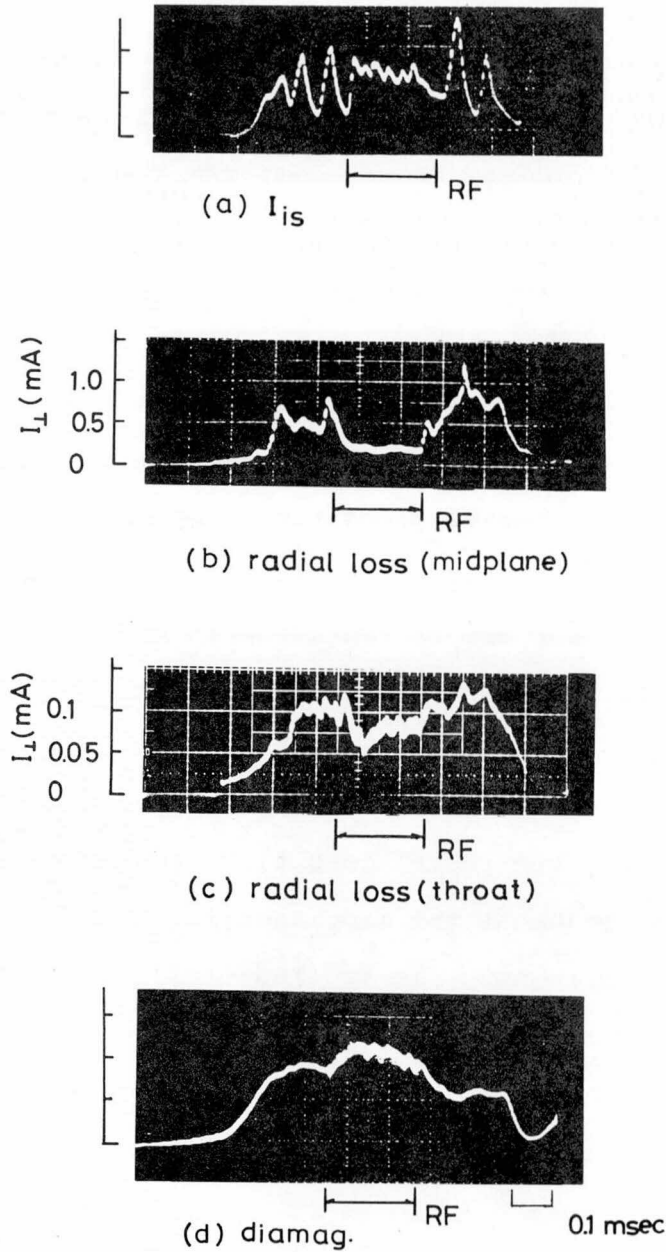


Fig.3-20 Oscilloscope traces of (a) the density fluctuation of the flute instability at $r/r_L = 0.7$, (b) the radial loss flux at the midplane, (c) the radial loss flux near the throat, and (d) diamagnetic signal for $n_0 = 5 \times 10^{10} \text{ cm}^{-3}$, $B_M = 0.26 \text{ T}$, and $R = 2.3$.

and parallel energy loss are not affected by the stabilizing RF, we obtain from Eq.(3-9) that $P_{in} - P_{||} = P_{in0} - P_{||0} = W_0/\tau_0 + P_{\perp 0}$, where suffix 0 represents the value in quasi-steady state. Then the energy rate equation during the RF pulse is given by

$$\frac{dW}{dt} = \frac{W_0}{\tau_0} + P_{\perp 0} - P_{\perp} - \frac{W}{\tau} . \quad (3-10)$$

If we further assume that $W_0/\tau_0 \sim W/\tau$ at the beginning of the RF pulse, we find that

$$\frac{dW}{dt} \sim P_{\perp} - P_{\perp}' . \quad (3-11)$$

We can estimate the value of $P_{\perp}$ and $P_{\perp}'$ from Figs.3-20(b) and (c) to be $P_{\perp} - P_{\perp}' \sim 1.2$ W. On the other hand, the value of the left hand side of Eq.(3-11) can be obtained from a rise of diamagnetic signal due to the stabilization of the flute instability shown in Fig.3-20(d). At a time immediately after switching on the RF pulse, $dW/dt = 2.3$ joule/sec (We have assumed parabolic profile of the plasma energy density.). We can see a fairly good agreement between these two values of "power gain" due to the stabilization obtained in different ways. This fact strongly indicates that the plasma confinement is improved through the stabilization of the flute instability. For later time, Eq.(3-11) no longer holds since the increase of the plasma energy enhances other power losses. This can be seen from a saturation and even decrease of the diamagnetic signal during the latter half of the RF pulse.

In Fig.3-20, the net RF power into the plasma was about 20 W, which is much larger than the power gain of the plasma. This is due to the lack of a mirror confinement in the present system in which the mean free paths for ions and electrons are much shorter than the mirror length. The electrons near the RF antenna are collisionally heated extracting the RF power and are lost through both ends. If the axial confinement is improved by increasing T_i and decreasing neutral gas density, the electron loss could be reduced and the power gain of the plasma would exceed the RF input power.

5. CONCLUSION

In this Chapter, we have investigated the instability arising in the axisymmetric mirror machine HIEI.

The instability observed in a wide range of density has been identified as the flute instability through the comparison of experimentally determined dispersion relation with the theoretical one.

We have demonstrated that the flute instability can be stabilized by applying $m = \pm 2$ RF field only in the bad curvature region. The stabilizing mechanism has been revealed to be the reversal of azimuthal drifts of ions by the radial ponderomotive force. The simple calculation shows that the stabilization occurs when the ponderomotive force exceeds the effective gravity.

It has also been observed that the radial confinement

of the plasma is improved with the stabilization of the flute instability. The increase of the plasma energy during the RF stabilization can be explained by the energy balance equation. At the present stage, however, the power gain of the plasma achieved through the flute stabilization is smaller than the input RF power. This difficulty will be easily overcome by improving the axial confinement in HIEI through reduction of residual neutral gas and ion heating, then the plasma power gain will exceed the RF power.

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CHAPTER 4

SCALING OF RF STABILIZATION

1. INTRODUCTION

Complete axisymmetrization of a tandem mirror requires a method of MHD stabilization other than minimum-B anchors. We have reported in the previous Chapter that the mode-controlled RF field in the ion cyclotron range of frequencies can stabilize the flute instability arising in a wide range of density, from 10^{10}cm^{-3} up to 10^{14}cm^{-3} , resulting in the reduction of radial diffusion [1]. The stabilization mechanism is well explained by the reversal of ion drift caused by the radial ponderomotive force for ions. This has been confirmed by the estimation of the ponderomotive force from the measured value of RF field strength and its comparison with the effective gravity force.

To clarify the utility of this method for stabilizing a fusion plasma, it is necessary to obtain a scaling relation of required RF electric field strength to plasma parameters.

In this Chapter, we perform the RF stabilization of the flute instability, varying the plasma density, the radius of curvature of the magnetic field line, and the ion temperature. It is to be described that the required strength of RF field in vacuum region for the flute stabilization scales as $n_0^{0.3\pm0.05}(T_i/R_c)^{0.4\pm0.1}$, predicting a reasonable value of

1.1 kV/cm for a plasma of $n_0 = 5 \times 10^{13} \text{ cm}^{-3}$, $T_i = 20 \text{ keV}$, and $R_c = 4.3 \text{ m}$.

Following the description of the experimental setup in the next Section, the experimental results of the RF stabilization in various plasma parameters are presented in Section 3A. The detailed measurement of the RF field penetration is shown in Section 3B. Section 4 gives the theoretical calculation and its comparison with the experiment. Conclusions are given in Section 5.

2. DESCRIPTION OF THE EXPERIMENT

The experiment was performed in the axisymmetric mirror machine HIEI, details of which are described in Chap. 3. The major modification of the setup is the installation of heating RF antenna at a place just inside the throat of the gun injection side as shown in Fig.4-1. This antenna produces the rotating RF field of $m = +1$ mode at a frequency slightly lower than the local ion cyclotron frequency. The $m = +1$ mode has been confirmed to give better heating efficiency than the $m = \pm 1$ mode used conventionally [2-4]. We use this antenna to change ion temperature through ion cyclotron heating. The resistance of the antenna was 0.3Ω with no plasma present, while it increased to 0.4Ω when the plasma of $n_0 = 1.5 \times 10^{13} \text{ cm}^{-3}$ was injected. This value was measured at an RF power level of 20 kW. It was observed that the plasma loading resistance decreased for higher RF

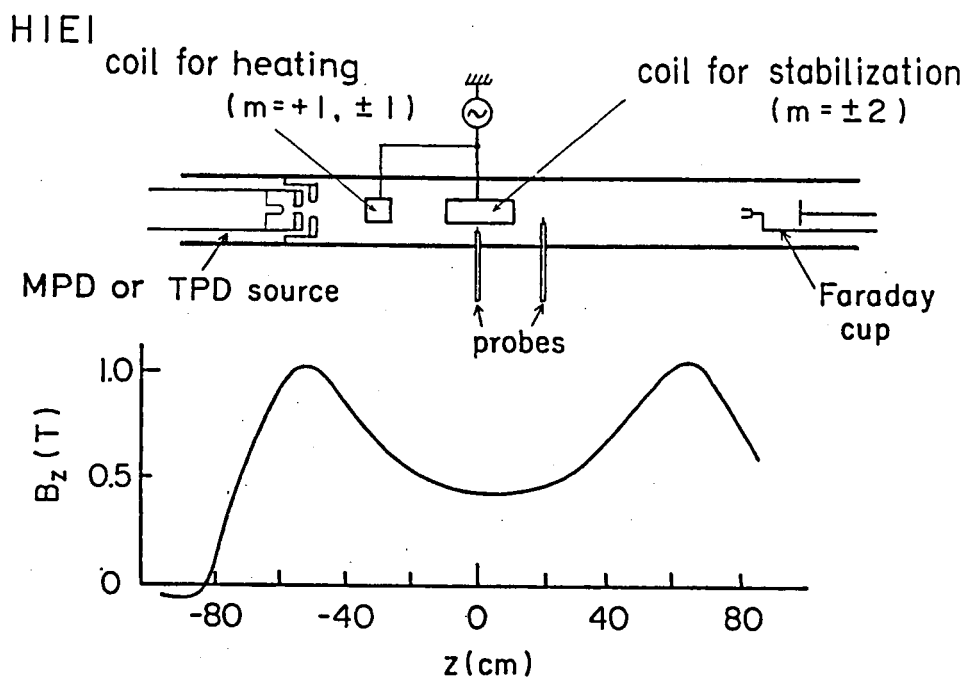


Fig.4-1 The experimental set up.

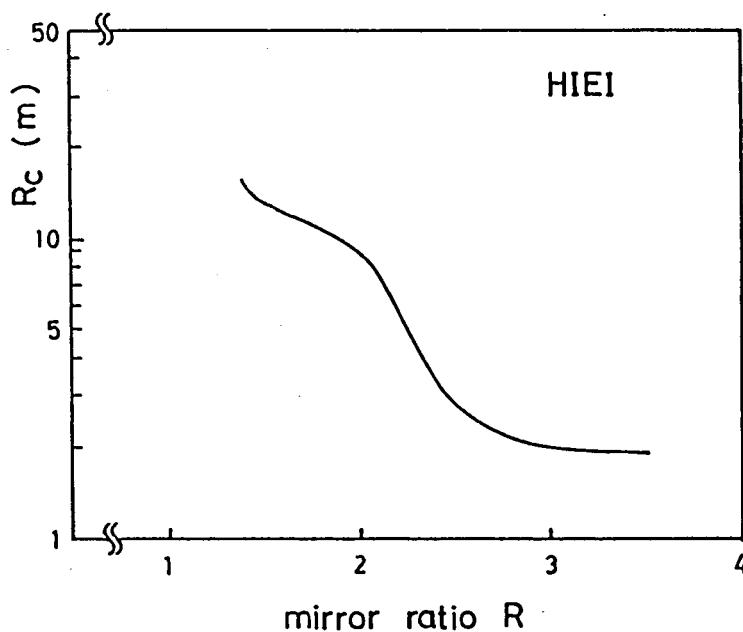


Fig.4-2 The radius of curvature of the magnetic field line at $z = 0$, $r = r_L$ versus the mirror ratio.

power. This fact as well as relatively large charge exchange loss restricted the increase of ion temperature up to 3 times the initial value. The increase of plasma energy was measured by a diamagnetic loop located close to the heating antenna. (12 cm downstream from the heating antenna and 20 cm upstream from the stabilizing antenna.) The RF stabilization was carried out with a density range from 10^{10} cm^{-3} up to 10^{14} cm^{-3} , ion temperature range from 10 eV to 30 eV, and the mirror ratio from 1.4 to 4. Figure 4-2 shows the radius of curvature of the magnetic field line at $(r,z) = (r_L, 0)$ as a function of the mirror ratio.

During the quasi-steady state of 1-1.5 msec after the plasma injection from TPD or MPD gun, the stabilizing RF field of $m = \pm 2$ mode was applied at the midplane for a duration of 0.5 msec. The frequency of the stabilizing RF was 10 MHz (6 MHz in some cases) and the ratio of it to the local ion cyclotron frequency at the midplane ranged 1.5 to 3.5. A change of this ratio affected little the stabilizing characteristics provided that the antenna current I_{RF} in the stabilizing antenna was kept constant. Stabilizing effects in various plasma parameters are investigated by measuring power spectrum of the flute fluctuation at various positions.

3. EMPIRICAL SCALING OF RF STABILIZATION

A. Dependence on Plasma Parameters

As described in Chap. 3, low frequency, large amplitude

fluctuation observed in HIEI has been identified as the flute instability for a density range 10^{10}cm^{-3} to 10^{14}cm^{-3} . We apply the $m = \pm 2$ RF field at the midplane, and observe power spectra of density fluctuations of the flute instability. The spectrum has a sharp peak at a certain frequency, and the peak decreases with increasing the RF current in the stabilizing antenna. Let us define δ as the ratio of the spectral amplitude of the instability with stabilizing RF to that without RF. Figure 4-3 shows the value of δ as a function of square of I_{RF} for a plasma with $n_0 = 2 \times 10^{10}\text{cm}^{-3}$, $T_i = 15\text{ eV}$, $B_M = 0.4\text{ T}$, and $R = 2.3$. For $\delta < 0.1$, the curve in Fig.4-3 has large ambiguity since the spectral amplitude of the instability is too small to be distinguished from background noise level. It is seen that δ can be expressed as $\delta = \exp(-C_0 \times I_{\text{RF}}^2)$ for $\delta > 0.1$, with C_0 being a constant.

We plot in Fig.4-4(a) the value of I_{RF} necessary for obtaining $\delta = 0.5, 0.3$, and 0.2 versus the plasma density in the low-density regime. The value of I_{RF}^2 increases proportionally to $n_0^{0.6}$ for this density range. In Fig.4-4(b), the value of I_{RF} for $\delta = 0.2$ in both the low- and high-density regimes are given versus the plasma energy density, where $T_i + T_e$ is approximately kept constant to be (20-25) eV. Here, the value 0.2 is chosen in order to determine δ within an accuracy of $\pm 10\%$. For $n_0 = 3 \times 10^{13}\text{cm}^{-3}$, the value of I_{RF} was 30 A_{rms} , and the net power into the plasma was 12 kW. Although there is no experimental point in the middle pressure range in Fig.4-4(b), it is found that the extrapolated line for $\delta = 0.2$ in Fig.4-4(a) connects smoothly

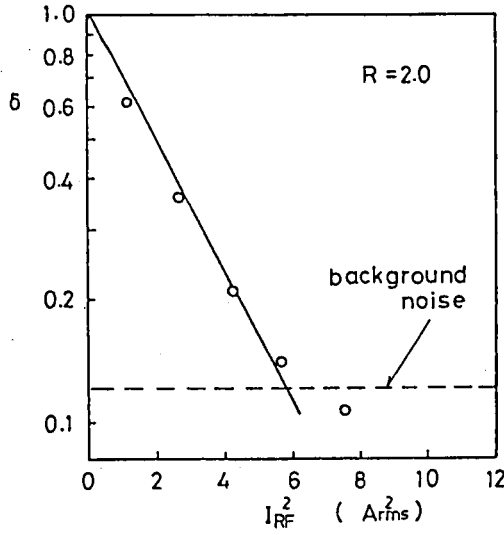


Fig.4-3 An example of semi-logarithmic plot of the suppression ratio versus I_{RF}^2 for the low-density regime.

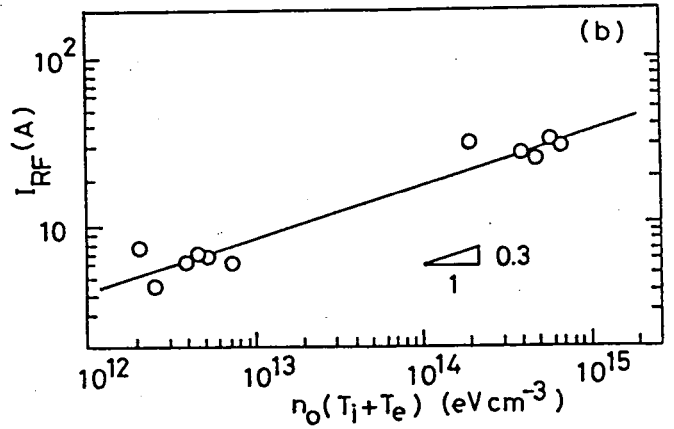
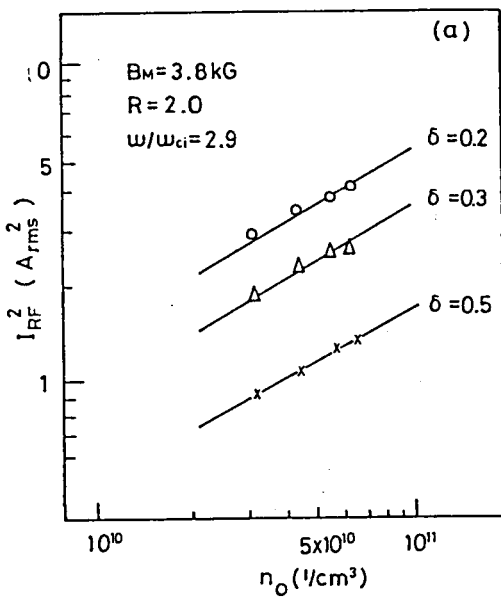


Fig.4-4 (a) RF current required for $\delta = 0.5, 0.3$, and 0.2 versus the density for the low-density regime. (b) RF current required for $\delta = 0.2$ versus the plasma energy density for low- and high-density regimes.

to the experimental data for $n_0(T_i + T_e) > 10^{14} \text{ eVcm}^{-3}$.

Therefore, from Figs.4-4 (a) and (b), we obtain a scaling law that I_{RF} necessary for $\delta = 0.2$ is proportional to $n_0^{0.3 \pm 0.05}$ for a density up to 10^{14} cm^{-3} .

The most important parameter that affects growth of the flute instability is the effective gravity, which is equal to eT_i/R_c with R_c being the radius of the curvature of the magnetic field line. We proceed to measure the effect of R_c and T_i upon the RF stabilization. The value of R_c is given in Fig.4-2 as a function of the mirror ratio.

In Fig.4-5, the suppression ratio δ is plotted versus I_{RF}^2 for the mirror ratio of 1.5, 2.0, and 2.5. Other parameters are $B_M = 0.63 \text{ T}$, $T_i = 10 \text{ eV}$, and $n_0 = (3.2-3.5) \times 10^{13} \text{ cm}^{-3}$. It is seen that larger I_{RF} is required for the stabilization of the plasma with larger mirror ratio. The value of δ is limited to 0.5 for $R = 2.5$ because the plasma radius at the midplane for this R is so large that the direct contact of the plasma to the stabilizing RF antenna causes a breakdown especially for larger I_{RF} .

In order to see the T_i -dependence of the stabilization, we perform ICRF heating as described in Section 2. The ion temperature was doubled for the net input power of 4.5 kW from the heating RF antenna. Figure 4-6 compares the characteristics of the RF stabilization with and without the ICRF heating. When $T_i \sim 10 \text{ eV}$, about $19 A_{\text{rms}}$ of I_{RF} is sufficient for obtaining the suppression ratio of 0.2. If T_i is increased to $\sim 25 \text{ eV}$, the required I_{RF} for $\delta = 0.2$ exceeds $25 A_{\text{rms}}$.

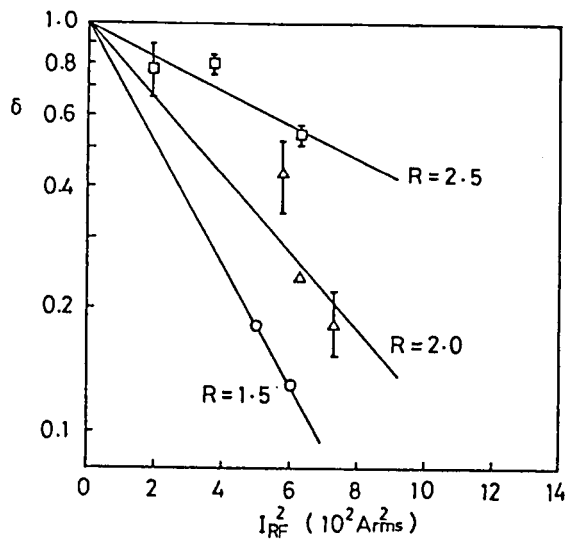


Fig.4-5 Suppression ratio versus I_{RF}^2 for $R = 1.5$, 2.0, and 2.5 in the high-density regime.

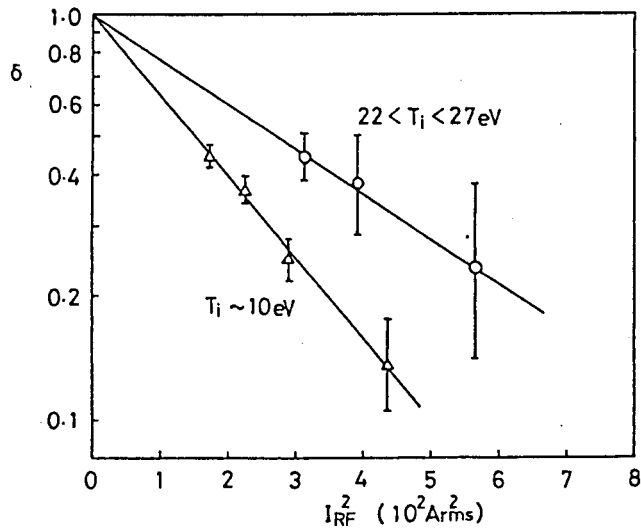


Fig.4-6 Suppression ratio versus I_{RF}^2 for $T_i = 10 eV$ and $T_i \approx 25 eV$ in the high-density regime.

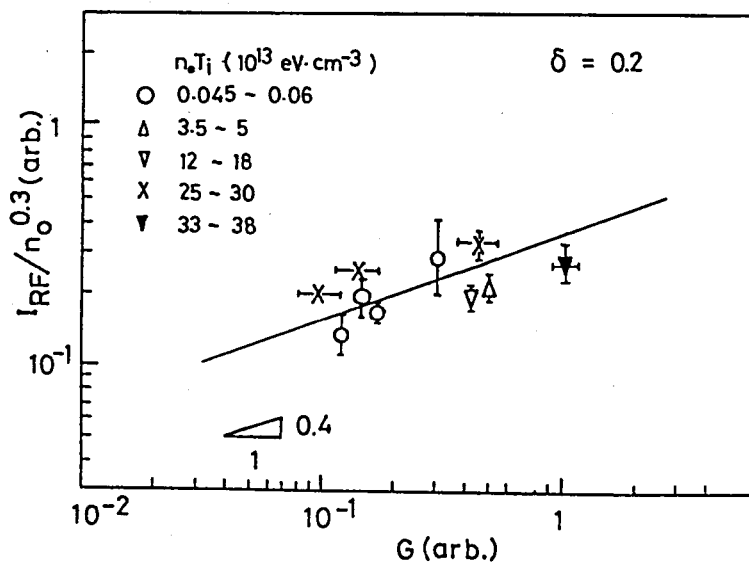


Fig.4-7 Normalized RF current necessary for $\delta = 0.2$ as a function of the effective gravity. The experimental points for different $n_0 T_i$ fall on the same line.

The experimental data in Figs.4-5 and 4-6 and those obtained in different conditions of $n_0 T_i$ are plotted in Fig.4-7, where the normalized value of I_{RF} necessary for obtaining $\delta = 0.2$ is given as a function of the effective gravity G . The value of I_{RF} is normalized by $n_0^{0.3}$ according to the result in Fig.4-4. It is clearly seen that the experimental data for different $n_0 T_i$ fall on the line. By combining the results in Figs.4-4 and 4-7, the empirical scaling formula;

$$I_{RF} \sim 0.3 n_0^{0.3 \pm 0.05} (T_i/R_C)^{0.4 \pm 0.1} \quad (4-1)$$

is obtained, where I_{RF} , n_0 , T_i , and R_C are in the unit of kA_{rms} , $10^{13} cm^{-3}$, keV, and meter, respectively.

B. Penetration of RF Fields

We measured the RF magnetic field generated by the $m = +2$ RF antenna by using a shielded loop of 2.5 mm in diameter covered by a glass tube, which is inserted at the midplane. In vacuum, $B_{\perp}/I_{RF}$ was maximum at $r = 0$ and decreases monotonously with r as shown in Fig.4-8(a). This kind of field pattern is different from that of $m = +2$ or -2 eigenmode, indicating that the antenna near field dominates the radiated field. When the plasma with $n_0 = 5 \times 10^{13} cm^{-3}$ is present, the measured field has a well-defined shape of $m = +2$ or -2 mode, as shown in Fig.4-8(b).

We calculate the dispersion relation and RF field distribution in a cylindrical, uniform, sharp boundary plasma

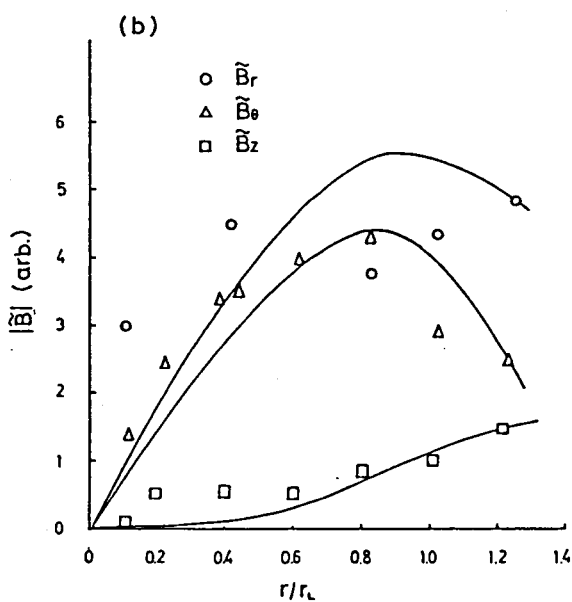
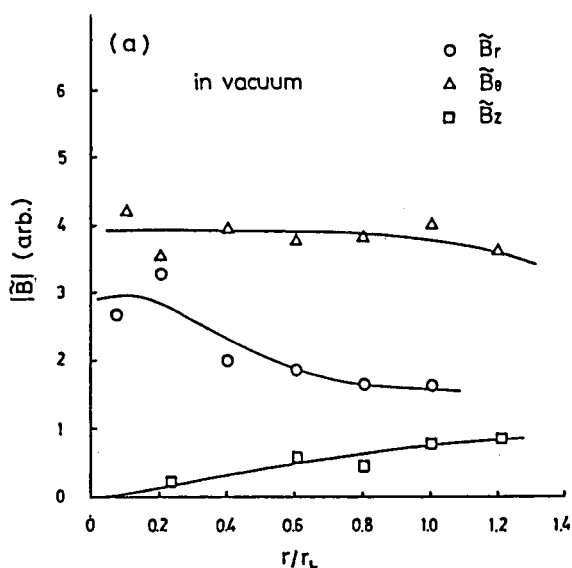


Fig.4-8 (a) Radial distribution of the RF magnetic fields in vacuum. (b) RF magnetic fields in the plasma; solid curves denote the calculated value of the field amplitudes.

with parameters $n_0 = 5 \times 10^{13} \text{ cm}^{-3}$, $\omega/\omega_{ci} = 2.4$, $T_e = 20 \text{ eV}$, $T_i = 0$, and the plasma radius $p = 3 \text{ cm}$. We assume that $2\pi/k_{\parallel}$ is constant and fixed to be 0.6 m , which is twice the axial length of the antenna. The dispersion relation shows that the RF field consists of the fast mode with $k_{\perp 1} = 80 \text{ m}^{-1}$ and the slow mode with $k_{\perp 2} = 4 \times 10^3 \text{ m}^{-1}$. The field distribution of the latter mode is highly oscillatory with respect to r , and is quite different from the measured pattern shown in Fig.4-8(b). Therefore, we neglect this mode, and calculate the field distribution of the fast mode. Figures 2-5 (d) and (e) in Chap. 2 show the electric field distribution of $m = +2$ and -2 fast mode, respectively. The solid curves in Fig. 4-8(b) represents the calculated magnetic field distribution for the $m = -2$ fast mode. The value of numerical constants of the field amplitude is so chosen that the amplitude of B_0 at the peak fits to the experimental value. It seems that the measured magnetic field pattern is well approximated by the calculated one. The discrepancy for small r may come from the fact that the $m = 0$ mode is partially excited in the real experiment. (Remember that the $m = \pm 2$ RF antenna generates the $m = 0$ mode simultaneously as described in Chap. 2.) The calculated magnetic field distribution for the $m = +2$ fast mode is very similar to that of the $m = -2$ fast mode, except that the phase relation between B_r and B_0 is inverted. So, we cannot identify the excited field whether it is of $m = -2$ or $+2$ mode, because the phase difference of the magnetic fields is not measured. There is, however, no experimental report showing that the fast wave of positive m is

excited in a laboratory plasma. The possibility of the excitation of the $m = +2$ fast wave would be excluded. There is another possibility that the measured RF field is formed by the $m = -2$ fast wave and the $m = +2$ quasi mode. In this case, B_r should be zero on $\theta = \text{const.}$ plane including the axial antenna current, and maximum if θ differs by 45° . The position of the magnetic probe was very close to the plane which included the axial antenna current. Therefore, B_r should be very small if the $m = \pm 2$ RF field is formed. This is not the case. We could say that the $m = -2$ fast wave is excited in the experimental condition shown in Fig. 4-8(b). For a density lower than $0.5 \times 10^{13} \text{ cm}^{-3}$, the fast wave becomes cutoff, and the excited field will be an evanescent mode. It should be emphasized that, for the RF stabilization, it is not essential whether the RF field is of $m = -2$ or $+2$ or ± 2 mode as will be shown later. So, it is sufficient to confirm from Fig. 4-8(b) that the stabilizing RF antenna indeed excites the RF field of $|m| = 2$ mode.

Figure 4-9 shows the relation between $|B_\theta|^2 / I_{\text{RF}}^2$ and the plasma energy density, where B_θ is measured at $r/r_L = 0.8$ for $\omega/\omega_{ci} = 2.4$. It is evident that the RF field is screened heavily for higher density. ($T_i + T_e$ is approximately constant.) If we take this effect into account, we obtain

$$B_\theta^P \propto n_0^{-0.35 \pm 0.05} I_{\text{RF}}^2, \quad (4-2)$$

where B_θ^P is the absolute value of B_θ at $r/r_L = 0.8$ (periphery of the plasma column).

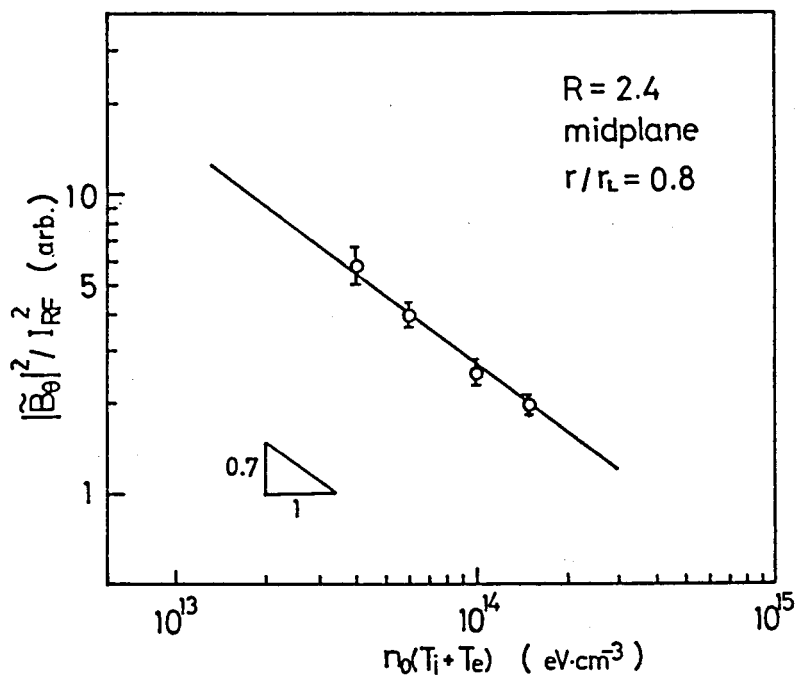


Fig.4-9 Squared RF magnetic field divided by I_{RF}^2 versus the plasma energy density. ($T_i + T_e$ is nearly constant.)

4. COMPARISON WITH THEORY

In the previous Section, an empirical scaling formula of the RF stabilization has been obtained. We, then, consider the RF stabilization scheme theoretically, starting from the simple model shown in Fig.4-10. (See also Fig.3-13 in Chap. 3.) Since we are much more interested in the result in the high-density regime than in the low-density regime, the following discussions are devoted to the former density regime, if otherwise noted.

In the case of no stabilizing RF field applied, the dispersion relation of the flute mode is given by Eq.(3-2). We approximate the effect of the RF field by a change in effective gravity in the region $r_s < r < r_0$ as shown in Fig.4-10. In this case, we replace G in Eq.(3-2) by $G + F_p/2$ with F_p being the radial ponderomotive force existing only in region 2 [5]. This situation is, of course, not exact because the radial ponderomotive force affects ions for almost whole of the plasma radii as shown in Fig.3-19. But, as the direction of F_{ir} in Fig.3-19 is favorable for the RF stabilization, it may not be an overestimate of the F_p to assume $F_p = F_{ir}(r = \frac{r_s + r_0}{2})$ with $r_0 = p$. Assuming that $\zeta \gg 1$, and neglecting the finite Larmor radius effect, we solve Eq.(3-2) to find that

$$\Omega_i^2 = \left[\frac{MeT_i}{m_i R_c} + \frac{MF_p}{2m_i} + \{(M-1) + 2\rho_2\} r_s \omega_{0s}^2 \right] \frac{\rho_1 - \rho_2}{r_s} \quad (4-3)$$

The stabilization of the flute mode requires $\Omega_i \rightarrow 0$. Since the first term in the square brackets in Eq.(4-3) is much

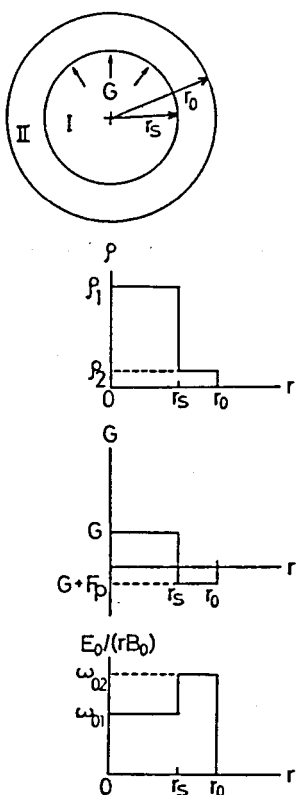


Fig.4-10 Theoretical model of the plasma cross section. The RF ponderomotive force acts as to reduce G in the region $r_s < r < r_0$.

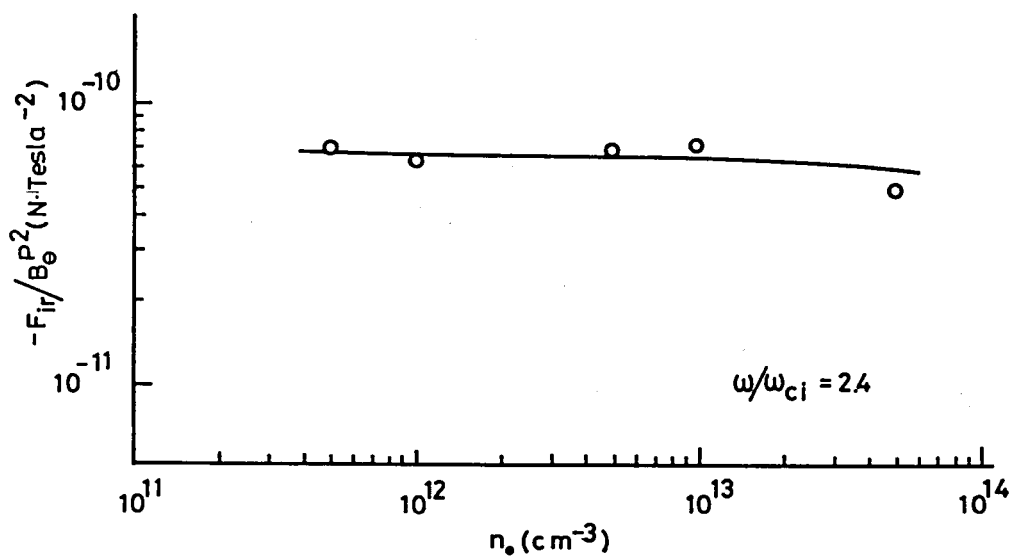


Fig.4-11 Calculated radial ponderomotive force divided by the square of B_0^p for $m = -2$.

larger than the last term under the present experimental condition, the marginal stability is achieved when $-F_p = 2eT_i/R_c$. As the ponderomotive force is considered to be proportional to the square of the RF magnetic (or electric) field in the plasma, F_p can be written as

$$F_p = F_{ir} \left(\frac{r_s + r_0}{2} \right) = -\xi |B_\theta^p|^2, \quad (4-4)$$

where B_θ^p is the azimuthal component of the RF magnetic field in the plasma region at $r/r_L = 0.8$ and ξ is a form factor, which will depend on n_0 , T_i , and m even if the geometries of the plasma and the stabilizing RF antenna are fixed. Here, there is no reason why r/r_L should be 0.8. We have chosen this value only because B_θ is measured at this position in the experiment.

The value of ξ is calculated as follows: We solve the dispersion relation for a given density and obtain the RF field distribution using the parameters $\omega/\omega_{ci} = 2.4$, $k_z = 10.5 \text{ m}^{-1}$, $p(=r_0) = 3.0 \text{ cm}$, $T_i = 0$, and $T_e = 20 \text{ eV}$. The value of F_{ir} is obtained from Eq.(3-5) for the calculated RF field distribution. Then, we divide this value by the square of B_θ^p which can be read from the field distribution. The result is given in Fig.4-11, where ξ for the $m = -2$ mode is plotted as a function of the density. We note that the density dependence of ξ is very weak though the RF field changes its mode from evanescent mode to propagating mode as n_0 exceeds $5 \times 10^{12} \text{ cm}^{-3}$. We also calculate the value of ξ for the $m = +2$ mode to find that it is smaller than the value in Fig.4-11

by only a factor of ~ 2 , and is very weakly dependent on n_0 . It is also expected that ξ is insensitive to T_i at least for the range used in the experiment. Thus, we can treat ξ as a constant value, and the stability condition yields

$$B_0^P = \left(\frac{2e}{\xi}\right)^{0.5} \left(\frac{T_i}{R_c}\right)^{0.5} . \quad (4-5)$$

This expression predicts that the RF magnetic field necessary to stabilize the flute instability is proportional to $G^{0.5}$ regardless of the density. It simply means that the larger ponderomotive force is required to reverse the azimuthal drift of ions due to larger effective gravity. We must note that Eq.(4-5) does not hold when the sum of the first two terms in the square brackets in Eq.(4-3) becomes comparable to the last term for large F_p .

On the other hand, Eqs.(4-1) and (4-2) give the empirical scaling formula for the flute stabilization. As the RF electric field in the vacuum region (outside the plasma column), E^V , is proportional to I_{RF} , Eq.(4-1) is rewritten as

$$E^V \sim 0.3 n_0^{0.3 \pm 0.05} (T_i/R_c)^{0.4 \pm 0.1} . \quad (4-6)$$

Here, the relation between E^V and I_{RF} has been calculated using Maxwell's equations. This equation gives the RF electric field in vacuum region necessary for the flute stabilization with $\delta = 0.2$.

Combining Eqs.(4-1) and (4-2), we obtain the empirical relation between B_0^P necessary for the flute stabilization with

$\delta = 0.2$ and the plasma parameters;

$$B_{\theta}^P \propto n_0^{-0.05 \pm 0.1} (T_i/R_c)^{0.4 \pm 0.1}. \quad (4-7)$$

Comparing Eq.(4-7) with Eq.(4-5), one finds that the powers of n_0 and (T_i/R_c) which are determined experimentally are very close to those predicted by the theory. [The powers of n_0 is zero in Eq.(4-5).] We have already obtained the semi-quantitative agreement of the experiment and the theory in Section 4 of Chap. 3 concerning the absolute value of B_{θ}^P necessary to stabilize the flute instability at $n_0 = 5 \times 10^{13} \text{ cm}^{-3}$. Therefore, it is concluded that Eq.(4-7) agrees with Eq.(4-5) within an accuracy of the experiment.

Having established the scaling formula of the flute stabilization [Eq.(4-6) or Eq.(4-7)], we can predict a necessary value of the RF electric field for the RF stabilization of a fusion-oriented plasma from Eq.(4-6) to be $E^V \sim 1.1 \text{ kV/cm}$ for $n_0 = 5 \times 10^{13} \text{ cm}^{-3}$, $T_i = 20 \text{ keV}$, and $R_c = 4.3 \text{ m}$. The value of E^V is easy to be realized and this method is hopeful for an application to a mirror fusion reactor. Furthermore, in a hot plasma the required RF field strength should be smaller by additional effects of finite Larmor radius and thin RF field penetration depth.

There may arise a question that Eq.(4-6) or Eq.(4-7) still holds or not for ion temperatures and the plasma radius larger than those in our experiment. It was reported that, in the tandem mirror Phaedrus, the central cell plasma could be stably sustained by the RF ($\omega/\omega_{ci} \gtrsim 1$) even when the plug

cell plasmas are turned off. If $\omega < \omega_{ci}$, large amplitude MHD instabilities were observed to reduce the central cell beta significantly. Having referred to our results in HIEI, Hershlowitz et al. [6] concluded that, in the case of $\omega > \omega_{ci}$, the stability of the axisymmetric mirror plasma in central cell stand alone mode was maintained by the radial ponderomotive force. The estimated value of the ponderomotive force was just enough to overcome the curvature force in a plasma of $n_0 = 3 \times 10^{12} \text{ cm}^{-3}$, $T_i \sim 60 \text{ eV}$, and $r_0 \gtrsim 15 \text{ cm}$. Although their RF antenna configuration is different from ours, the results in Phaedrus strongly support our scaling formula, and confirm that the RF stabilization method is still effective for these values of T_i and r_0 . It is, however, necessary to confirm the validity of Eq.(4-6) or Eq.(4-7) for much higher T_i .

There are two more problems which remain to be resolved. First, which mode of RF is the best for the stabilization? In the experiments described in Chapters 3 and 4, the modes of the RF were the propagating fast wave for the high-density regime, and the evanescent mode for the low-density regime. If we use the propagating mode, the RF field can be large in the plasma producing large ponderomotive force, but care should be taken not to lose the RF energy as it propagate away towards the ends. For the evanescent mode, although the efficiency of the RF field excitation is smaller, the problem of the RF energy loss will be insignificant.

We have used the $m = \pm 2$ RF field, which is linearly polarized, for stabilization. As will be shown in the next Chapter,

the linearly polarized field produces radial convective loss of the plasma. The convective loss is, however, negligibly small as compared with the loss caused by flute instabilities. Therefore, the convective loss has no effect on the RF stabilization experiments described in Chaps. 3 and 4. But once a complete stabilization of the flute instabilities is achieved, we must use a circularly polarized RF field for further improvement of radial confinement. The RF antenna shown in Fig.2-8(b) in Chap. 2 is suitable for this purpose.

Second, can the RF stabilization method improve the plasma confinement time in the axisymmetric mirror to a classical value ? As described earlier, the residual neutral gas pressure was relatively high and T_i was less than several tens of eV in HIEI, resulting in the lack of a mirror confinement. Therefore, we could not discuss about overall particle and energy confinement times using the experimental results. We did observe the decrease of the radial diffusion loss followed by the increase of the diamagnetic signal when the RF stabilization was applied to HIEI. It was also observed in Phaedrus [6] that the central cell beta increased with the RF stabilization compared to the unstable case. Judging from these experimental results, we can surely expect an improvement of the particle and energy confinement times by the RF stabilization. The experiment of the RF stabilization of a mirror-confined plasma is now in progress.

5. CONCLUSION

The RF stabilization of a plasma in the axisymmetric mirror HIEI was investigated in various plasma parameters.

We have obtained the empirical scaling formula of the RF stabilization that the required RF field strength increases proportionally to $n_0^{0.3 \pm 0.05} (T_i/R_c)^{0.4 \pm 0.1}$. This scaling formula has been revealed to agree fairly well with the theoretical prediction that the radial RF ponderomotive force must overcome the curvature force for stability.

From this scaling, a reasonable value of the RF electric field of 1.1 kV/cm is estimated for the RF stabilization of a fusion-oriented plasma with $n_0 = 5 \times 10^{13} \text{ cm}^{-3}$, $T_i = 20 \text{ keV}$, and $R_c = 4.3 \text{ m}$.

This method of stabilization will be effective for increasing the central cell beta of a tandem mirror and could be an alternative for non-axisymmetric MHD anchors.

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CHAPTER 5

CONVECTIVE PLASMA LOSS CAUSED BY AN ION CYCLOTRON RF FIELD AND ITS ELIMINATION BY MODE CONTROL

1. INTRODUCTION

RF fields in ion cyclotron range of frequencies (ICRF) have been successfully employed in heating ions in linear and toroidal devices. In most experiments, slow and fast waves are launched by half-turn coil(s) [1-3] or a Kharkov-type coil [4], which produces RF fields of $m = +1$ and -1 [and 0 for half-turn coil(s)] azimuthal modes simultaneously. Here, RF fields are assumed to vary as $\exp[j(m\theta - k_z z + \omega t)]$ in cylindrical geometry. For large amplitude waves, there may arise deleterious effects due to nonlinear interactions between these fields and a plasma. Chen and Etievant [5] calculated particle drifts caused by ponderomotive force of excited waves. The ponderomotive effects will become important if the excursion velocity of particles in wave fields exceeds the thermal velocity.

Although nonlinear processes in lower hybrid heating [6, 7] have been investigated in detail, few experimental reports are available concerning the nonlinear effects in ICRF heating.

In this Chapter, we wish to present the experimental evidence that the RF field of $m = \pm 1$ azimuthal mode, which is conventionally used in ICRF heating, produces convective cross-field plasma loss.

The convective loss is produced through particle drifts due to the ponderomotive force of the large amplitude $m = +1$ RF field. By utilizing the $m = +1$ circularly polarized RF field, we can achieve higher overall heating efficiency in consequence of eliminating the convective loss.

Theoretical consideration of the convective motion in an RF field is presented in Sec. 2. Section 3 includes the description of the experimental setup. The experimental results which demonstrate the convective loss in ICRF heating are given in Sec. 4 together with the comparative measurement of the heating efficiency. Section 5 contains conclusions.

2. THEORY OF CONVECTIVE LOSS

The mechanism producing ion convection may be understood from the following fluid theory in which the quasilinear correction to the linear solution is taken into account [5]. Let us consider a cylindrical, homogeneous, collisionless plasma with radius r_L immersed in a static magnetic field B_0 in the positive z direction. If the current density in a Kharkov-type RF antenna (see Fig. 2-7 in Chap. 2) is approximated by $J_\theta = J_0 \delta(r-s) \sin\theta \cos(-k_z z + \omega t)$ with s being the coil radius, the RF electric field in the plasma is a superposition of the fields of $m = +1$ and -1 modes, and may be represented by

$$E_r^{(1)} = [E_1(r) + E_2(r)] \cos\phi_1 + [E_3(r) + E_4(r)] \cos\phi_2$$

$$E_{\theta}^{(1)} = [-E_1(r) + E_2(r)] \sin \phi_1 + [-E_3(r) + E_4(r)] \sin \phi_2$$

$$E_z^{(1)} = 0, \quad (5-1)$$

where $\phi_1 = \theta - k_z z + \omega t$, $\phi_2 = -\theta - k_z z + \omega t$ with k_z being the axial wave number, and E_1 (E_3) and E_2 (E_4) are, respectively, the amplitudes of the left and right hand circularly polarized field components of the $m = +1$ (-1) mode. Using Eq.(5-1), we can obtain linear and second order (in the field amplitude) solutions for particle motions. To second order, the ponderomotive force gives rise to a quasi steady drift of ions in the direction perpendicular to both the force and $\vec{B}_0$ as given by Eq.(2-12) in Chap. 2. For the field of Eq.(5-1), no charge separation arises due to the second order drifts of ions and electrons. We can calculate the quasi steady flux of ions;

$$\vec{\Gamma}_i = \langle n^{(1)} \vec{v}_i^{(1)} \rangle + n_0 \vec{v}_i^{(2)}, \quad (5-2)$$

where n_0 and $n^{(1)}$ are the zeroth- and first-order density, respectively, $\vec{v}_i$ is the ion velocity, and the brackets denote a high-frequency average. After some algebra, we obtain the radial component of the quasi steady ion flux, which is given by

$$\Gamma_{ir} = 2 \left(\frac{e}{m_i} \right)^2 \frac{n_0}{\omega_{ci}} \frac{1}{r} \sin 2\theta \left[\frac{E_1 E_3}{(\omega - \omega_{ci})^2} + \frac{E_2 E_4}{(\omega + \omega_{ci})^2} \right], \quad (5-3)$$

where m_i is the ion mass and ω_{ci} is the ion cyclotron angular frequency. This equation has been derived under the condition

that the RF field is traveling in the z direction. In the case of an axially standing RF field, Γ_{ir} is represented by Eq.(5-3) multiplied by $[\omega^2 + \omega_{ci}^2/(2\omega\omega_{ci})]\cos^2 k_z z$. Equation (5-3) shows that the large cross-field ion flux is produced by the $m = \pm 1$ RF field at fundamental ion cyclotron resonance. We must note that the $m = +1$ and -1 RF fields in Eq.(5-1) are not necessarily the wave fields.

Since we have assumed that $n_0(r) = 0$ for $r > r_L$, the net plasma loss is obtained by integrating Eq.(5-3) over that half-wave of the sine which leads to flux toward the outside.

In order to eliminate this convective plasma loss, we design the RF antenna which can produce the $m = +1$ or -1 RF field selectively. [see the antenna (c) in Fig.2-7 of Chap. 2] In this case, we have $E_3 = E_4 = 0$ (for the $m = +1$ mode) or $E_1 = E_2 = 0$ (for the $m = -1$ mode), then Eq.(5-3) predicts that no radial flux of ions is produced.

3. EXPERIMENTAL SETUP

The experiment was performed in the single-ended Q-machine HIEI-Q illustrated in Fig.5-1. A potassium plasma is continuously produced by thermal ionization of the neutral K-beam on a tungsten plate heated to about 2200°K, and forms a column of about 1.2 m long and 3 cm in diameter. A uniform static magnetic field up to 1.0 T confines the plasma radially. Initial temperatures of ions and electrons are both about 0.25 eV, and electron density is in the range of 10^9

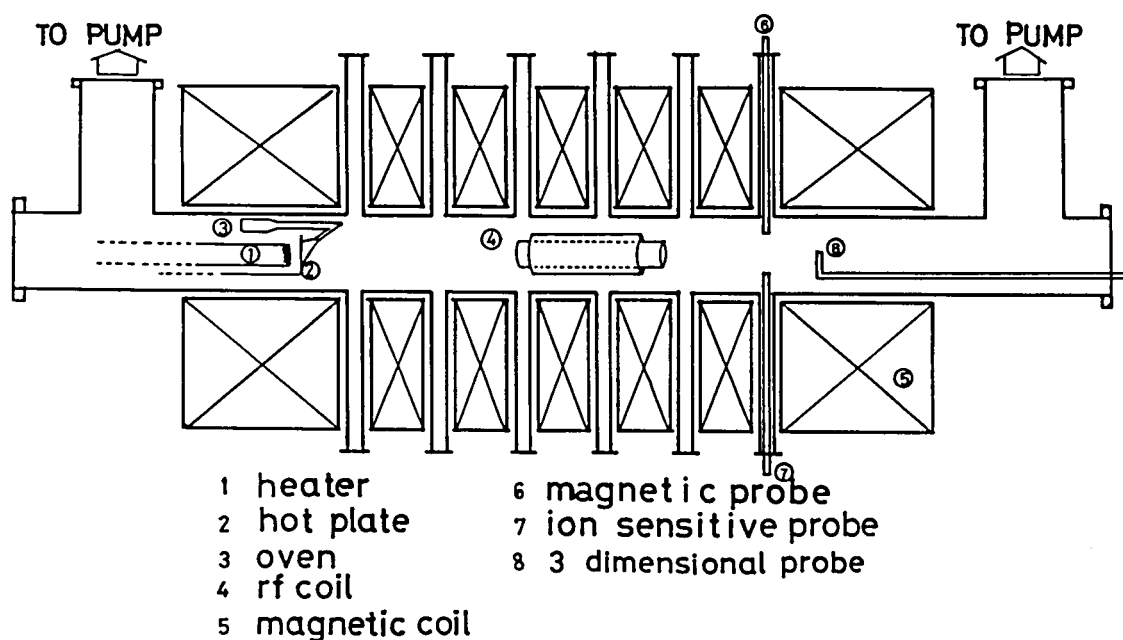


Fig.5-1 Schematic view of HIEI-Q.

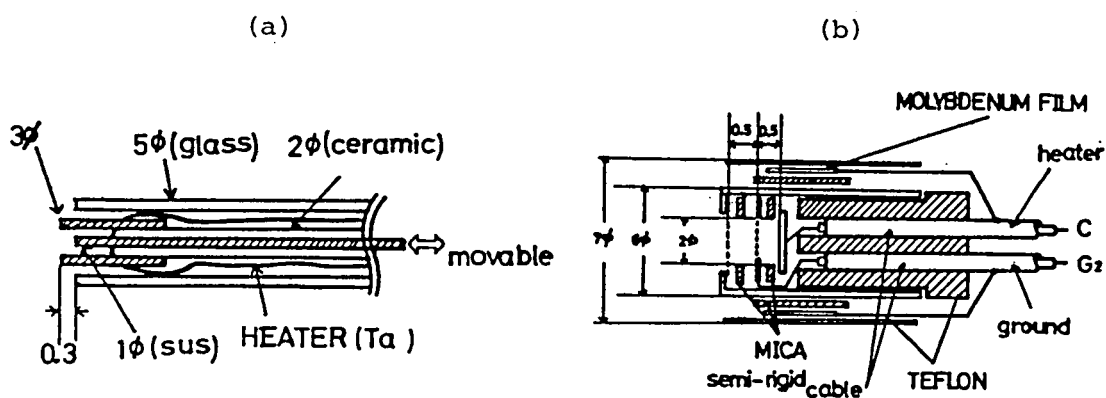


Fig.5-2 Structures of the ion sensitive probe (a) and the Faraday cup (b).

to 10^{11} cm^{-3} with the background neutral pressure less than 2×10^{-6} Torr. For the plasma parameters described above, mean free paths of ion-neutral and ion-electron collisions are much longer than the length of the plasma column.

In order to measure the perpendicular and parallel ion temperatures, we constructed an ion sensitive probe and a Faraday cup. These probes should be as small as possible not to disturb the plasma, and should have a heater to avoid contaminative attachment of potassium atoms. The ion sensitive probe which is shown in Fig.5-2(a), has the structure so as to receive ions selectively using the difference of the Larmor radii. The configuration of the Faraday cup is shown in Fig.5-2(b). Considering the Debye length of the plasma, we adopted grids with 100 meshes per inch to apply a uniform retarding potential.

In the experiments described later, it is sometimes necessary to measure two dimensional profiles of the plasma parameters. We developed a probe mechanism which enables us to set the probe tip at any position in the plasma column. Figure 5-3(a) shows the bird's-eye view of the three dimensional probe (3D probe) mechanism, where the single probe is mounted on a pantograph which slides back and forth along the guide rail laid in the axial direction. The lower pivots of the pantograph are coupled with the ends of the inner and outer shafts which are extended out of the vacuum chamber through a gauge port. The relative translational movement of the inner and outer shafts makes the probe go up or down, and the rotation of the shafts results in the probe tip move-

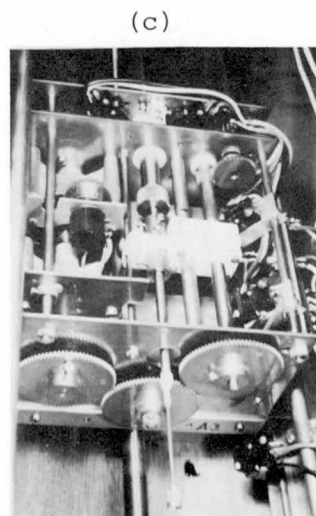
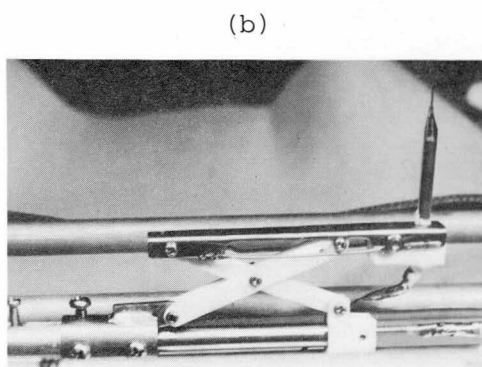
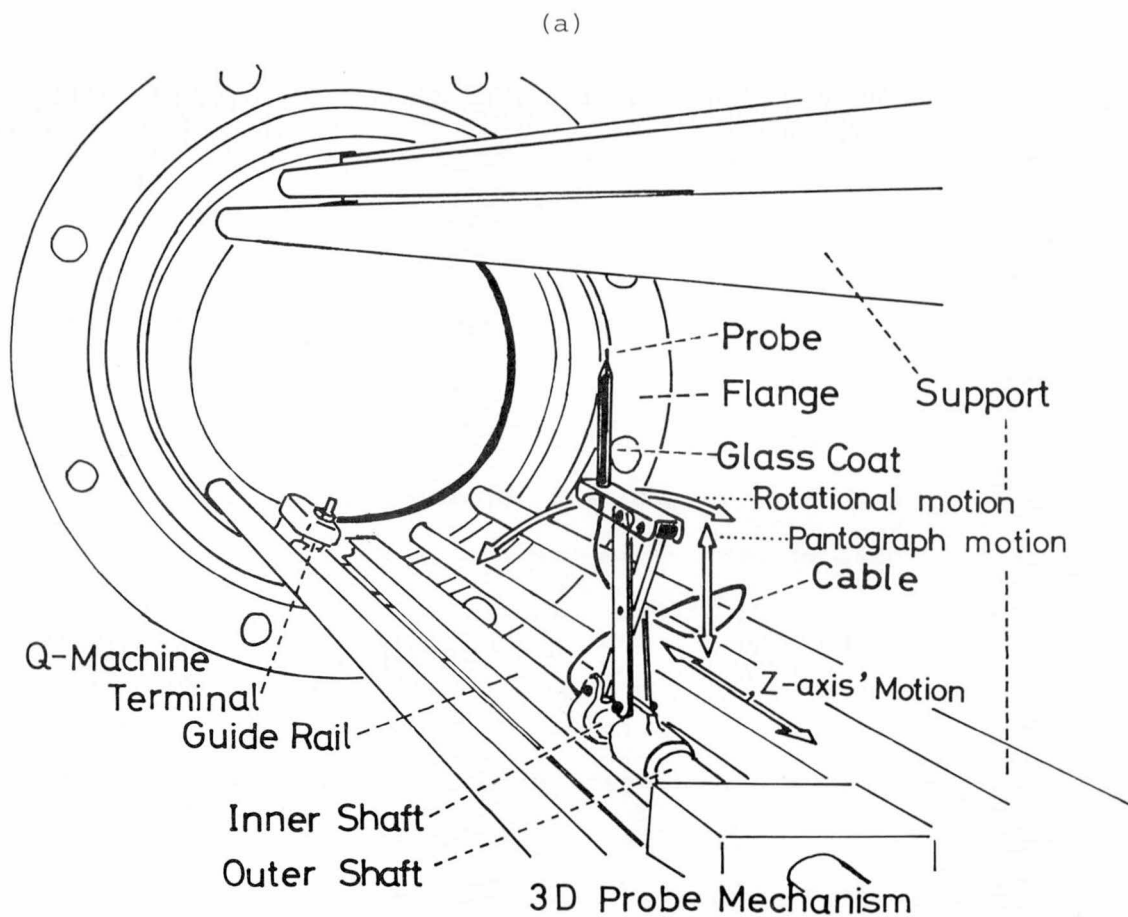


Fig.5-3 (a) Bird's-eye view of the 3D probe,
 (b) Photograph of the probe tip, and (c) Photograph
 of the gear mechanism.

ment along an arc. The axial position of the probe is easily changed by sliding the shafts. By combining these three driving mechanisms, we can set the probe tip at a desired position. The photograph of the probe and the gear box which drives the inner and outer shafts are shown in Figs.5-3(b) and (c), respectively.

The RF antennae are installed 50 cm downstream from the hot plate with the electrostatic shield inside. The antenna A in the inset of Fig.5-4 is the Kharkov-type, and consists of two half-turn loops. Each loop is made of 8 turn copper wire of 2 mm in diameter and has a length of 26 cm, which determines the axial wave length of the excited RF field to be $\lambda_{||} = 52$ cm. The loops are connected in series to an RF oscillator which can deliver pulsed RF power up to 200 W at a frequency of 0.35 MHz. The RF antenna A excites the linearly polarized RF field given by Eq.(5-1), which we call the $m = \pm 1$ field.

The RF antenna which produces circularly polarized RF field consists of ℓ half turn loops placed around the plasma column so that they forms an angle of $360^\circ/\ell$ with each other in azimuthal direction. In the experiments, we use both the $\ell = 3$ and $\ell = 4$ RF antennae. The RF antenna B in Fig.5-4 shows the case of $\ell = 3$, while the antenna (c) in Fig.2-7 is essentially the $\ell = 4$ type. Each loop has the same dimension as that used in the antenna A. The loops are connected in parallel to the RF oscillator through an LC-phase divider and impedance matching circuits. By adjusting the phase divider so that ℓ -phase RF current flows in the

antenna with proper sequence, we can establish the $m = +1$ or $m = -1$ RF field selectively.

4. EXPERIMENTAL RESULTS

A. Antenna Loading and Wave Excitation

The series loading resistance of the antenna, P_{in}/I_T^2 , was measured as a function of ω/ω_{ci} and n_0 , where P_{in} is the incident RF power into the RF circuits and the plasma as measured with a directional coupler connected at the output of the oscillator and

$$I_T^2 = \left\langle \frac{1}{2\pi} \int_0^{2\pi} (u I_\theta)^2 d\theta \right\rangle, \quad (5-4)$$

which gives total spectral power of azimuthal RF current I_θ in each u-turn loop. The axial RF current is of no consideration because it does not couple to the plasma as long as the electrostatic shield is installed. The resistance was resonantly peaked near $\omega/\omega_{ci} = 1$ for the $m = +1$ and the $m = \pm 1$ operations, and the increment of the series loading resistance at resonance relative to its vacuum value, ΔR , is plotted against n_0 in Fig. 5-4 for several modes. The value of ΔR increases linearly with n_0 . It should be noted that the $m = +1$ data points obtained by using different types of antennae ($\ell = 3$ and $\ell = 4$) fit on the same line. The ratio of resistive load for the $m = +1$ operation to that for the $m = -1$ operation is about 5. For the $m = \pm 1$ operation, the loading resistance should be the arithmetic mean of the $m = +1$ and $m = -1$ values because I_T^2 is equally divided into the $m = +1$

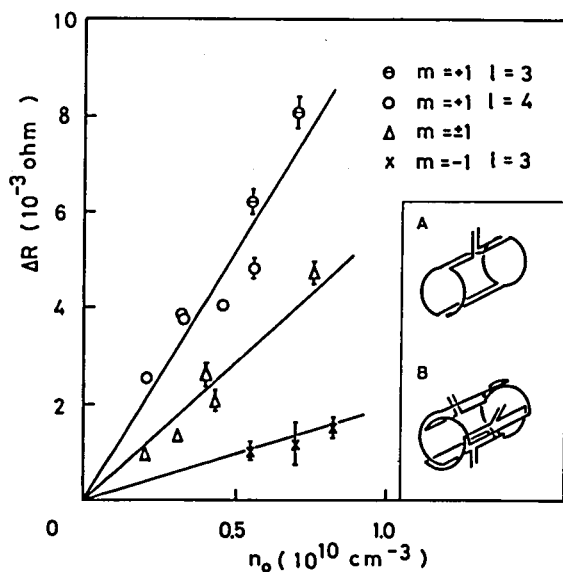


Fig.5-4 Increment of the equivalent series loading resistance vs. the density for several types of RF antennae. The inset shows simplified schematics of RF antennae.

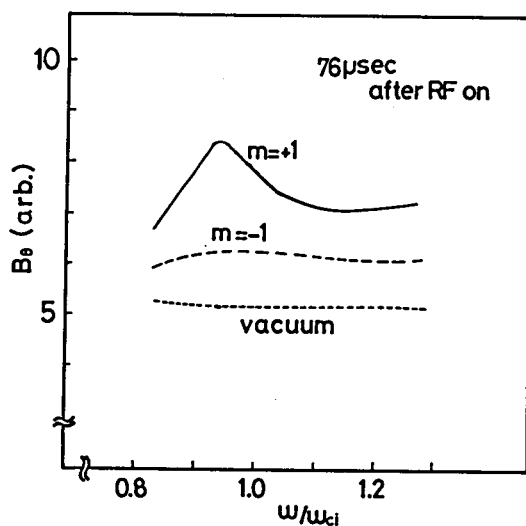


Fig.5-5 The amplitude of B_θ as a function of ω/ω_{ci} for the $m = +1$ (solid line), $m = -1$ (broken line), and $m = +1$ in vacuum (dotted line).

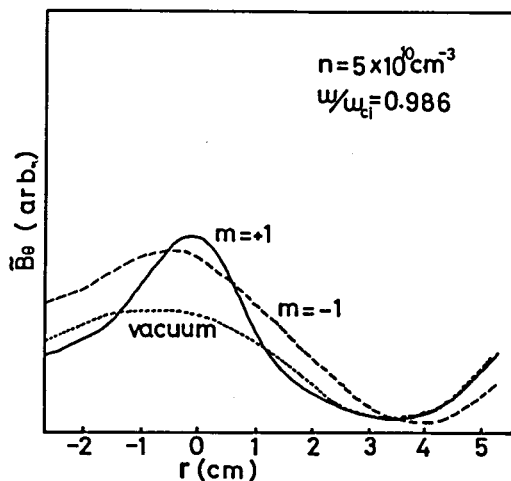


Fig.5-6 Radial profiles of $|B_\theta|$ for $m = +1$ and $m = -1$ modes.

and $m = -1$ spectral power. This is clearly seen from the result shown in Fig.5-4. We have thus confirmed the reasonableness of Eq.(5-1).

The RF field was measured by a magnetic probe located 44 cm downstream from the center of the RF antenna. The amplitude of the azimuthal component of the magnetic field is shown in Fig.5-5 as a function of ω/ω_{ci} for the $m = +1$ operation (solid line), the $m = -1$ operation (broken line), and the $m = +1$ operation with no plasma (dotted line). It is clearly seen that, for the $m = +1$ mode, the slow wave is resonantly excited near the ion cyclotron frequency ($\omega/\omega_{ci} = 0.96$ for $n_0 = 6 \times 10^{10} \text{ cm}^{-3}$). If the density was more decreased, the peak shifted closer to the point where $\omega/\omega_{ci} = 1$. When the RF antenna was operated in the $m = -1$ mode, the RF magnetic field amplitude had no dependence on ω/ω_{ci} . The measured phase difference between the radial and the azimuthal magnetic fields showed that the field was almost left-hand circularly polarized on axis. These results indicate that the $m = +1$ field generates the $m = +1$ slow wave, while the $m = -1$ field generates only antenna near field rather than a propagating wave field. For the antenna A, the $m = \pm 1$ field is equally divided into the $m = +1$ and -1 spectral powers, and both the $m = +1$ slow wave and the $m = -1$ antenna near field are generated.

Radial profiles of $|B_\theta|$ for $m = +1$ and -1 modes are compared in Fig.5-6. It is found that the $m = +1$ wave field is sharply peaked as compared with the $m = -1$ field. The axial components of the RF magnetic field was about 10% of the other components.

B. Measurement of Convective Loss

The temporal evolution of two-dimensional profile of ion saturation current I_{is} during rf pulse was measured with the Langmuir probe movable in three degrees of freedom. The equi- I_{is} contours are displayed in Fig. 5-7 for the $m = +1$ and ± 1 operations with mean squared rf current $I_T^2 = 0.51 \times 10^3 \text{ A}^2$, $\omega/\omega_{ci} = 0.99$, and $n_0 = 1.2 \times 10^{10} \text{ cm}^{-3}$ for both modes. The contours are almost concentric at $t=0$ (before the RF pulse) and continue to be so with time for the $m = +1$ operation. On the contrary, the profile for the $m = \pm 1$ operation shows a convective drift pattern at $t=0.4$ msec such that higher density, hot ions in the plasma core drift radially outward and lower density, cold ions in the periphery radially inward in adjacent quadrants alternatively. From the boundary condition which relates $B_z^{(1)}$ to J_θ , we note that the azimuthal position where $J_\theta = 0$ corresponds to $\theta=0$ or π in the theory described before. The arrows in Fig. 5-7 designate the direction of Γ_{ir} predicted from Eq. (5-3). There is a good agreement between the prediction and the experimental result.

We also note from Eq. (5-3) that the rotation of the RF antenna A in Fig. 5-4 by 90° in azimuth will result in the reversal of the direction of Γ_{ir} . This is demonstrated by the difference of the measured I_{is} -contours in Fig. 5-8. Only the antenna A was rotated by 90° between the two traces, all other parameters were the same. The convective drift patterns well corresponds with the theoretically predicted direction of Γ_{ir} which is again indicated by arrows. No enhanced cross-field loss was observed when the $m = -1$ field was applied, showing that both

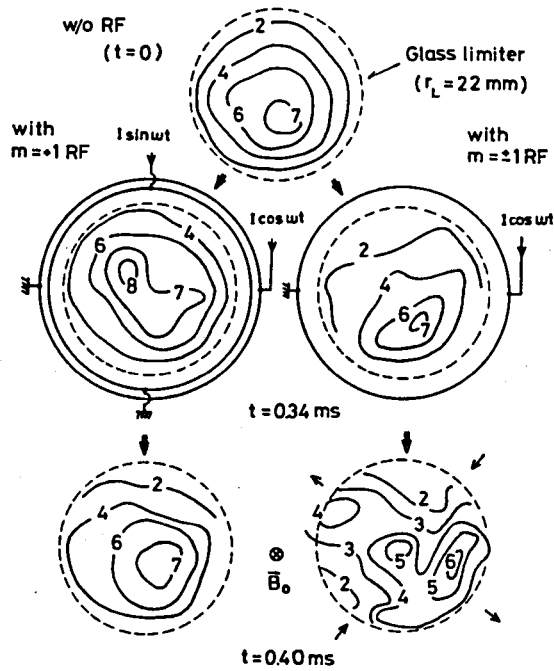


Fig.5-7 Two-dimensional profile of ion saturation current I_{is} (in arbitrary unit) for $n_0(r=0) = 1.2 \times 10^{10} \text{ cm}^{-3}$, $\omega/\omega_{ci} = 0.99$, and $I_T^2 = 0.51 \times 10^3 \text{ A}^2$. Temporal evolution of I_{is} profile for the $m = +1$ mode (left traces) and for the $m = \pm 1$ mode (right traces) is shown.

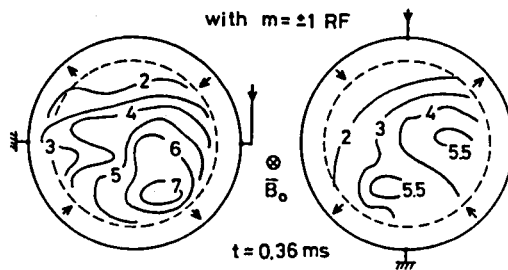


Fig.5-8 Two-dimensional I_{is} -profile for the $m = \pm 1$ mode with two different antenna positions.

E_1 and E_3 (or E_2 and E_4) are necessary for the convective motion.

The cross-field ion flux was obtained in such a way that the axial ion saturation current I_c to the negatively biased cold end plate of radius r_p was measured as a function of axial distance z to give Γ_{ir} which is equal to $(2\pi r_p e)^{-1} (\partial I_c / \partial z)$ [8]. Figure 5-9 shows Γ_{ir} divided by the density averaged over the cross section, $\bar{n}$, as a function of I_T^2 . The drift wave activities; an increase of the amplitude and the phase velocity and changes in the frequency spectrum, were observed as I_T^2 was increased. The dotted line in Fig. 5-9 is drawn theoretically using the Bohm diffusion coefficient with measured T_e and the density scale length for the $m = +1$ case. The numerical factor of the coefficient, so chosen as to fit the experimental data at $I_T^2 = 0$, is 5.2×10^{-2} . The theoretical curve agrees well with the experimental points for the $m = +1$ mode, showing that the increase of radial loss for this mode is due to the enhancement of Bohm diffusion. Since the drift wave activities were nearly the same for both modes with the same I_T^2 , the excess of $\Gamma_{ir}/\bar{n}$ for the $m = +1$ mode (solid curve) over that of the Bohm diffusion (dotted curve) is attributed to the radial ion convection peculiar to the $m = +1$ field. This fractional $\Gamma_{ir}/\bar{n}$ corresponding to the convective loss is proportional to I_T^2 , which is consistent with Eq. (5-3) since E_1 to E_4 are proportional to I_T . From the measurement of wave magnetic fields with probes, the magnitudes of E_1 and E_3 are estimated to be 6.3 V/m and 3.8 V/m, respectively, at $r = 2.0$ cm for $I_T^2 = 0.62 \times 10^3 \text{ A}^2$. Substituting these values and $\omega/\omega_{ci} = 0.98$ into Eq. (5-3), we obtain $\Gamma_{ir}/\bar{n} = 2\Gamma_{ir}/n_0(r=0)$ to be 6.7 m/sec with an error of $\pm 30\%$ involved in the estimation of the field amplitudes. The directly measured

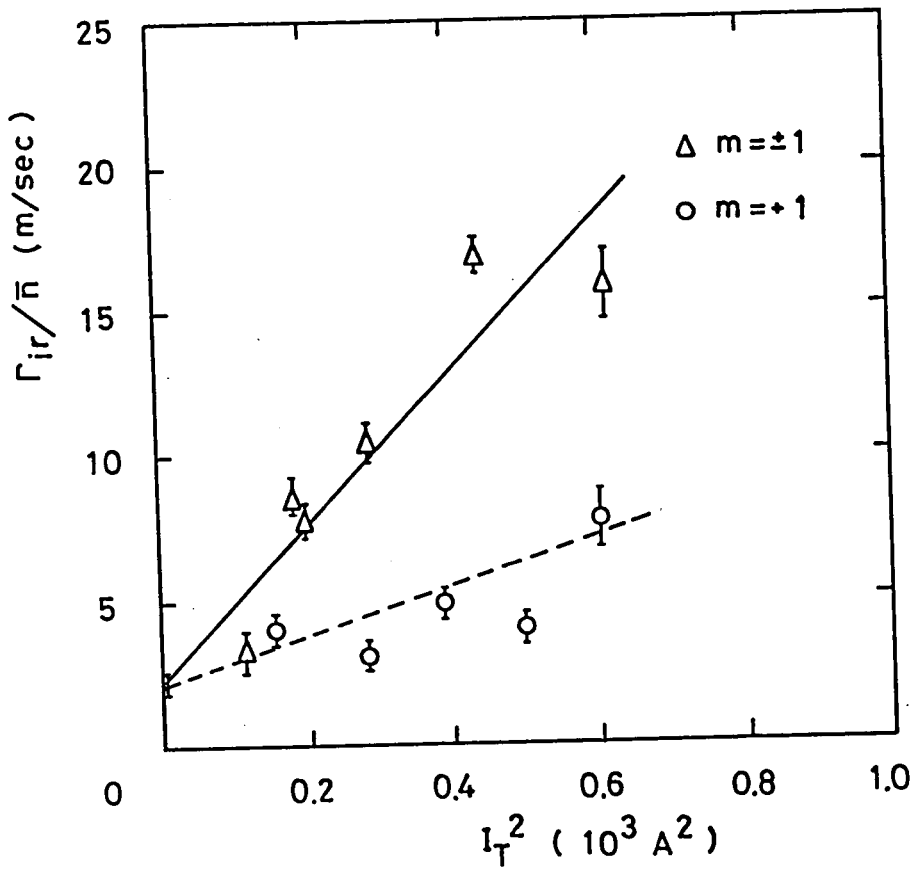


Fig.5-9 Radial ion flux Γ_{ir} divided by the averaged plasma density $\bar{n}$ versus I_T^2 for the $m = +1$ and $m = \pm 1$ modes. [$n_0 = (1-3) \times 10^{10} \text{ cm}^{-3}$ and $\omega/\omega_{ci} = 0.98$.]

value of $\Gamma_{ir}/\bar{n}$ for the $m = \pm 1$ mode above the Bohm diffusion value is (6.9-9.3) m/sec for the same I_T^2 as shown in Fig. 5-9. The agreement between theory and experiment is seen to be very good.

There may arise a question that the convective drifts are produced by the azimuthal asymmetry of equilibrium potential due to asymmetric heating of ions or electrons by the $m = \pm 1$ field. The difference in equilibrium potential of (0.6-0.8) V between two points separated by 90° in azimuth at $r \approx 1.5$ cm is necessary to produce $\Gamma_{ir} \approx 9 \times 10^{15} \text{ m}^{-2} \text{ sec}^{-1}$, which is measured at $\bar{n} = 0.5 \times 10^9 \text{ cm}^{-3}$, $n_0(r=1.5 \text{ cm}) = 0.3 \times 10^9 \text{ cm}^{-3}$, and $I_T^2 = 0.48 \times 10^3 \text{ A}^2$. However, the measured potential difference is always less than 0.3 V, which is much weaker than that required.

C. Comparison of Ion Heating

The perpendicular ion temperature $T_{i\perp}$ was measured with the multi-grid energy analyzer and the ion sensitive probe. In Fig. 5-10 are shown the temporal evolutions of $T_{i\perp}$, T_e , and n_0 at a position of 44 cm downstream from the RF antenna. The envelope of RF current in the antenna is also shown in the Figure. Significant difference in ion heating is observed between the $m = +1$ and $m = -1$ cases although the RF electric fields in the plasma have little difference in amplitude. This is the very confirmation of the theoretical prediction that ions are heated by the RF field which is left-hand circularly polarized. A little ion heating for the $m = -1$ mode may be due to the presence of the left-hand component in outer radii. For the $m = \pm 1$ mode, it is the fractional component of $m = +1$ mode that heats ions.

The electron heating may be due to the electron Landau damping of the slow wave (the thermal velocity of the electron is nearly equal to the phase velocity of the wave) or RF Joule heating. The density decrease for the $m = +1$ mode and for the $m = -1$ mode is attributed to the enhanced Bohm diffusion as discussed in the previous Subsection. The convective loss peculiar to the $m = \pm 1$ mode produces very rapid decrease of the density as seen from Fig.5-10.

The reciprocal of the density decay time, τ_n^{-1} , for the $m = \pm 1$ operation with $I_T^2 = 1.27 \times 10^3 \text{ A}^2$ was $9.6 \times 10^3 \text{ sec}^{-1}$, which is larger than the sum of $\tau_n^{-1} = 5.4 \times 10^3 \text{ sec}^{-1}$ for the $m = +1$ operation with $I_T^2 = 9.79 \times 10^2 \text{ A}^2$ and $\tau_n^{-1} = 3.05 \times 10^3 \text{ sec}^{-1}$ for the $m = -1$ operation with $I_T^2 = 3.33 \times 10^2 \text{ A}^2$. This result shows that the density decay for the $m = \pm 1$ operation is enhanced over the $m = +1$ case by an amount more than what is brought by the fractional $m = -1$ spectral power in the $m = \pm 1$ field, supporting the existence of nonlinear effect in the plasma loss mechanism, that is, convective plasma loss.

Figure 5-11 shows the increment of the ion temperature as a function of the static magnetic field. The RF current reached its maximum value of $I_T = 30 \text{ A}$ at 200 μsec after switching on the RF, and $T_{i\perp}$ was measured at 180 μsec . The ions are heated resonantly at a frequency near the ion cyclotron frequency for the $m = +1$ mode. The $\pm 10\%$ discrepancy in ω/ω_{ci} value yields the reduction of $T_{i\perp}$ by a factor of 3. For the $m = -1$ mode, ion heating is very little as expected.

Figure 5-12 shows the dependence of ion heating on the plasma density for $m = +1$ and $\omega/\omega_{ci} = 1.0$. The value of $T_{i\perp}$

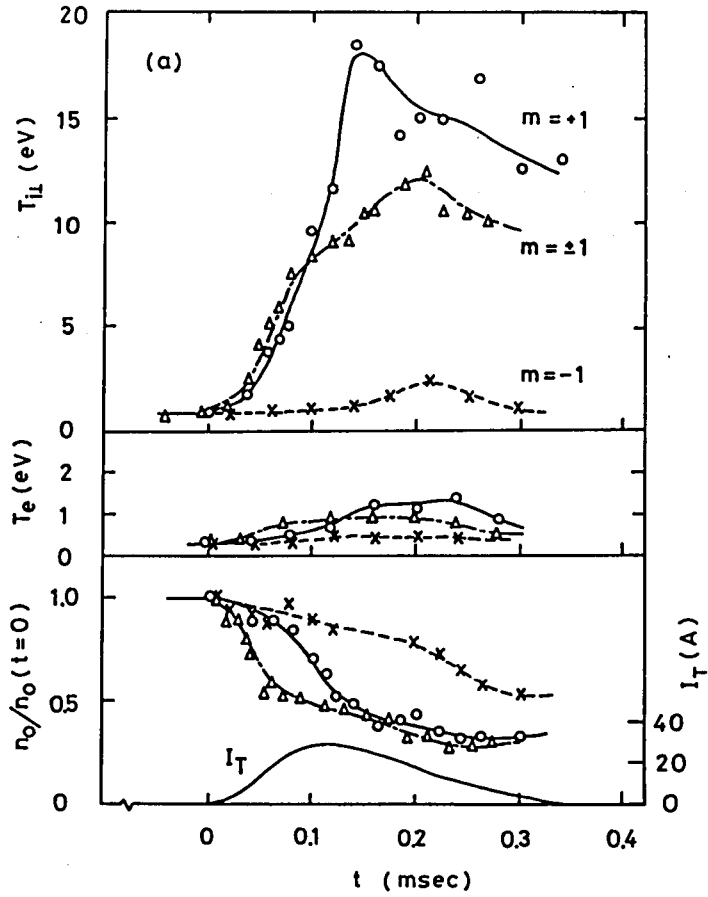


Fig.5-10 Temporal evolution of the perpendicular ion temperature, the electron temperature, and the plasma density for three RF modes.

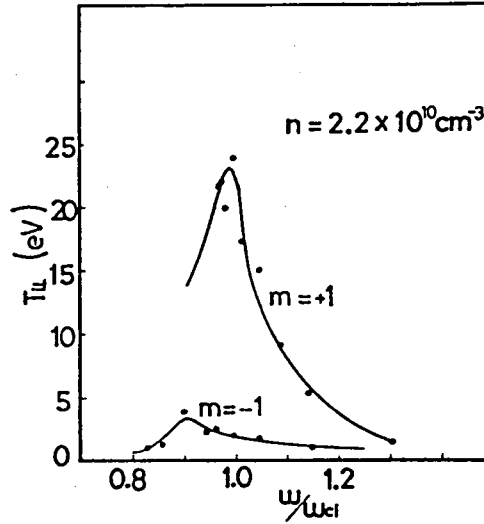


Fig.5-11 $T_{i\perp}$ versus ω/ω_{ci} for the $m = +1$ and ± 1 modes.

is plotted versus the initial plasma density at different times after RF is switched on. The RF current, which reaches 49 A at $t = 200 \mu\text{sec}$, is kept constant throughout the measurement. No obvious dependence of $T_{i\perp}$ on the initial plasma density is observed. This result implies that the RF power absorbed by ions per unit volume is proportional to the plasma density as expected from the hot plasma theory. The result is also consistent with the plasma loading resistance in Fig.5-4 where ΔR is proportional to the density. The value of $T_{i\perp}$ is maximum on axis and becomes markedly smaller for larger radii in accordance with the radial profile of the left-hand component of the $m = +1$ slow wave field. These facts indicate that the observed ion heating results from fundamental cyclotron damping of the slow wave.[9].

The heating efficiency for each RF mode is given in Fig. 5-13, where the ordinate is the maximum increment of ion energy density, Δp , and the abscissa is the incident RF energy into the plasma from $t = 0$ to the time when the maximum Δp is attained. The increased ion energy density is proportional to the incident RF energy. We can see that the $m = +1$ RF field indeed improves the overall heating efficiency over the $m = \pm 1$ RF field by more than 70 %. From the antenna loading measurement shown in Fig. 5-4, we found that, in the $m = \pm 1$ operation, about 20 % of the incident RF power is dissipated in the $m = -1$ field which heats ions little as shown in Fig.5-13. So, the elimination of the $m = -1$ field component should result in improvement in heating efficiency of no more than 20 %. The observed improvement far more than expected is due to the fact that the heating by the $m = \pm 1$ field is accompanied by the convective cross-field loss of accelerated ions.

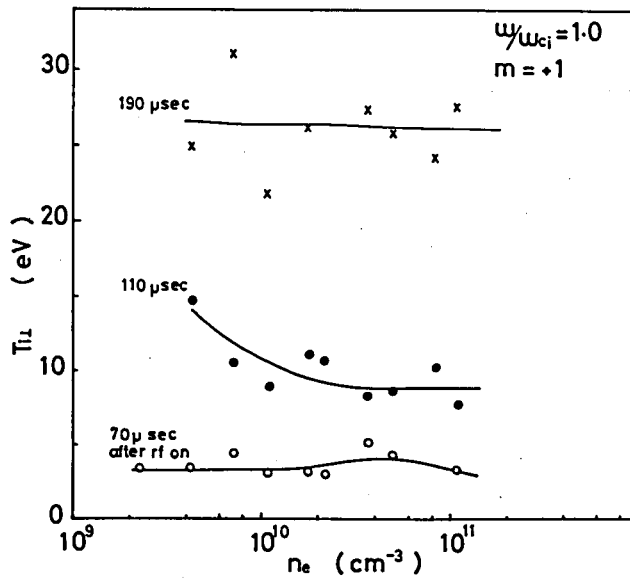


Fig.5-12 Dependence of the ion heating on the plasma density for the $m = +1$ mode.

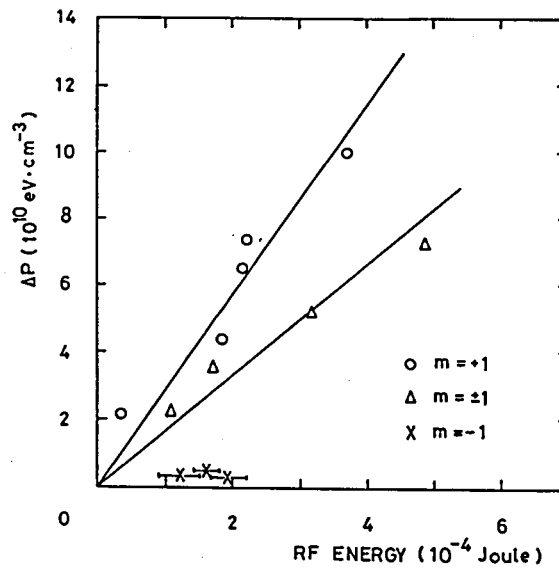


Fig.5-13 Maximum increment of ion energy density versus the incident RF energy for three RF modes. ($n_0 = 5.0 \times 10^{10} \text{ cm}^{-3}$ and $\omega/\omega_{ci} = 0.97$.)

5. CONCLUSION

We have investigated nonlinear effects in the Q-machine plasma subjected to large amplitude RF fields in ion cyclotron range of frequencies.

We have observed for the first time that the convective plasma loss occurs due to the azimuthal ponderomotive force of the $m = +1$ ICRF field. This result is in quantitative agreement with the theory based on the fluid model in which quasi-linear corrections are taken into account.

In order to eliminate the convective plasma loss, we have developed the RF antenna which can produce $m = +1$ circularly polarized RF field. By utilizing the $m = +1$ RF field, the convective loss is eliminated and the improvement in heating efficiency by more than 70 % is achieved.

Careful measurements of the wave propagation and characteristics of the ion heating have revealed that the power absorption is entirely due to the cyclotron damping of ion cyclotron waves.

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CHAPTER 6

APPLICATIONS OF IMPROVED ICRF HEATING TO TOROIDAL AND MIRROR PLASMAS

1. INTRODUCTION

ICRF heating is recognized to be one of the most promising methods for additional heating of a fusion plasma. A number of experiments performed in tokamaks have confirmed the effectiveness of ICRF heating with fast waves up to MW level. For open ended systems, power level of ICRF heating with fast or slow waves is yet within a few hundred of kW, and a great deal of efforts are devoted into the increase of input RF power and the improvement of heating efficiency.

In 1978, we first verified experimentally that the RF field of $m = +1$ mode was much more effective in slow wave heating than the $m = \pm 1$ mode employed conventionally [1]. The $m = +1$ RF field is purely left-hand circularly polarized on axis and can accelerate ions effectively, while for the $m = \pm 1$ field, a fractional RF power of the $m = -1$ component is lost without heating ions. Furthermore, we have observed in Chap. 5 that, for the $m = \pm 1$ mode, the nonlinear interaction between the $m = +1$ and -1 RF fields produces convective motion of ions in radial direction resulting in the reduction of radial confinement. The pure $m = +1$ RF field has no such problem, and can give the heating efficiency 1.7 times higher than the $m = \pm 1$ RF field [2].

In this Chapter, in order to clarify the advantage of the $m = +1$ RF field in more general, we investigate the ICRF heating by slow waves in two types of devices, i.e., a small axisymmetric torus and a simple mirror. ICRF heating of the toroidal plasma is performed in Sec. 2 using slow waves generated by $m = +1$ circularly polarized RF fields and $m = \pm 1$ linearly polarized RF fields. From direct comparison between two field-modes on wave excitation, ion heating, and effects on plasma confinement, it is revealed that the circularly polarized fields are far more effective than the linearly polarized fields in both achieving higher heating efficiency and reducing plasma loss. The ion heating by the linearly polarized fields is accompanied with enhanced cross-field plasma loss, which is again quantitatively explained by the two-fluid theory.

Section 3 contains the experimental results of ICRF heating in the simple mirror. It is also confirmed that the use of the $m = +1$ RF field improves the heating efficiency. Discussions and conclusions are given in Sec. 4.

2. TOROIDAL PLASMA HEATING

A. Experimental Apparatus

The experiment is performed in a small axisymmetric torus named Synchromak [3,4] shown in Fig.6-1. A hydrogen plasma of 4 cm in minor radius and 25 cm in major radius is produced by Joule heating in a toroidal field up to 6 KG. Typical toroidal current and loop voltage are 3 kA and 20 V, respective-

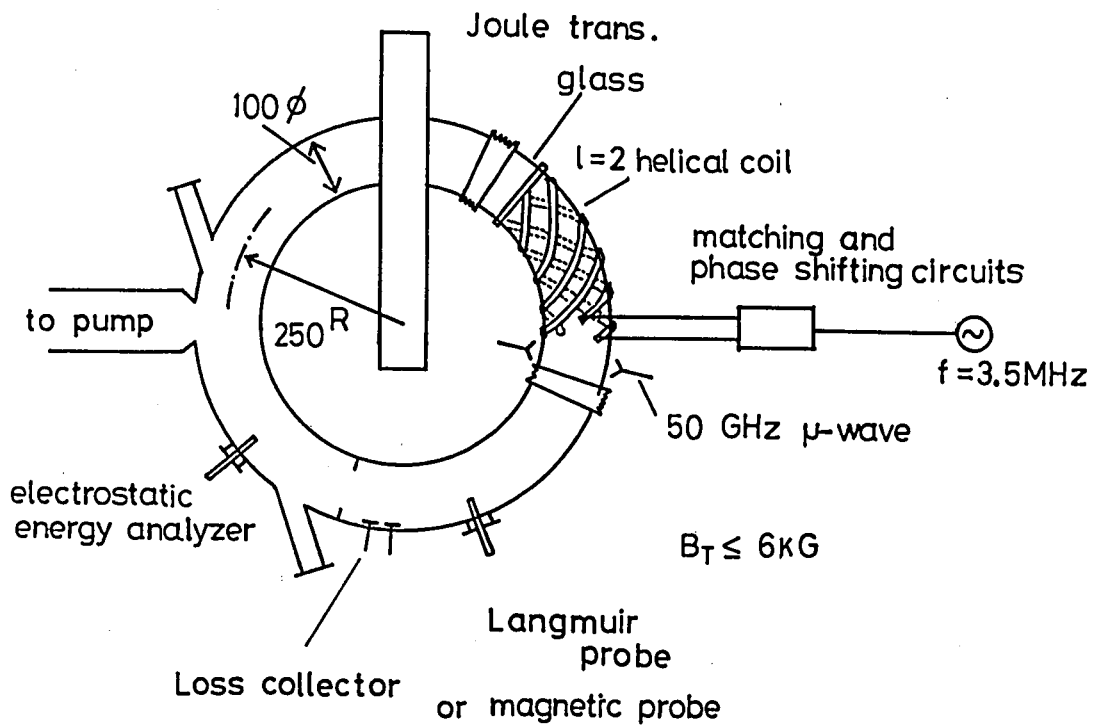


Fig.6-1 Schematic of the experimental apparatus.

ly. The plasma density as measured with a 6 mm microwave interferometer is initially $1 \times 10^{13} \text{ cm}^{-3}$ and decays at a time constant of about 0.5 msec. Ion and electron temperatures without RF are (12-14)eV and (20-25)eV, respectively.

ICRF field is produced with an $l = 2$ helical coil of one-field-period which consists of four helices wound at 90° intervals around the glass section of the torus with an electrostatic shield inside [5]. The pitch length of each helix is 0.3 m, hence, giving the axial wave number $k_{||}$ to be 21 m^{-1} . An RF oscillator which can deliver the maximum power of 200 kW at a frequency of 3.5 MHz is connected to the helical coil through a phase shifter and impedance matching circuits. Four helices are driven in sequential phase difference of 90° . We can thus establish the RF field of $m = +1$ mode (field rotation in the left-hand sense with respect to the toroidal magnetic field) or the $m = -1$ mode, selectively. These modes are circularly polarized on minor axis. The $m = \pm 1$ linearly polarized RF field, which is the superposition of the $m = +1$ and $m = -1$ fields, is obtained when two helices placed opposite to each other are driven out of phase, with the other two being in open-circuit. This implies that direct comparison between circularly and linearly polarized fields is possible using the same RF coil in the same machine.

B. Wave Excitation

The series loading resistance of the RF coil, R , is obtained by the relation; $R = P_{in} / I_T^2$, where P_{in} is the incident RF power and I_T^2 is the sum of the squared RF current in every

helix. The resistance without the plasma, R_{vac} , is about 1.2Ω , which corresponds to the circuit and wall losses. In Fig.6-2 is shown the series plasma loading resistance ΔR ; the difference between R with and without the plasma, as a function of ω/ω_{ci} , where ω_{ci} is the ion cyclotron frequency on minor axis. Data were taken at a density $n_0 = 6.0 \times 10^{12} \text{ cm}^{-3}$ and with P_{in} less than 3 kW. For the $m = +1$ and $m = \underline{+1}$ operating modes, ΔR has a peak at $\omega/\omega_{ci} \approx 0.95$, while, for the $m = -1$ mode, no resonant coupling is seen.

The wave field is detected by a small magnetic probe away from the RF coil. As the RF power is applied to the density-decaying plasma, the wave field on minor axis is peaked once at a certain time after switching on the RF power for a fixed ω/ω_{ci} . The plot of the density and ω/ω_{ci} at that time on the parameter space gives the dispersion relation. The result for the $m = +1$ operation is shown in Fig.6-3. We also calculate the dispersion relation of the slow wave in a cylindrical, bounded, cold plasma using the procedure described in Chap. 2. (The toroidal effect is not taken into account.) The solid curves in Fig.6-3 show the dispersion relation of the $m = +1$ slow wave with the radial $n = 0, 1$, and 2 modes for $k_{||} = 21 \text{ m}^{-1}$. The experimental points seem to agree semi-quantitatively with the calculated dispersion.

The radial distribution of the RF electric field for the present experimental condition is shown in Fig.2-5(a) in Chap.2 for $m = +1$, $n = 1$, and $\omega/\omega_{ci} = 0.96$. We compared the measured radial magnetic field distribution with the theoretical one and obtained reasonable agreement between them.

The $m = +1$ slow wave was also excited by the $m = \underline{+1}$

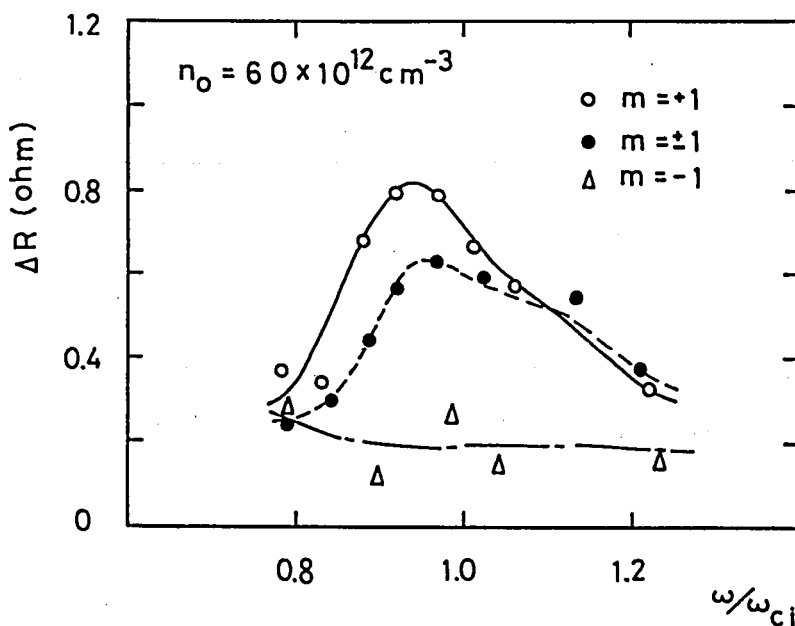


Fig.6-2 Increment of equivalent series loading resistance of the RF coil as a function of ω/ω_{ci} for $n_0 = 6.0 \times 10^{12} \text{ cm}^{-3}$. Open circles for the $m = +1$, closed circles for the $m = \pm 1$, and triangles for the $m = -1$ mode.

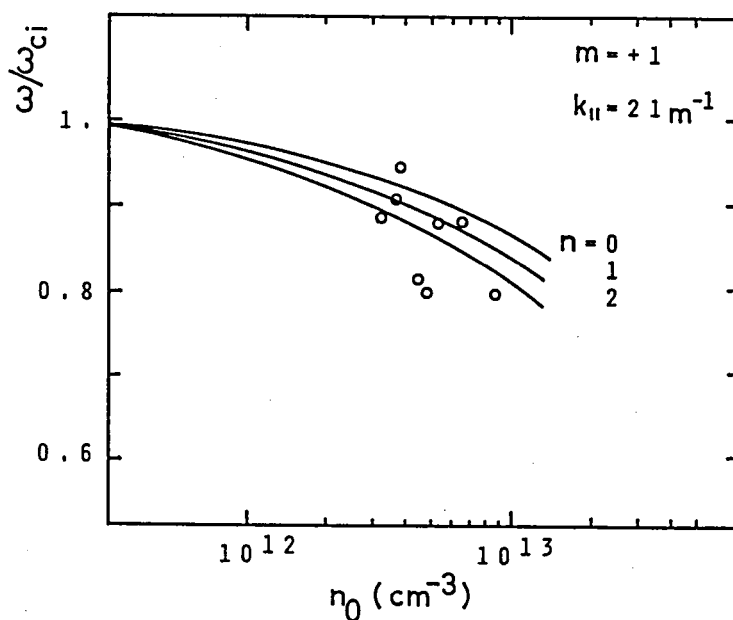


Fig.6-3 Dispersion relation of the excited wave. Solid curves are the calculated dispersion relation for the $m = +1$ slow waves with $n = 0, 1$, and 2 radial mode.

operation of the RF coil. As explained in Sec. 4 of Chap.2, the spectral power of the $m = \pm 1$ field would be equally divided into the $m = +1$ and -1 spectral powers, and it is the $m = +1$ component that generates the $m = +1$ slow wave. In this sense, the value of ΔR for the $m = \pm 1$ operation should be the arithmetic mean of those for the $m = +1$ and -1 operations. In Fig.6-2, this relationship is fulfilled in the region $\omega/\omega_{ci} < 1.0$ (propagating region).

We note that $(\Delta R)_{-1} = (0.2-0.25)(\Delta R)_{+1}$, where the subscript denotes the mode of operation. This means that, in the $m = \pm 1$ operation, (20-25)% of the RF power is lost as $m = -1$ field. We define η_w as the efficiency of generating left-circularly polarized RF fields, which accelerate ions. For the $m = +1$ operation, $(\eta_w)_{+1} = 1$, while, for the $m = \pm 1$ operation, $(\eta_w)_{\pm 1} = 0.8-0.83$.

C. Comparison of Ion Heating

A typical oscillogram of perpendicular ion temperature, $T_{i\perp}$, and density n_0 with and without RF are shown in Fig.6-4. Here, $T_{i\perp}$ is measured with a multi-grid energy analyzer of 6 mm in diameter which is inserted at $r = 1.5$ cm. The RF coil is operated in the $m = \pm 1$ mode, and the RF pulse of duration of 0.5 msec starts at $t = 0.24$ msec. As the plasma density decays from $1 \times 10^{13} \text{ cm}^{-3}$ to the half during RF pulse, the toroidal field is chosen to be $\omega/\omega_{ci} = 0.88$ so that optimum coupling is maintained.(see Fig.6-3) The temperature of bulk ions is raised from 12 eV to about 20 eV. The electron temperature without rf is (20-25) eV. No change in both the loop voltage and the toroidal current is observed

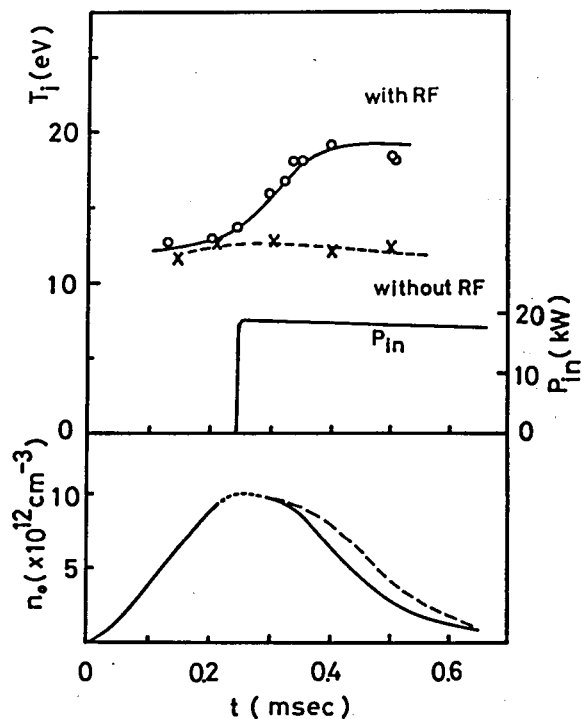


Fig.6-4 Temporal evolution of the plasma density and the perpendicular ion temperature with (solid line) and without RF (dotted line) of $m = +1$ mode. The incident RF power is also shown as a reference.

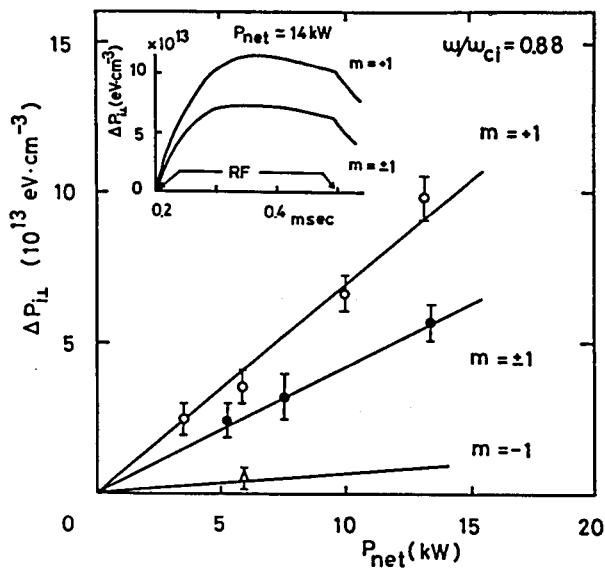


Fig.6-5 Increment of the perpendicular ion energy density as a function of the net RF power for three RF modes. The inset shows typical time evolution of the ion energy density.

when the RF pulse is applied, showing little electron heating, if any. Here, the effect of potential fluctuation caused by the RF pulse on probe characteristics of the analyzer must be considered. The fluctuation level in potential, ϕ , is several volts at a position of the measurement. Since $e\phi/kT < 1$ for both ions and electrons, an error involved in the determination of the ion and electron temperatures is expected to be less than a few % [6].

It is found from Fig.6-4 that the decay of the plasma density is enhanced by the RF pulse.

Figure 6-5 gives the increment of perpendicular ion energy density, $\Delta P_{i\perp}$, as a function of P_{net} for three operating modes, where $\Delta P_{i\perp}$ is defined as the difference of $n_0(t)T_{i\perp}(t)$ with RF from that without RF, and P_{net} is the net RF power into the plasma, which is given by $P_{in}\Delta R/(R_{vac} + \Delta R)$. An example of temporal evolution of $\Delta P_{i\perp}$ for $P_{net} \approx 14$ kW is shown in the inset. The value of $\Delta P_{i\perp}$ reaches a maximum within first (0.1-0.15) msec and remains approximately constant during RF pulse. We have used the value at $t = 0.5$ msec in the plot of Fig.6-5 in order to eliminate an error in $T_{i\perp}$ measurements caused by RF fields. The value of $\Delta P_{i\perp}$ increases linearly with P_{net} for the $m = +1$ and $m = \pm 1$ modes. When the RF coil is operated in the $m = -1$ mode, there is little ion heating. The $m = -1$ field is right-circularly polarized on minor axis and continues to be so for most plasma radii, so it does not accelerate ions. It is clearly seen from Fig. 6-5 that the energy gain of ions per unit RF power for the $m = +1$ mode is larger than for the $m = \pm 1$ mode; The value of $\Delta T_{i\perp} / P_{net}$ is 1.5 eV/kW for the

$m = +1$ mode and 0.9 eV/kW for the $m = \pm 1$ mode. In this case, $T_{i\parallel}$ is not measured, but in a separate experiment, it is observed that $T_{i\parallel} = (0.4-0.6) T_{i\perp}$. So, it may be reasonable to set $P_i = n_0 k (T_{i\perp} + \frac{1}{2} T_{i\parallel}) \approx \frac{5}{4} n_0 k T_{i\perp}$. Thus, we can say that the use of the $m = +1$ mode improves $\Delta P_i / P_{\text{net}}$ by a factor of 1.7.

D. Enhanced Plasma Loss

When the RF power is applied, the decay time of the plasma density appears to be shorter than that without RF. In Fig. 6-6 is plotted the reciprocal of the decay time, τ_n^{-1} , versus the applied RF power. The increase of τ_n^{-1} with P_{in} for the $m = \pm 1$ mode is significant compared with that for the $m = +1$ mode. We suspect that the confinement of the plasma is seriously influenced by the RF field of $m = \pm 1$ mode. In our plasma, the density is partially maintained by ionization of incoming neutral gas. So, the value of τ_n in Fig. 6-6 is different from the particle confinement time τ_p . In order to obtain τ_p , we measure ion flux toward the outer wall with plane collectors set just outside the radius of the limiter ($r_L = 4.0$ cm). The value of τ_p is calculated by the relation; $\tau_p = N/F$, where N is the total number of plasma particles within the radius of the limiter estimated from the microwave measurement, and F is the radial particle flux toward the whole wall estimated from the flux to the collectors. To do this we must assume toroidal and poloidal symmetry of ion flux and n_0 . The value of τ_p thus obtained is (0.22-0.29) msec when no RF is applied. Using the relation;

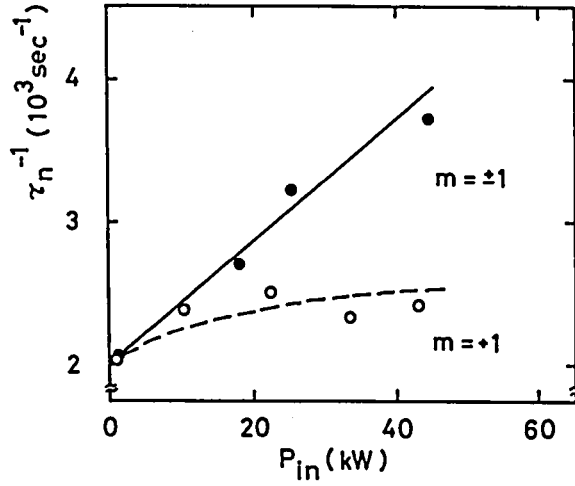


Fig.6-6 Reciprocal of the density-decay time versus the incident RF power for the $m = +1$ and $m = \pm 1$ modes.

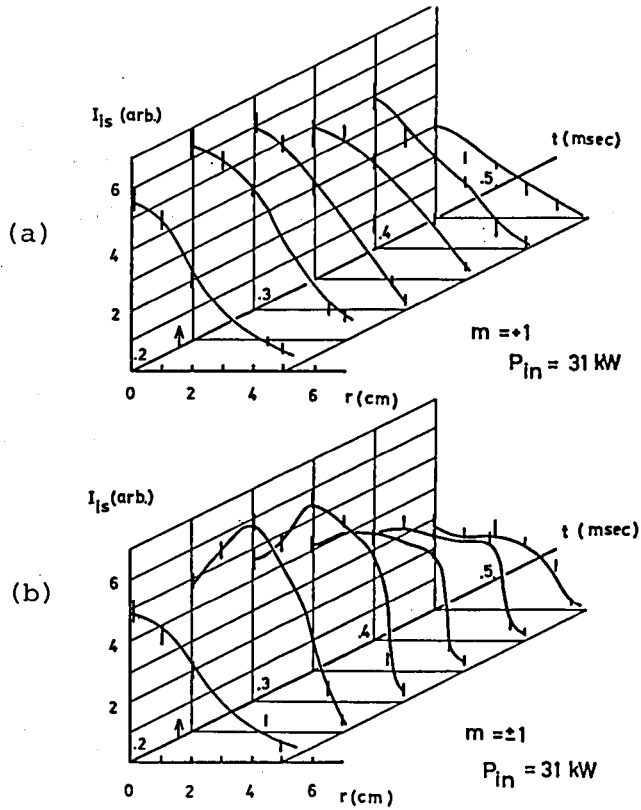


Fig.6-7 Radial profile of the ion saturation current for the $m = +1$ mode (a) and for the $m = \pm 1$ mode (b). The start of RF pulse is indicated by an arrow.

$$\frac{dn_0}{dt} = - \frac{n_0}{\tau_p} + \langle v \rangle n_0 ,$$

with $\langle v \rangle$ being the average ionization rate, we find $\langle v \rangle$ to be $(1.5-2.5) \times 10^3 \text{ sec}^{-1}$. The value of τ_p with RF is calculated by the relation; $\tau_p^{-1} = \tau_n^{-1} + \langle v \rangle$, under the assumption that $\langle v \rangle$ is unaffected by RF power. This assumption is confirmed to be true as discussed in the following. When the electrostatic shield beneath the RF coil was removed, the density increased by more than 30 % with $P_{in} \approx 20 \text{ kW}$. It is evident that the axial RF electric field enhances ionization. In the present experiment with the shielded coil, the density rather decreases with RF and no electron heating is observed. Hence, the effect of RF field on ionization is negligible. We finally obtain the particle confinement time with RF to be $(\tau_p)_{+1} = (0.20-0.25) \text{ msec}$ for the $m = +1$ mode and $(\tau_p)_{\pm 1} = (0.15-0.18) \text{ msec}$ for the $m = \pm 1$ mode. [$P_{in} \approx 45 \text{ kW}$ ($P_{net} \approx 14 \text{ kW}$) for both modes.]

Figure 6-7 shows temporal evolution of the radial profile of ion saturation current to a Langmuir probe for the $m = +1$ and $m = \pm 1$ modes. This curve roughly gives the radial density profile since T_e varies within 40 % or less over the minor radius. In the case of the $m = +1$ mode, the density decays holding an initial shape of the radial distribution after the start of RF. The density profile for the $m = \pm 1$ mode becomes much flatter with time indicating enhanced radial loss.

It has been shown that the application of the $m = \pm 1$ ICRF field results in decrease of particle confinement time

accompanied with the modification of density profile. It should also be noted that such a reduction in τ_p is little when the $m = +1$ field is launched. As the excursion velocity of ions in RF fields exceeds the thermal velocity in the present experiment, we must take into account ponderomotive effects.

Let us consider a cylindrical, homogeneous, collisionless plasma with radius r_L immersed in a static magnetic field B_0 in z direction. If the azimuthal current density in the helical RF coil in $m = \pm 1$ operation is approximated by

$$j_\theta = j_0 \delta(r - s) \cos(\theta - k_{\parallel} z) \cos \omega t \quad (6-1)$$

with s being the coil radius, the RF electric field in the plasma, which is a superposition of the $m = +1$ and -1 modes, may be represented by

$$\begin{aligned} E_r &= \frac{E_+}{2} [AJ_0 - BJ_2] \cos \phi_1 + \frac{E_-}{2} [-BJ_0 + AJ_2] \cos \phi_2 \\ E_\theta &= \frac{E_+}{2} [-AJ_0 - BJ_2] \sin \phi_1 + \frac{E_-}{2} [-BJ_0 - AJ_2] \sin \phi_2 \\ E_z &= 0, \end{aligned} \quad (6-2)$$

where $\phi_1 = \theta - k_{\parallel} z + \omega t$, $\phi_2 = -\theta + k_{\parallel} z + \omega t$, E_+ and E_- are, respectively, $m = +1$ and $m = -1$ field amplitudes which are proportional to the coil current, J_v is the v -th order Bessel function of the first kind with the argument $k_{\perp} r$, $A = k_{\perp}^2 / (-k_{\parallel}^2 + k_0^2 L)$, and $B = -k_{\perp}^2 / (-k_{\parallel}^2 + k_0^2 R)$.

Following the procedure used in Chap.5, we obtain the radial component of the quasi-steady ion flux, which is given by

$$\Gamma_{ir} = n_0 \left(\frac{e}{m_i} \right)^2 E_+ E_- \frac{(\omega^2 + \omega_{ci}^2)}{2\omega\omega_{ci}} \left[\frac{A^2}{(\omega - \omega_{ci})^2} - \frac{B^2}{(\omega + \omega_{ci})^2} \right] \frac{1}{r} J_0(k_{\perp} r) J_2(k_{\perp} r) \sin 2(\theta - k_{\parallel} z), \quad (6-3)$$

for the $m = \pm 1$ RF field. For the case of the $m = +1(-1)$ RF field in which $E_-(E_+)$ is identically zero, Eq.(6-3) predicts that $\Gamma_{ir} = 0$. This equation shows that a radial ion flux is produced by the interaction of field components of the same polarization but different propagation direction with respect to one another. The net loss flux of ions is obtained by integrating Eq.(6-3) over the whole surface at $r = r_L$. Since $n_0 = 0$ for $r > r_L$, the integration should be taken over that half-wave of the sine which leads to flux towards the outside.

The reciprocal of characteristic ion loss time τ_L^{-1} , which is given by $\Gamma_{ir}/n_0 r_L$, is proportional to the field amplitudes E_+ and E_- and becomes very large at ion cyclotron resonance. This prediction is consistent with the experimental result shown in Fig.6-6, in which the reciprocal of the density decay time for the $m = \pm 1$ mode is proportional to the incident RF power, i.e., to the square of the field amplitude. To calculate Γ_{ir} in Eq.(6-3), we measure wave fields with a calibrated magnetic probe. When $P_{in} \approx 45$ kW, $E_+ A$ and $E_- B$ are estimated to be $28 V_{pp}/cm$ and $11 V_{pp}/cm$, respectively. Substituting plasma parameters into Eq.(6-3) with $A/B \approx 1$, we obtain $\Gamma_{ir}(r = r_L)/n_0 \approx 3.1 \times 10^2$ m/sec. The characteristic

loss time is, therefore, 0.13 msec. On the other hand, a reduction in τ_p , from (0.22-0.29) msec to (0.15-0.18) msec, is observed when a 45-kW, $m = +1$ -mode RF pulse is applied. This requires that $\tau_L = (0.3-0.4)$ msec. The agreement between the predicted τ_L from Eq.(6-3) and the experimental value is fairly good. The prediction may be an underestimate of τ_L because the convective loss is considered to be localized near the RF coil.

The fact that τ_p is affected little by the $m = +1$ RF field is also explained by Eq.(6-3). We can thus say that the enhanced loss observed when the $m = +1$ RF field is applied is due to the ion convection caused by ponderomotive effects.

E. Power Balance

We consider power balance of ions by using the following equation:

$$\frac{dW_i}{dt} = \eta_w \eta_h P_{\text{net}} - \frac{W_i}{\tau_{Ei}} - \frac{W_i - W_e}{\tau_{eq}}. \quad (6-4)$$

Here, W_i and W_e are, respectively, ion and electron energy, τ_{Ei} is the ion energy confinement time, τ_{eq} is the equipartition time, and η_h is the efficiency of power transfer from left-circularly polarized RF fields to ions. (η_w is defined in Subsection 2B. The value of τ_{Ei} can be estimated from $\tau_{Ei}^{-1} = \tau_p^{-1} + \tau_{cx}^{-1} + \tau_{\text{cond}}^{-1}$, where τ_{cx} and τ_{cond} are loss times of energy due to charge exchange and thermal conduction, respectively. Under the experimental condition shown in the inset in Fig.6-5, we obtain $\tau_{cx} \approx 0.2$ msec, hence, $(\tau_{Ei})_{+1} \approx$

110 μsec and $(\tau_{\text{Ei}})_{+1} \approx 85 \mu\text{sec}$. [τ_{Ei} without RF is (110-160) μsec .] The last term on the right hand side of Eq.(6-4) is negligible since τ_{eq} is an order-of-magnitude larger than τ_{Ei} . From Eq.(6-5), we find that, at equilibrium, $\Delta W_i \equiv W_i - W_{i0} = \eta_w \eta_h P_{\text{net}} \tau_{\text{Ei}}$, where W_{i0} is the value of W_i without RF. Recalling that $(\eta_w)_{+1} = 1$ and $(\eta_w)_{-1} = 0.8-0.83$, and noting that η_h should be the same for both modes, we can predict that $\Delta W_i / P_{\text{net}}$ for the $m = +1$ mode is larger than for the $m = -1$ mode by a factor of 1.6. Direct measurement of ΔT_{i+} exhibits that the improvement factor of $\Delta P_i / P_{\text{net}}$ is 1.7. The agreement between the prediction and the experiment is very good. The value of η_h is calculated to be 23 % under the assumption that the radial profile of P_i is parabolic and that the heating is uniform along the torus.

3. MIRROR PLASMA HEATING

A. Experimental Apparatus

RF heating of a mirror plasma at $\omega \lesssim \omega_{\text{cia}}$, where ω_{cia} is the ion cyclotron frequency at an RF antenna position and ω is the applied RF frequency, was carried out in the axisymmetric mirror TPD-III. The magnetic field could be varied up to 1.3 T at the throat with the mirror ratio of 2-3. Hydrogen plasmas were injected through one mirror throat by the TPD plasma gun. The initial plasma parameters were as follows: $n_0 = (3-8) \times 10^{11} \text{cm}^{-3}$, $T_e = (10-15) \text{eV}$, and $T_i = 10 \text{eV}$. In Fig.6-8, the schematic diagram of TPD-III is shown

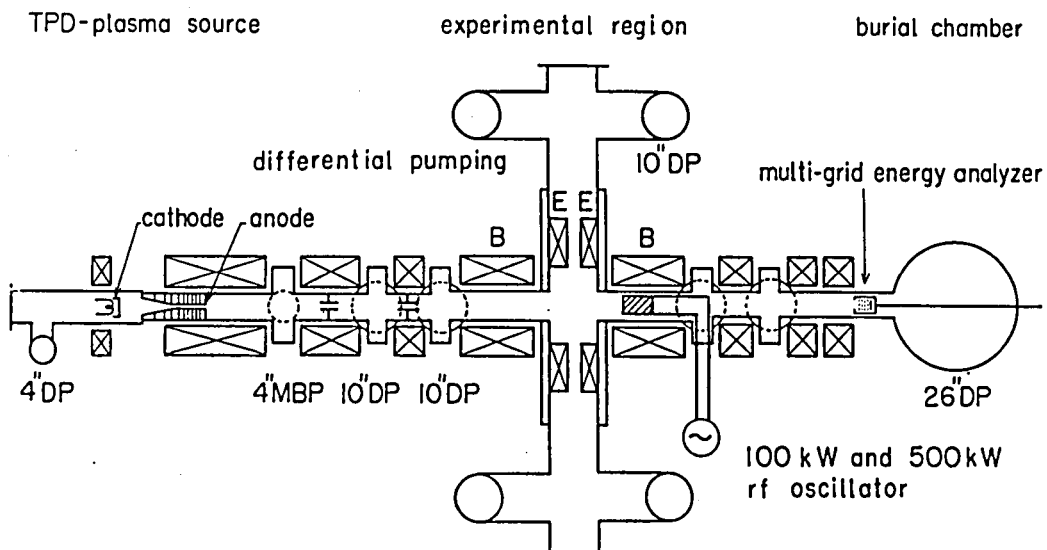


Fig.6-8 Schematic view of TPD-III.

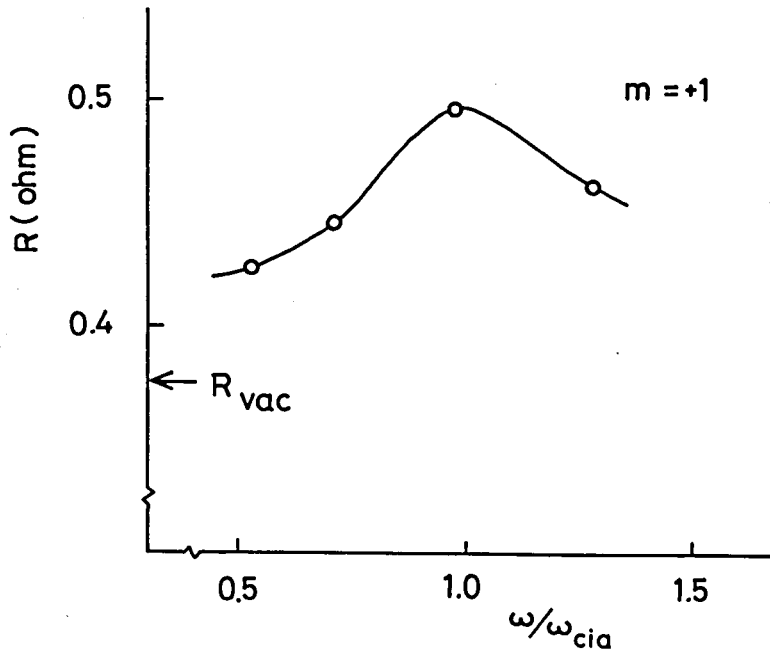


Fig.6-9 Series loading resistance of the RF antenna versus ω/ω_{cia} for the $m = +1$ mode.

together with RF antenna and some diagnostics. It is operated in quasi-steady mode of 30 msec duration and at 5 sec intervals. The pressure of the neutral particles in the experimental region is kept less than 1×10^{-5} Torr.

The RF antenna similar to the one shown in Fig.2-7(c) in Chap. 2 is installed at 27 cm downstream from the midplane. The length and the radius of the antenna are 6 cm and 3 cm, respectively. Beneath the antenna, copper straps of 0.5 cm width and 20 cm length are aligned with 1 cm interval along azimuthal direction. These straps work as a Faraday shield. Feeding RF currents 90° out of phase into two coil units of the antenna, we can establish the azimuthally rotating RF field either of $m = +1$ mode or of $m = -1$ mode. The in-phase excitation of two units gives the $m = \pm 1$ RF field. The antenna is excited by a self-oscillator with output power up to 500 kW at a frequency of 6.35 MHz. The resistance of the antenna with no plasma present was 0.38Ω .

B. Heating Results

After the plasma injection from the TPD gun, the quasi-steady state of plasma density lasts for about 30 msec. The MHD stability of the plasma is maintained through a line-tying effect of the metal plate of a Faraday cup which is located outside the right throat. The heating RF is applied for (0.5-1) msec during this period. Figure 6-9 shows the series loading resistance of the RF antenna as a function of ω/ω_{cia} for $m = +1$ mode. The resistance is sharply peaked

near the fundamental ion cyclotron resonance. We measured the polarization of the excited RF field by a magnetic probe to find that the field was almost left-hand circularly polarized on axis when the $m = +1$ field was generated.

Unlike the case of the toroidal plasma, the slow wave launched from the RF antenna propagates toward the midplane and cyclotron-damps at the beach. Ion temperature was measured at the midplane by using a Faraday cup, which is inserted perpendicularly with respect to the static magnetic field. Figure 6-10 shows the increment of $T_{i\perp}$ just after the RF pulse of 0.54 msec, $\Delta T_{i\perp}$, as a function of the net RF power P_{net} for the $m = +1$ and $m = \pm 1$ modes. The value of P_{net} was almost constant during the RF pulse. The data for $m = -1$ mode is not shown here because the increment of $T_{i\perp}$ is very little for this mode. The value of $\Delta T_{i\perp}$ increases linearly with P_{net} for both the $m = +1$ and ± 1 modes. With 0.7 kW of RF of $m = +1$ mode, the ion temperature doubled, from 15 eV to about 33 eV. In the toroidal plasma heating, it was confirmed that the increment of $T_{i\perp}$ per unit RF power for the $m = +1$ mode was 1.7 times greater than that for the $m = \pm 1$ mode. This result is reproduced in the mirror plasma heating as shown in Fig. 6-10. It is, however, seen that the improvement factor of the heating efficiency is 1.5, which is smaller than in the toroidal plasma.

We have observed in the previous Section that the convective loss caused by the $m = \pm 1$ RF field yields a flattening of the radial density profile. We also measured the density profile in TPD-III. Figure 6-11 displays the ion saturation

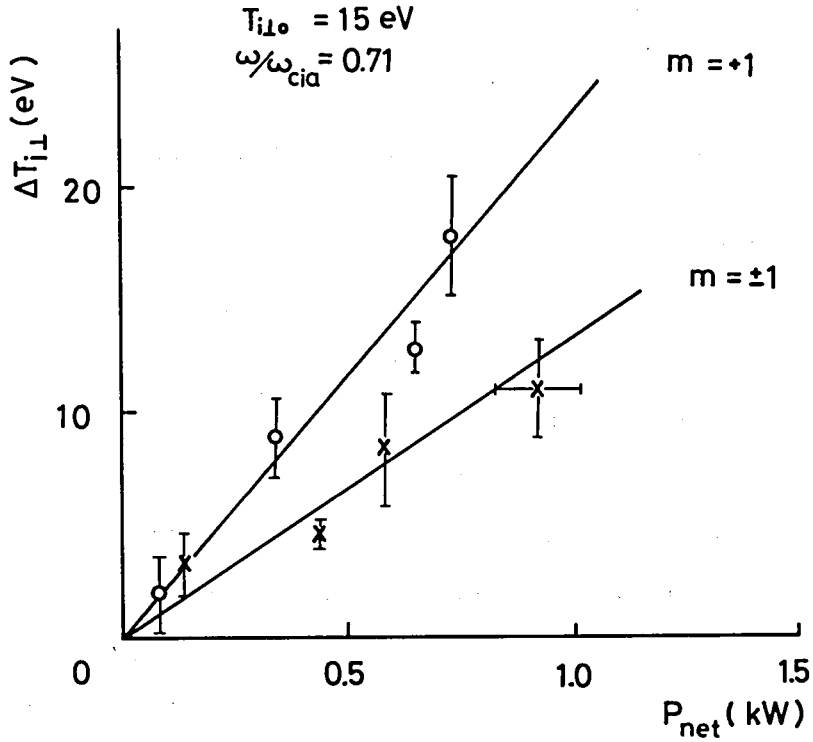


Fig.6-10 Increment of $T_{i\perp}$ as a function of the net RF power for the $m = +1$ and $m = \pm 1$ mode.

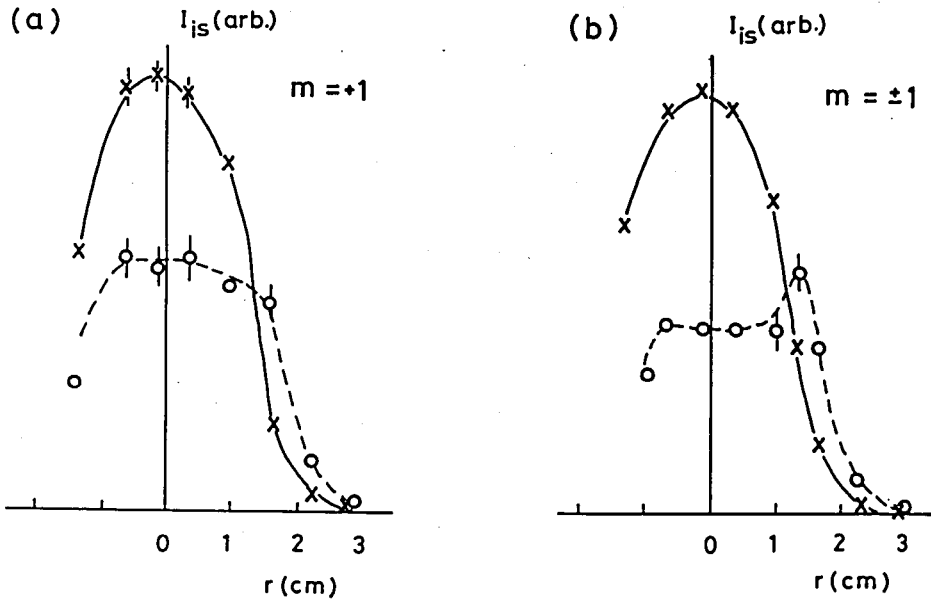


Fig.6-11 Radial profile of the ion saturation current with (dotted line) and without RF (solid line) for the $m = +1$ (a) and $m = \pm 1$ mode (b).

current of a single probe I_{is} located at the midplane versus the radius for $n_0 = 1.8 \times 10^{11} \text{ cm}^{-3}$, $\omega/\omega_{cia} = 0.71$, and $P_{net} = 0.5 \text{ kW}$. The solid line in the figure denotes the case without RF, while the dotted line shows I_{is} at 0.3 msec after switching on the RF pulse. We note that the density decreases significantly by the application of RF pulse even for the $m = +1$ mode. The measurement of the injected ion flux from the TPD gun with a directional probe at the midplane showed that the amount of plasma influx itself decreased and broadened during RF regardless of the RF mode. So, the observed decrease in density in Fig.6-11 may come from the change of plasma influx from the gun as well as the enhancement of diffusion due to RF heating [7]. Comparing the I_{is} -profile for the $m = \pm 1$ mode with that for the $m = +1$ mode, we find larger decrease of I_{is} on axis and more remarkable distortion of the profile especially near the periphery. These characteristics indicate the existence of the convective radial loss for the $m = \pm 1$ mode.

The profile of Fig.6-11(b) was measured on a plane 45° offset from the polarization plane of the RF electric field. The theory described in Chap. 5 predicts that the convection leads the flux toward outside on that plane. This may be in qualitative agreement with the measurement. In a plasma produced by a TPD source, a deep potential well tends to be formed within a radius of the anode hole. The resultant steady radial electric field gives rise to a plasma rotation, and ions would feel the ponderomotive force averaged over θ . If the period of $E \times B$ rotation is much shorter than the time for ions to traverse convectively the plasma radius, the averaged ponderomotive force

is zero, and hence no convection occurs. For a moderate $E \times B$ rotation frequency, there would appear a flattened or hollow density profile instead of the convective pattern as observed in HIEI-Q. This is the case of Fig.6-11(b). The convective loss for this case is considered to be smaller than expected from Eq.(5-3), and this may be a reason for the smaller difference of heating efficiencies of $m = +1$ and $\underline{+1}$ modes observed in Fig.6-10.

4. DISCUSSION AND CONCLUSION

In Sec. 2, ICRF heating at $\omega \lesssim \omega_{ci}$ has been investigated in Synchromak using the helical RF coil which can generate both the $m = +1$ circularly polarized field and the $m = \underline{+1}$ linearly polarized field.

It is found from the comparative measurements that the $m = +1$ RF field greatly improves the overall efficiency of ion heating over the $m = \underline{+1}$ RF field used conventionally.

The ion heating by the $m = \underline{+1}$ RF field is accompanied with the enhanced cross-field plasma loss. For an explanation of this result, is presented the theoretical prediction that the radial ion convection is caused by the ponderomotive force of the applied $m = \underline{+1}$ RF field. The prediction agrees quantitatively with the experimental result.

Of course, the slow wave heating used here is not applicable to larger tokamaks. It should be noted, however, that if poloidally asymmetric RF fields of moderate amplitude (a

few hundred volts / cm for present-day tokamaks) exist, resulting ponderomotive force might give rise to convective loss of surface plasmas, at least, near an RF coil.

In fast wave heating of a two-ion plasma, radial loss of the minority species would take place, resulting in increase of recycling or decrease of energy confinement time.

In Sec. 3, we have applied the improved ICRF heating to a mirror plasma produced in TPD-III.

Although the anomalous convective loss caused by the $m = \pm 1$ RF field was not clearly observed as in HIEI-Q and Synchromak, we again confirmed that the $m = +1$ RF field could improve the heating efficiency by a factor of 1.5. The smaller improvement factor as compared with other two experiments may be attributed to the much larger $E \times B$ rotation frequency in TPD-III.

Through the experiments described in Chaps. 5 and 6, we have established the superiority of the $m = +1$ rotating RF field over the $m = \pm 1$ RF field used conventionally.

Recently, Adati et al.[8] performed ICRF heating experiments in the double-cusp machine RFC-XX following in the wake of our experiments. They changed the polarization of the excited RF field at $\omega \lesssim \omega_{ci}$, from linear to circular continuously, by varying the phase of RF currents, ϕ , in two pairs of the Type-III coils located at the central cell. The increment of plasma energy measured from the diamagnetic signal is plotted in Fig. 6-12 as a function of ϕ for a constant RF coil current and $n_0 \approx 10^{12} \text{ cm}^{-3}$. In the abscissa, the point where $\phi = -\pi/2$ ($+\pi/2$) corresponds to the RF field of $m = +1$ (-1) mode. Their result shows that the energy increment for the $m = +1$ mode is

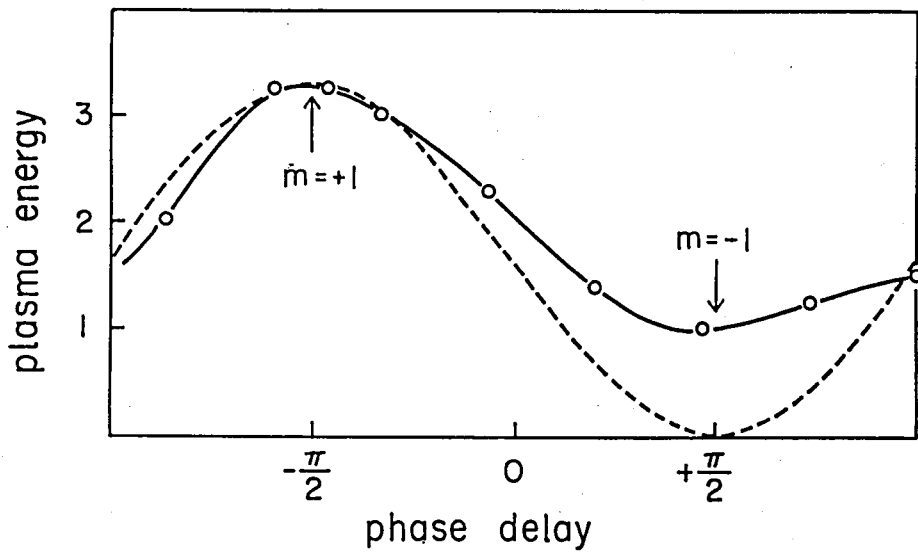


Fig.6-12 Increment of plasma energy as a function of the phase difference ϕ in RFC-XX. [After Adati et al. (Ref. 8)]

1.6 times higher than for the $m = \pm 1$ mode ($\varphi = 0$ or π), representing the exact reproduction of our preceding results.

In the tandem mirror Phaedrus is obtained an experimental result in which the convective drift due to the azimuthal ponderomotive force may play an important role. Hershkowitz et al.[9] observed that the radial potential profile was hollow when the plasma was RF-sustained and heated by the RF field launched from half-turn antennae. As the RF field includes the $m = \pm 1$ component, the resultant convective drift of hot ions toward outside may possibly form larger potentials in larger radii. They started to investigate the mechanism of the potential formation along this interpretation [10].

We finally conclude that the superiority of the $m = \pm 1$ RF field in slow wave heating observed in HIEI-Q for the first time has been generally confirmed by the experiments described in this Chapter.

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CHAPTER 7

PLUGGING OF OPEN ENDS IN LINEAR DEVICES BY ROTATING ION CYCLOTRON RF FIELDS

1. INTRODUCTION

RF plugging is one of the methods for enhancing the energy amplification factor of an open ended system. A number of experimental and theoretical studies have been carried out especially in IPP, Nagoya [1]. In the experiment performed by Watari et al.[2], it was shown that the $m = \pm 1$ RF field produced by the type-III coil was much more effective for end loss reduction than the $m = 0$ RF field in the sense that the required RF electric field strength remained relatively small for higher densities. They also explained that the plugging effect was due to the axial ponderomotive force for ions, and electrons were confined by a resultant ambipolar potential.

Fader et al.[3] concluded that the ponderomotive effect on electrons was dominant in their plugging experiment where the cross-sectional variation of end loss reduction was compared with the $m = \pm 1$ RF field distribution of the type-III coil to find that the maximum reduction occurred along the magnetic field line on which the axial electric field E_z was large.

We have shown in Chaps. 5 and 6 that the $m = \pm 1$ RF field

is far more effective in ICRF heating than the $m = \pm 1$ RF field. Since the perpendicular acceleration of ions by RF fields is responsible also for RF plugging in which the ponderomotive force for ions is dominant, it is predicted that a use of the $m = \pm 1$ RF field will enhance the end loss reduction.

In this Chapter, it is reported that the plugging effect is greatly improved when the $m = \pm 1$ RF field (rotating in the sense of ion gyration) is used instead of the $m = \pm 1$ RF field generated by the type-III coil. This result suggests that the plugging is attributed to the ponderomotive effect on ions because the left-hand polarized component of RF field of $m = \pm 1$ mode is larger than that of $m = \pm 1$ mode, while E_z remains the same.

We further present the detailed experimental investigation of effects of RF mode on the plugging and resultant plasma confinement in an axisymmetric mirror stabilized by line-tying. The particle and energy confinement times with and without the RF plugging are compared for different RF modes. The best confinement occurs when the $m = \pm 1$ mode is used. The confinement time is observed to be limited by an increase of the radial loss of the plasma due to associated heating and MHD activities.

If the line-tying is removed, the flute instability is observed resulting in a significant reduction of the confinement time. In this case, both the RF stabilization and the RF plugging are necessary for confinement. We also present the first experimental result on simultaneous application of the RF stabilization and the RF plugging. It is observed

a combined effect that the plugging efficiency is enhanced with less total RF power as compared with the case of RF plugging only.

In the latter half of this Chapter is described the application of the plugging by the $m = +1$ RF field to a point cusp of RFC-XX. RFC-XX is the open-ended machine in IPP, Nagoya, which consists of a central section and two cusp regions in both ends [4]. By applying the RF voltage to parallel electrodes located at two line cusps, the end loss flux through the line cusps is reduced to less than a few % of the initial value. The energy confinement time of the central-section plasma is improved by a factor of 2 through the line cusp plugging [5]. The end loss flux through the point cusps is about one-tenth of that through the line cusps, and hence up to now, has minor effects on total plasma confinement. However, once almost complete end loss reduction is achieved at the line cusps, an efficient RF plugging of the point cusps is necessary for further improvement of the energy confinement time. The type-III coil is adopted for this purpose so far, but we apply the $m = +1$ RF field to the point cusp plugging. The $m = +1$ RF field is confirmed again to have better plugging effect than the $m = \pm 1$ RF field.

The RF plugging of ICRF-heated plasma is also performed. The streaming plasma with T_i up to 250 eV can be plugged by the $m = +1$ RF field with an RF voltage less than 20 kV. A scaling formula of the RF plugging of the ICRF-heated plasma is obtained experimentally for the first time.

2. BRIEF THEORY OF RF PLUGGING

The RF plugging scheme has been investigated by many authors [3,6-8]. If a ponderomotive quasi-potential ψ_i acts on ions, there would arise an ambipolar potential ϕ that maintains plasma quasi-neutrality. The local densities of electrons and ions are given by

$$n_e \approx n_0 \exp[\phi / T_e] \quad (7-1)$$

$$n_i \approx n_0 \exp[(-e\phi - \psi_i) / (eT_i)], \quad (7-2)$$

where n_0 is a density at a reference position and T_e and T_i are, respectively, the electron and ion temperatures in the unit of eV. The quasi-neutrality requires that $n_e \approx n_i$. It is found from Eqs.(7-1) and (7-2) that

$$n_i = n_0 \exp[-\psi_i / \{e(T_i + T_e)\}]. \quad (7-3)$$

The ion flux Γ_i passing over the ponderomotive potential barrier is written as

$$\Gamma_i = \Gamma_{i0} \exp[-\psi_i / \{e(T_i + T_e)\}], \quad (7-4)$$

where Γ_{i0} is the ion flux in the absence of RF. Even when the electrons are directly plugged by a ponderomotive potential Eq.(7-4) applies provided ψ_i is replaced by ψ_e .

Then we proceed to obtain an explicit form of ψ_i in a

cylindrical plasma. We assume the RF electric field in the plasma launched from the antenna in Fig.2-7(c) in Chap. 2 to be

$$\begin{aligned}
 E_r^{(1)} &= \frac{1}{\sqrt{2}} [\epsilon_+(E_1+E_2)\cos\phi_1 + \epsilon_-(E_3+E_4)\cos\phi_2]e^{-K_{\parallel}z} \\
 E_{\theta}^{(1)} &= \frac{1}{\sqrt{2}} [\epsilon_+(-E_1+E_2)\sin\phi_1 + \epsilon_-(-E_3+E_4)\sin\phi_2]e^{-K_{\parallel}z} \\
 E_z^{(1)} &= 0 \quad (z \geq 0), \quad (7-5)
 \end{aligned}$$

where $K_{\parallel}$ is the axial scale length of the localized RF field, $E_1(r)$ to $E_4(r)$ are the field amplitudes, $\phi_1 = \theta + \omega t$, $\phi_2 = -\theta + \omega t$, and E_2 and E_3 are zero at $r = 0$. In Eq.(7-5), $\epsilon_+ = \sqrt{2}$ and $\epsilon_- = 0$ ($\epsilon_+ = 0$ and $\epsilon_- = \sqrt{2}$) for the $m = +1$ (-1) mode, and $\epsilon_+ = \epsilon_- = 1$ for the $m = \pm 1$ mode if the antenna current is kept constant. The axial ponderomotive force can be calculated from Eqs.(2-11) and (2-12) in Chap. 2. Assuming that B_0 is uniform and the plugging is adiabatic, we find that

$$\begin{aligned}
 \psi_i &= \frac{e^2}{4m_i} \frac{1}{\omega} \left[\frac{(\epsilon_+E_1)^2 + (\epsilon_-E_3)^2}{\omega - \omega_{ci}} + \frac{(\epsilon_+E_2)^2 + (\epsilon_-E_4)^2}{\omega + \omega_{ci}} \right. \\
 &\quad \left. + 2\epsilon_+\epsilon_-\left(\frac{E_1E_3}{\omega - \omega_{ci}} + \frac{E_2E_4}{\omega + \omega_{ci}}\right)\cos 2\theta \right] e^{-2K_{\parallel}z} \quad (7-6)
 \end{aligned}$$

For the case of the $m = +1$ (-1) mode, the ponderomotive potential near $(r,z) = (0,0)$ is represented as

$$\psi_i = \frac{e^2}{2m_i} \frac{1}{\omega} \frac{E_1^2}{\omega - \omega_{ci}} \quad \left(\psi_i = \frac{e^2}{2m_i} \frac{1}{\omega} \frac{E_4^2}{\omega + \omega_{ci}} \right),$$

and hence the conventional expression is recovered. For the $m = \pm 1$ mode, the dominant term in ψ_i is

$$\frac{e^2}{4m_i} \frac{1}{\omega} \left[\frac{E_1^2}{\omega - \omega_{ci}} + \frac{E_4^2}{\omega + \omega_{ci}} \right],$$

which is smaller than ψ_i for the $m = \pm 1$ mode near $r = 0$. Furthermore, for larger r , the ponderomotive potential is azimuthally asymmetric producing an asymmetric density distribution. The resultant density gradient may bring about an electric field $E_\theta^{(2)}$ giving rise to a convective plasma loss due to $E_\theta^{(2)} \times B_0$ drift. Convective plasma loss may be produced also through $S_{i\theta} \times B_0$ drift with $S_{i\theta}$ being the azimuthal component of the ponderomotive force of the $m = \pm 1$ RF field divided by e , as described in Chaps. 5 and 6. Therefore, it is predicted that the $m = \pm 1$ RF field is superior for the RF plugging to the $m = \pm 1$ RF field in viewpoints not only of efficient plugging but of overall plasma confinement.

3. BASIC EXPERIMENTS IN HIEI

A. Experimental Apparatus

The experiments are carried out in the axisymmetric mirror machine HIEI in which a hydrogen plasma of density from 10^{10} cm^{-3} to $5 \times 10^{13} \text{ cm}^{-3}$ is produced with TPD or MPD plasma guns, in a static magnetic field up to 1.1 T at the throat. The length between two mirror points is 1.2 m, and the mirror ratio is variable from 1.4 to 4. Figure 7-1 shows the sche-

HIEI

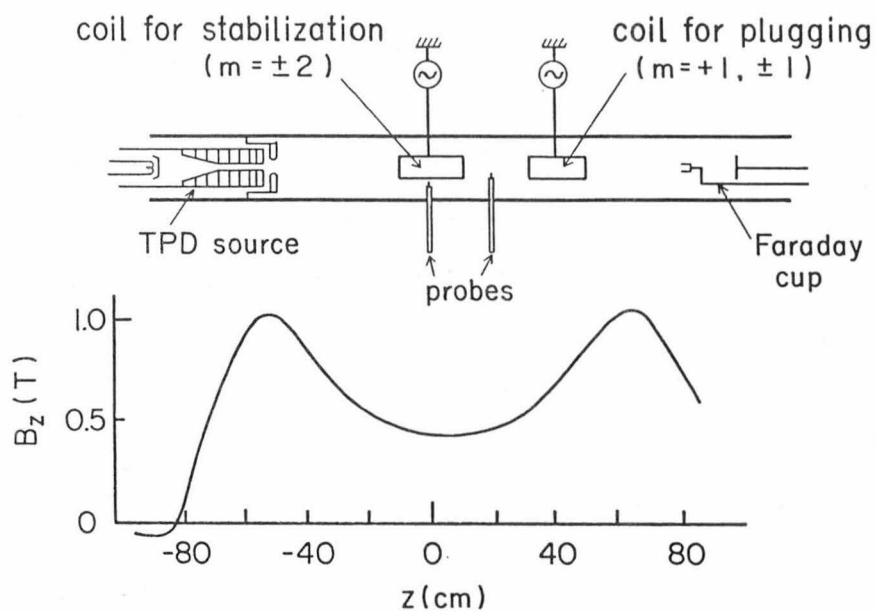


Fig.7-1 Experimental setup.

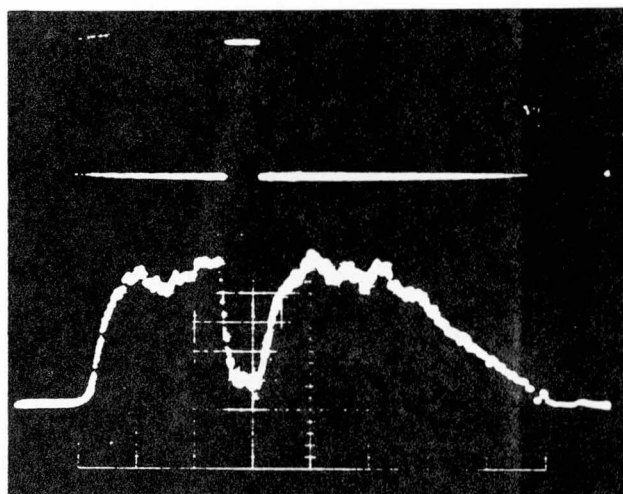


Fig.7-2 Typical oscillogram of the end loss flux.
The upper trace is the gate signal to the RF oscillator.
Time scale is 0.2 msec/div.

matic diagram of HIEI. The plasma diameter at the throat is 4 cm (FWHM), which is fixed by the limiter made of pyrex glass. Typical electron temperatures are $T_e = (15-20)$ eV and ion temperatures without heating are $T_i \approx 15$ eV. The plasma injected through the left mirror point is terminated by the metal end plate which is located outside the right mirror point and has an entrance slit to an end loss analyzer. This plate ensures MHD stability of the plasma through line-tying effects. (The plate is removed in the experiment in Subsection 3E.)

The RF antenna for the plugging is essentially the same as that used in Chap. 5. The configuration of the antenna is shown in Fig.2-7(c) in Chap. 2. We install this antenna 40 cm downstream from the midplane and generates $m = +1, -1$, and ± 1 RF fields selectively. An RF oscillator which can deliver the RF power of 200-kW, 0.5-msec pulse at a frequency of 8 to 10 MHz is connected to the antenna through a phase shifter and impedance matching circuits.

In Subsection 3E, we apply the RF stabilization together with the RF plugging. The antenna at the midplane generates $m = \pm 2$ RF field and suppresses the flute instability as described in Chaps. 3 and 4.

B. End Loss Reduction

We first perform the RF plugging by the $m = \pm 1$ RF field to confirm the result in TPD-III[2]. The end loss flux I_e measured with the end loss analyzer is shown in Fig.7-2 for $n_0 = 2 \times 10^{11} \text{ cm}^{-3}$, $\omega/\omega_{cia} = 1.1$, and $\omega/2\pi = 9.0$ MHz with ω_{cia}

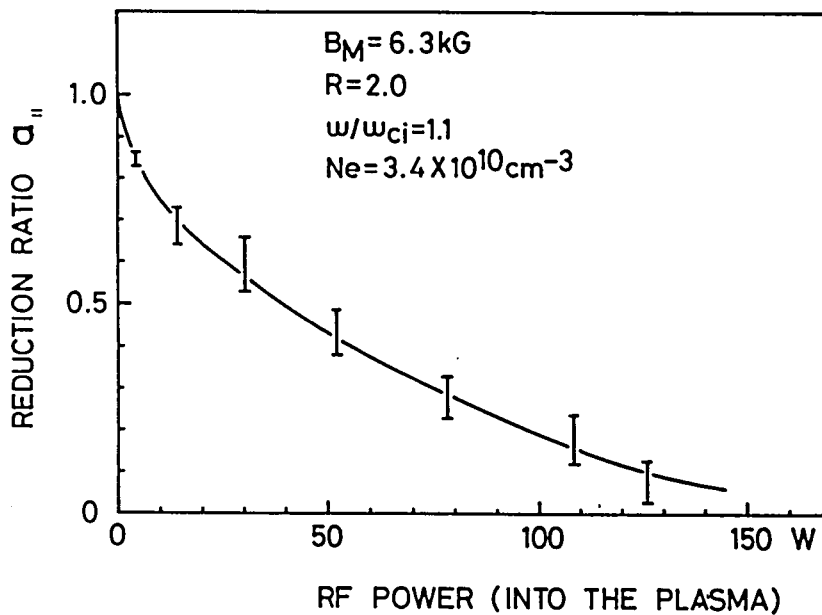


Fig.7-3 End loss reduction ratio versus the net RF power for $n_0 = 3.4 \times 10^{10} \text{ cm}^{-3}$, $B_M = 0.63 \text{ T}$, $R = 2$, $m = +1$, and $\omega/\omega_{cia} = 1.1$.

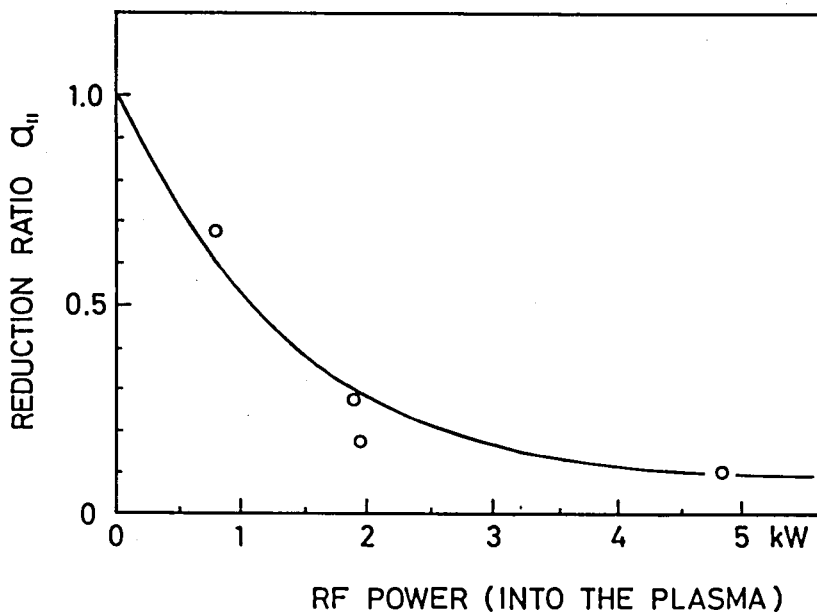


Fig.7-4 End loss reduction ratio versus the net RF power for $n_0 = (5-8) \times 10^{13} \text{ cm}^{-3}$. Other parameters are the same as in Fig.7-3.

being the ion cyclotron frequency at the plugging antenna. The value of I_e reduced to less than 20 % of the initial value with a net RF power of 200 W into the plasma. The ratio of end loss flux with RF to that without RF, $\alpha_{||}$, is plotted in Fig.7-3 as a function of the net RF power P_{net} for $n_0 = 3.4 \times 10^{10} \text{ cm}^{-3}$, $B_M = 0.63 \text{ T}$, $R = 2$, $m = \pm 1$, and $\omega/\omega_{cia} = 1.1$. The value of $\alpha_{||}$ decreases with increasing P_{net} . Figure 7-4 shows the same plot for $n_0 = (5-8) \times 10^{13} \text{ cm}^{-3}$. (Other parameters are the same as in Fig.7-3.) These characteristics in Figs. 7-2 to 7-4 are qualitatively the same as in TPD-III. We have also found that the net RF power necessary for $\alpha_{||} = 0.1$ increases proportionally to $n_0^{0.5}$ up to a density of 10^{14} cm^{-3} .

We then investigate the effect of RF mode on the end loss reduction. The value of $\alpha_{||}$ is given in Fig.7-5 as a function of the axial magnetic field B_0 for the RF fields of $m = +1$, -1 , and ± 1 modes. The plasma density is $(2-3) \times 10^{11} \text{ cm}^{-3}$ with $R = 4$, and the RF power is kept constant for three operating modes. The end loss flux is resonantly decreases near the cyclotron resonance for the $m = +1$ mode. Since B_0 varies more than 15 % under the RF antenna of 15 cm long, the resonance is rather broad. The value of $\alpha_{||}$ for the $m = -1$ mode is little dependent on B_0 . We also note that the resonant decrease of $\alpha_{||}$ for the $m = \pm 1$ mode is weaker than for the $m = +1$ mode. If the plasma density was decreased less than $1 \times 10^{11} \text{ cm}^{-3}$, the resonant point shifted to weaker B_0 indicating the change from electromagnetic to electrostatic coupling.

We measured the perpendicular component of the RF magnetic

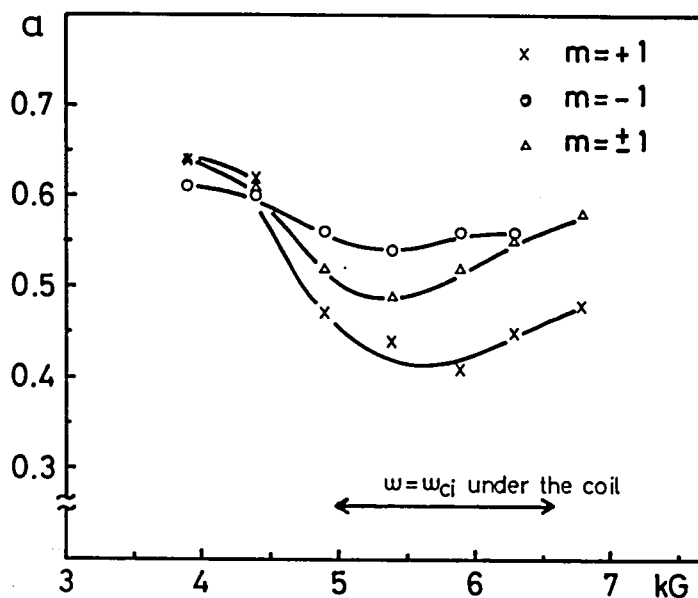


Fig.7-5 End loss reduction ratio versus the axial magnetic field strength for $m = +1$, -1 , and ± 1 modes. The horizontal line indicates that $\omega = \omega_{cia}$ point lies inside the RF antenna.

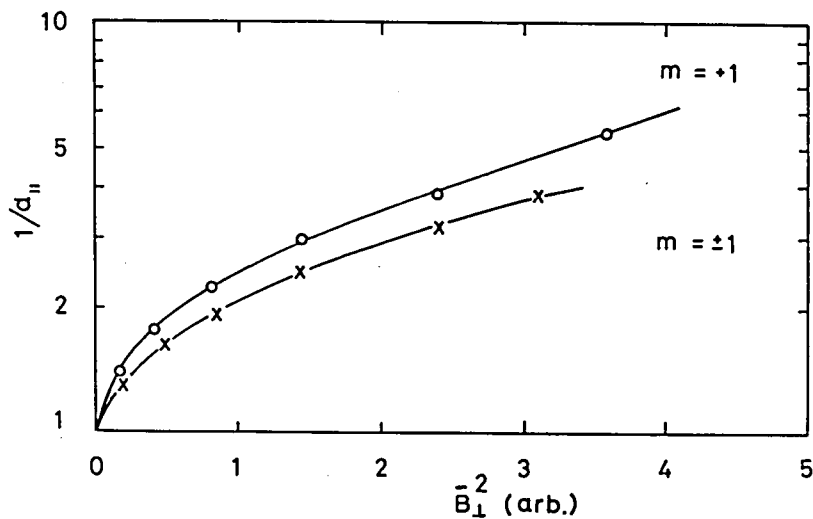


Fig.7-6 The reciprocal of the end loss reduction ratio versus the square of amplitude of the RF magnetic field for $m = +1$ and ± 1 modes.

field underneath the RF antenna, $B_{\perp}$, by using a magnetic probe. In Fig.7-6, the value of $\ln(1/\alpha_{\perp})$ is given versus $B_{\perp}^2$ for $m = +1$ and ± 1 modes. It should be noted that $B_{\perp}^2 = B_r^2$ (or B_0^2) for the $m = +1$ mode and $B_{\perp}^2 = B_r^2 + B_0^2$ for the $m = \pm 1$ mode. We can see that the same value of $\alpha_{\perp}$ is obtained with less $B_{\perp}^2$ for the $m = +1$ mode.

It is predicted from Eq.(7-6) that ψ_i becomes resonantly large near $\omega = \omega_{ci}$ for the $m = +1$ and $m = \pm 1$ modes and that ψ_i for the $m = +1$ mode is larger than for the $m = \pm 1$ mode since the numerator of the first term in the square brackets for the former is larger than for the latter. These characteristics are just what Fig.7-5 shows. The non-resonant behavior of $\alpha_{\perp}$ for the $m = -1$ mode is also explained by Eq.(7-6).

The ponderomotive potential given in Eq.(7-6) is expressed in an alternative form in terms of the perpendicular RF magnetic field at $r = 0$. The results are

$$\psi_i(r \approx 0, z=0) \sim \frac{e^2}{4m_i} \frac{\omega}{K_{\perp}^2} \frac{B_L^2}{\omega - \omega_{ci}} \quad (m = +1) \quad (7-7)$$

$$\psi_i(r \approx 0, z=0) \sim \frac{e^2}{4m_i} \frac{\omega}{K_{\perp}^2} \left[\frac{B_L'^2}{\omega - \omega_{ci}} + \frac{B_R'^2}{\omega + \omega_{ci}} \right] \quad (m = \pm 1), \quad (7-8)$$

where B_L is the amplitude of left-hand component of the $m = +1$ RF field and B_L' and B_R' are the amplitudes of left- and right-hand components of the $m = \pm 1$ RF field, respectively. Experimentally observed amplitude $B_{\perp}$ is related as $(B_{\perp}^2)_{+1} = B_L^2$ for the $m = +1$ mode and $(B_{\perp}^2)_{\pm 1} = B_L'^2 + B_R'^2$ for the $m = \pm 1$ mode. In Fig.7-6, we varied $B_{\perp}^2$ holding $(B_{\perp}^2)_{+1} = (B_{\perp}^2)_{\pm 1}$. In this case, Eqs.(7-7) and (7-8) predict that $(\psi_i)_{+1} > (\psi_i)_{\pm 1}$

for $\omega \approx \omega_{ci}$ since $B_L^2 > B_L'^2$. The ordinate in Fig.7-6 is proportional to ψ_i as seen from Eq.(7-4). So, the experimental result agrees qualitatively with this prediction. Using the ratio of $(\psi_i)_{+1}$ to $(\psi_i)_{-1}$ which is obtained from Fig.7-6, we find from Eqs.(7-7) and (7-8) that $B_L'/B_R' \sim 1.9$ for the $m = +1$ mode. The value of B_L'/B_R' obtained directly from the measured amplitudes of the RF magnetic field was about 1.2. The agreement between both values of B_L'/B_R' is fairly good showing the reasonableness of Eqs.(7-7) and (7-8).

We have thus verified both theoretically and experimentally that the $m = +1$ RF field is more effective for the end loss reduction than the $m = \pm 1$ RF field used conventionally.

C. Particle Confinement Time

The particle confinement time of a plasma is given by

$$\tau_p = \frac{N}{\Gamma_u + \Gamma_d + \Gamma_\perp}, \quad (7-9)$$

where N is the total number of particles in the mirror region, Γ_u and Γ_d are the parallel loss fluxes through the mirror points of upstream and down stream sides, respectively, and $\Gamma_\perp$ is the total perpendicular loss flux. We estimate these values as follows; From the measured density profile at the midplane, we approximate that $N = L \int_0^{r_L} 2\pi r n_0(r) dr$, where r_L is the radius of the limiter and L is the length between the positions where Γ_u and Γ_d are measured. (L is nearly equal to the length between the mirror points.) The value of Γ_d is measured by the end loss analyzer used in the previous

Subsection. The loss flux toward the plasma source is obtained by scanning a small Faraday cup over the plasma radius. Three loss collectors which measure perpendicular loss flux are set at three different axial positions, left and right mirror points, and the midplane. Each collector consists of negatively biased plane electrodes and a grounded mesh in front of them. We estimate the total $\Gamma_{\perp}$ by integrating $\Gamma_{\perp}(z)$ which is obtained from an interpolation of measured values.

Figure 7-7 shows the ratio of Γ_d with RF to that without RF, $\alpha_{\parallel}$, and the ratio of $\Gamma_{\perp}$ with RF to that without RF, $\alpha_{\perp}$, as a function of the net RF power divided by the density for $B_M = 0.71$ T, $R = 2$, and $\omega/\omega_{cia} = 1.2$. The plasma density is varied as indicated in the figure. The value of $\alpha_{\parallel}$ decreases with the normalized RF power for both the $m = +1$ and ± 1 modes, while $\alpha_{\parallel}$ for the $m = +1$ mode is smaller than for the $m = \pm 1$ mode indicating better plugging efficiency for the former. These results are the same as those described in the previous Subsection. The perpendicular loss flux increases proportionally to the normalized RF power regardless of the RF mode. The increase of $\alpha_{\perp}$ may be due to flute instabilities caused by the reduced line-tying effect through the RF plugging. The value of Γ_u decreased to (70-80) % of the initial value for $P_{net}/n_0 \sim 1 \text{ kW}/10^{11} \text{ cm}^{-3}$, and was insensitive to the RF mode. The ratios of $\Gamma_{\perp}$, Γ_d , and Γ_u to the total loss flux were ~ 0.15 , ~ 0.8 , and ~ 0.05 , respectively in the case of no RF applied. It was also observed that the value of N increased by a factor of 2.5 when the $m = +1$ RF field with $P_{net}/n_0 \sim 0.9 \text{ kW}/10^{11} \text{ cm}^{-3}$ was applied. For the $m = \pm 1$ mode, the increment factor was

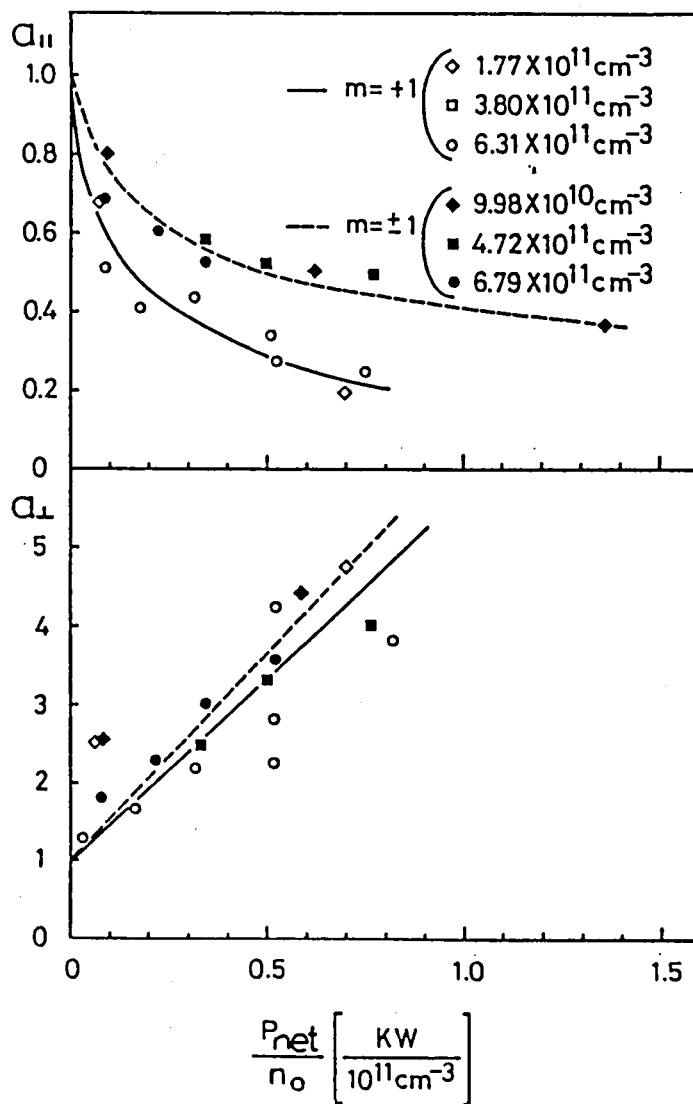


Fig.7-7 The values of $\alpha_{||}$ and $\alpha_{\perp}$ as a function of the net RF power divided by the density for $m = +1$ and ± 1 modes. The density is varied from $1 \times 10^{11} cm^{-3}$ to $6.8 \times 10^{11} cm^{-3}$.

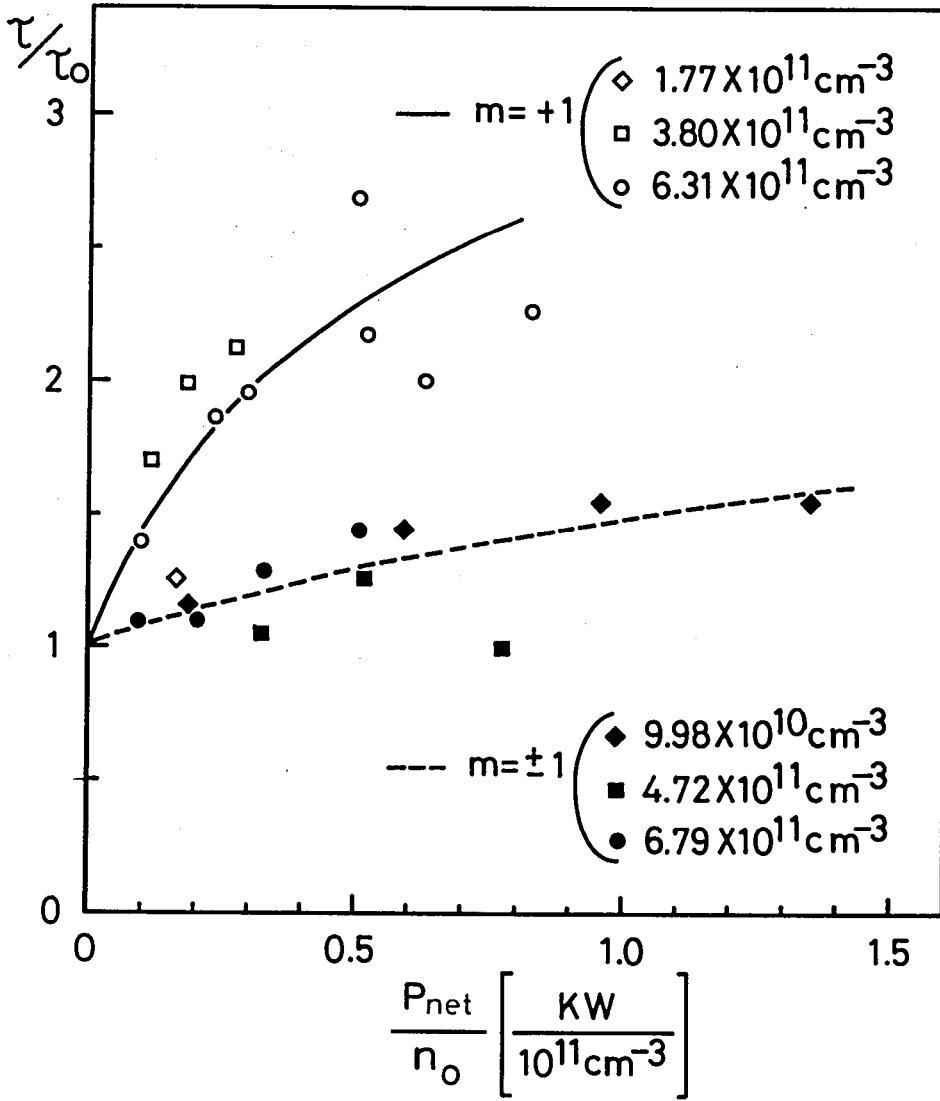


Fig.7-8 Normalized particle confinement time versus the net RF power divided by the density for the $m = +1$ and ± 1 plugging RF fields. The value of τ_{p0} is 0.45 msec.

about 2. The increase of N is brought about not only by the RF plugging but also by RF ionization of residual neutral gas. This is the very reason why we use Eq.(7-9) to estimate the particle confinement time. The value of the particle confinement time without the RF plugging, τ_{p0} , was obtained from Eq.(7-9) to be 0.45 msec. We plot the value of τ_p normalized by τ_{p0} in Fig.7-8 as a function of P_{net}/n_0 for the $m = +1$ and ± 1 modes. The particle confinement time is indeed improved by more than a factor of 2 by applying the RF plugging of $m = +1$ mode. The use of conventional $m = \pm 1$ RF field results in an improvement of τ_p by only a factor of 1.5.

It seems that the value of τ_p/τ_{p0} saturates for larger P_{net}/n_0 . This is due to the increase of Γ_1 which amounts to (60-80) % of the total loss flux for larger P_{net}/n_0 . Since the increase of Γ_1 is caused by the flute instabilities, it may be easy to reduce Γ_1 by using the RF stabilization method simultaneously. We will describe the preliminary experiment of the RF plugging of RF-stabilized plasma in Subsection 3E.

D. Effect on Energy Confinement

Since the RF field near the fundamental ion cyclotron frequency is used for plugging, fractional ions are heated through ion cyclotron resonance. Figure 7-9 shows the diamagnetic signal measured at 20 cm upstream from the plugging antenna. The $m = +1$ 270-W RF power is applied to a plasma of $n_0 = 2.2 \times 10^{11} \text{ cm}^{-3}$ in a magnetic field of $B_M = 0.69 \text{ T}$ and $R = 1.7$. The increment of the diamagnetic signal depends strongly on the magnetic field strength. In Fig.7-10 is

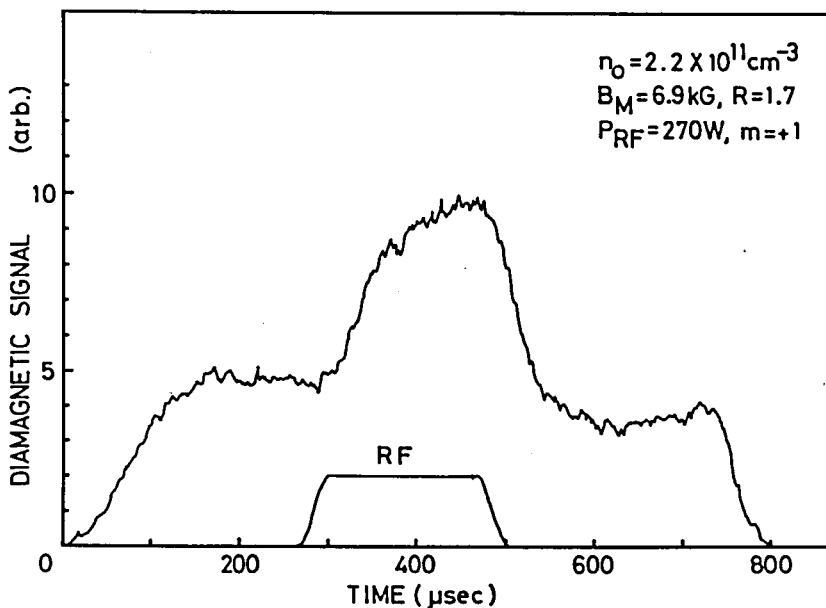


Fig.7-9 Raw data of the diamagnetic signal versus time when the plugging RF field of $m = +1$ mode is applied.

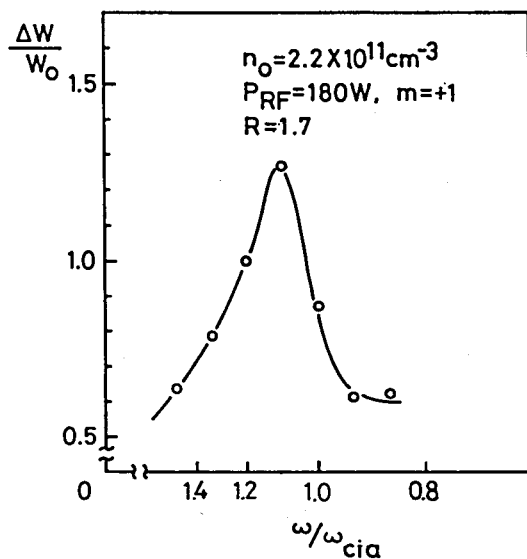


Fig 7-10 Increment of the plasma energy density divided by the initial value versus the axial magnetic field strength. The peak of the curve corresponds to the point where $\omega/\omega_{cia} = 1.1$.

shown the increment of the perpendicular plasma energy density ΔW divided by the initial value W_0 versus the magnetic field strength.

The value of ΔW is resonantly peaked at $\omega/\omega_{cia} = 1.1$, indicating that the heating is due to the cyclotron acceleration. The deviation of the ω/ω_{cia} value for the maximum ΔW from unity is because the increase of W is partly produced by the RF plugging effect which is maximum at $\omega/\omega_{cia} = (1.2-1.3)$.

The perpendicular component of the plasma energy W is expected to vary as

$$W(t) = W_0 + P_{net}\tau_E[1 - \exp(-t/\tau_E)], \quad (7-10)$$

where t is a time after the start of the RF pulse and τ_E is the energy confinement time with the RF plugging. If the RF pulse is switched off after W reaches a maximum value W_{max} ($\approx W_0 + P_{net}\tau_E$), W decays as

$$W(t) = W_{max}\exp(-t'/\tau_{E0}), \quad (7-11)$$

where t' is a time after switching off the RF pulse and τ_{E0} is the energy confinement time without the RF plugging.

Figure 7-11 shows the product of plasma energy and energy confinement time as a function of P_{net}/n_0 . The value of $W_{max}\tau_E$ is increased by a factor of 5 with the $m = +1$ RF field with $P_{net}/n_0 \approx 0.5 \text{ kW}/10^{11} \text{ cm}^{-3}$, while for the $m = \pm 1$ RF field with the same power, $W_{max}\tau_E$ is increased by a factor of 3. The absolute value of τ_E with the $m = +1$ RF field is 40

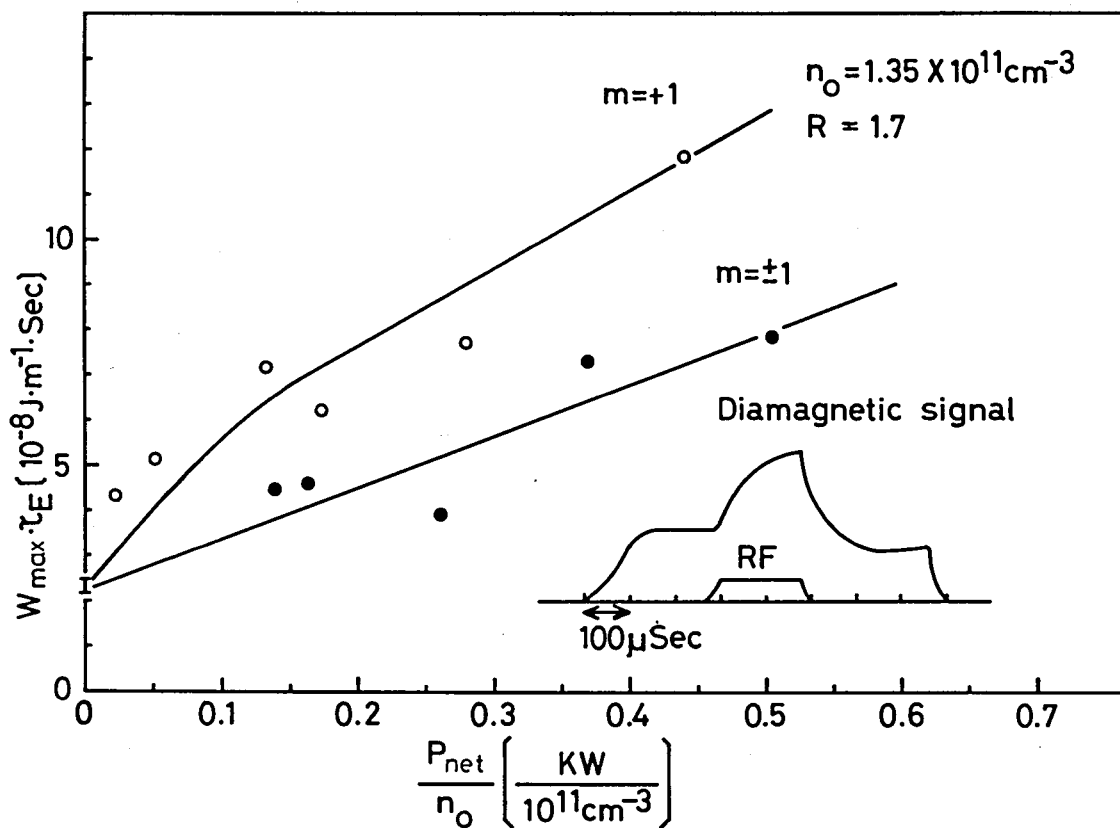


Fig.7-11 Product of maximum energy density achieved and energy confinement time versus the net RF power divided by the density for the $m = +1$ and ± 1 plugging RF fields.

to 52 μ sec, which is about twice as long as that without RF.

E. RF plugging of an RF-Stabilized Plasma

It has been observed that, if the RF power is increased to obtain larger end loss reduction, the perpendicular loss flux increases due to the flute instabilities. The onset of the instability is attributed to a reduction of line-tying effects caused by the RF plugging. The flute instabilities can be easily stabilized using the method described in Chaps. 3 and 4. So, the simultaneous application of the RF stabilization and the RF plugging will surely improve the particle confinement time more than observed in Subsection 3C. Before performing such an experiment, we think it is important to investigate whether any combined effect other than additive effects exists or not when the plugging RF and the stabilizing RF are both applied.

We now replace the metal end target by a target made of pyrex glass to remove the line-tying effects. The end loss is measured by an end loss analyzer of 6 mm in diameter and 10 mm long which is movable along the radius. The analyzer is located at 10 cm upstream from the target. The RF antenna for stabilization is installed at the midplane as shown in Fig.7-1. The structure and RF circuits of the stabilizing antenna are the same as those described in Chap. 3.

The end loss flux measured at a radius $r/r_L = 0.4$ is shown in Fig.7-12 together with the gate signal to the RF oscillator. By applying the $m = \pm 1$ RF field with $\omega/\omega_{cia} = 1.2$, the loss flux is reduced to about 50 % of the initial

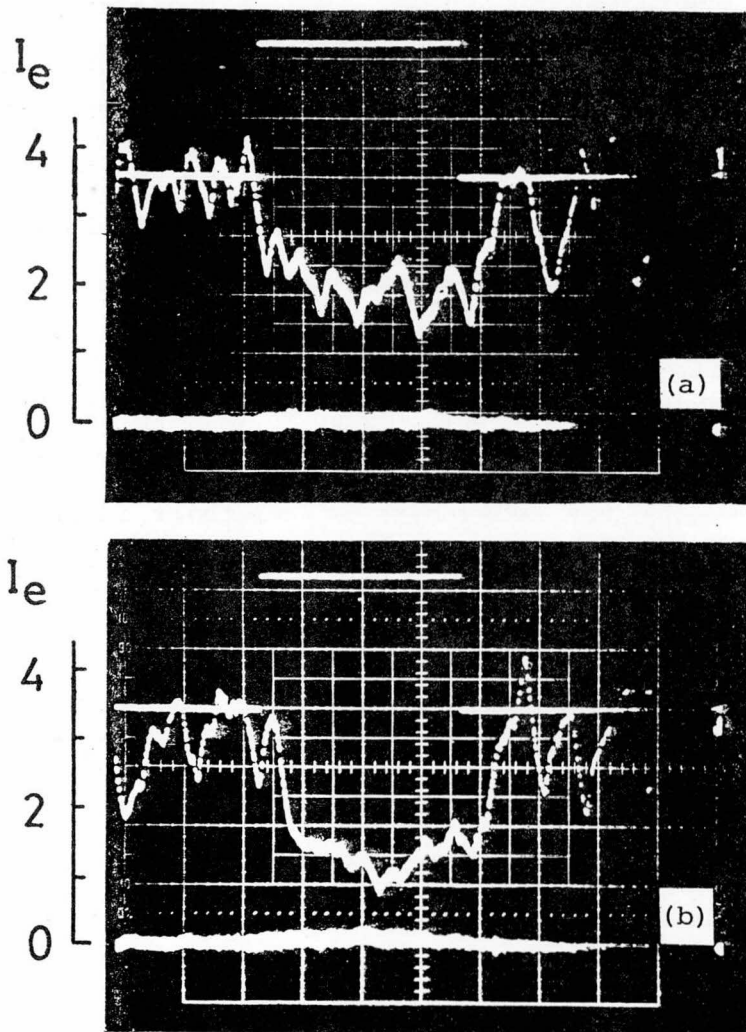


Fig.7-12 Typical oscillogram of end loss flux at $r/r_L = 0.4$ for the cases of (a) RF plugging only and (b) simultaneous application of RF plugging and RF stabilization. The upper trace in each oscillogram is the gate signal to the RF oscillator and the time scale is 50 $\mu\text{sec}/\text{div}$.

value for $n_0 \approx 10^{11} \text{ cm}^{-3}$, $B_M = 0.69 \text{ T}$, and $R = 2.3$. It is, however, found from Fig.7-12(a) that the flute instability leads to bursts of the end loss flux even during the plugging RF pulse. No stabilization effect was observed both at the midplane and at the throat. Figure 7-12(b) shows the end loss flux for the case of simultaneous application of the plugging RF and the stabilizing RF. The total RF power is the same as in (a). As the stabilizing RF antenna is excited by the same RF oscillator, the value of ω/ω_{cia} is 2.4 at the midplane. We find not only the stabilization of the flute mode but also the improved RF plugging with $\alpha_{||} = 0.35$ during the RF pulse. We also note in the separate experiment that no reduction of the end loss flux occurs by the $m = \pm 2$ stabilizing RF field.

In Fig.7-13, the reciprocal of $\alpha_{||}$ is given versus the RF power for two cases. The cross points show the case with the plugging RF only, while the open circles show the case in which the plugging and stabilizing RF are applied simultaneously. The abscissa is the total RF power to the two antennae. Larger reduction of the end loss flux is achieved with less RF power in the latter case.

The ion saturation current at the midplane corresponding to the case in Fig.7-12 is shown in Fig.7-14. The Langmuir probe was inserted at $r/r_L = 0.2$, where the fluctuation of flute instability was small. The increase in the ion saturation current for (b) the simultaneous application of two RF fields is larger than for (a) the plugging RF only.

Comparing Figs.7-12(a) and (b), we note that the decrease of the fluctuation amplitude in (b) is due to an additive

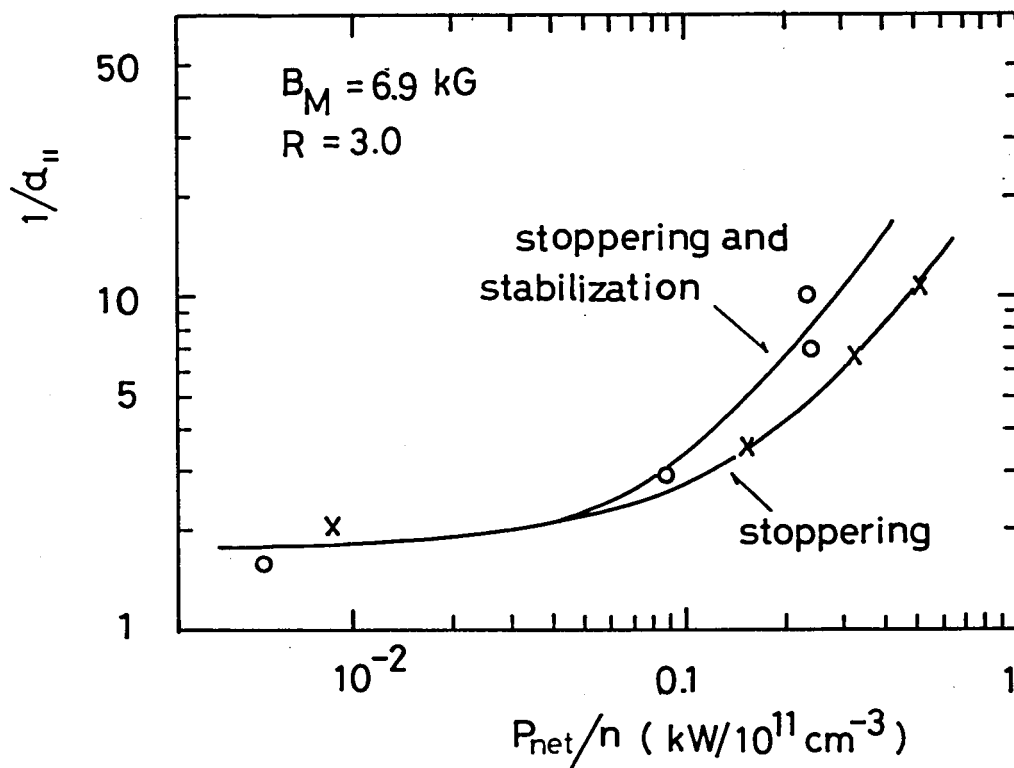
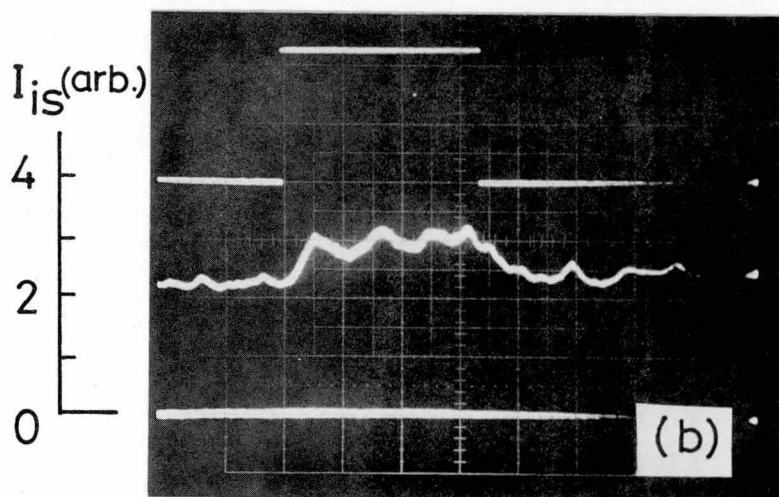
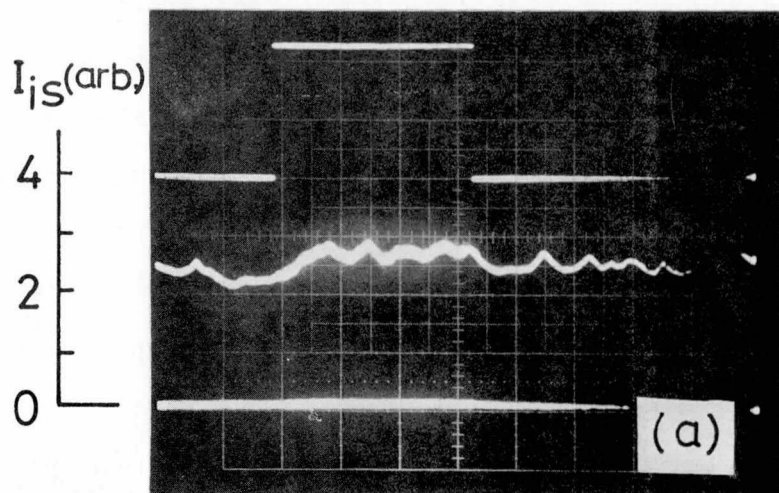


Fig.7-13 The reciprocal of $\alpha_{\parallel}$ versus the net RF power divided by the density for two cases of RF plugging only and the simultaneous application of RF plugging and RF stabilization.



50 μ sec/div.

Fig.7-14 Typical oscillogram of the ion saturation current at the midplane for the cases of (a) RF plugging only and (b) simultaneous application of RF plugging and RF stabilization. The upper trace in each oscillogram is the gate signal to the RF oscillator and the time scale is 50 μ sec/div.

effect of the RF stabilization. But the further reduction of the end loss flux is considered to be a combined effect of the RF plugging and the RF stabilization. The larger increase of the midplane density observed in Fig.7-14(b) may be caused by larger decreases of Γ_d and Γ_i as compared with the case in (a). A rough estimation of τ_p from Eq.(7-9) in the cases of Figs. 7-12 and 7-14 shows that τ_p is increased by a factor of 1.4 by the superposition of the RF stabilization even if we assume that Γ_i is the same for both cases. As shown in Fig.3-20 in Chap. 3, Γ_i is significantly reduced by the RF stabilization. Hence, we can expect an improvement factor much larger than estimated.

The mechanism responsible for the combined effect, i.e., the further reduction of α_e with the RF stabilization, is not clearly understood at present. But a possible explanation is as follows; When the RF stabilization is superimposed, the density perturbation is reduced and the density profile changes especially near the periphery. If this change of the density profile is such that the region of the left-hand polarization of the plugging RF field extends to outer radii, the end loss tends to reduce further in outer radii. This explanation is consistent with the observation that the combined effect is significant in the periphery.

4. APPLICATION OF PLUGGING BY THE $m = +1$ RF FIELD TO RFC-XX

A. Experimental Apparatus

Figure 7-15 shows the schematic view of RFC-XX, details of which are described elsewhere [4]. A hydrogen plasma is injected from the TPD plasma gun to the central section through differential pumping sections. Typical plasma parameters in the central section are as follows; $n_0 = 2 \times 10^{12} \text{ cm}^{-3}$, $T_i = 20$ eV, and $T_e = 10$ eV. The magnetic field strengths are 2 T at the line cusp, 3.7 T at the point cusp, and 0.95 T at the central section. The cusp field ensures MHD-stability. The ion temperature is increased up to 250 eV through the ion cyclotron heating at the central section.

Two line cusp ends and two point cusp ends are plugged by RF fields in ICRF [9]. In the present experiment, only the point cusp end of downstream side is RF plugged to investigate effects of RF mode precisely. The RF antenna for the point cusp plugging, which is shown in Fig.7-16, consists of two pairs of type-III coils displaced azimuthally by 90° with respect to each other. Two 200-kW RF amplifiers tuned at $\omega/2\pi = 51$ MHz are excited by two signal generators with the same frequency but with a different phase angle Δ . Each output of the RF amplifier is connected to one of the type-III coils through an impedance matching line. Since the value of Δ is variable continuously, we can establish the RF field of any desired polarization. The ratio of the applied RF frequency to the local ion cyclotron frequency at the center of the RF coil is chosen to be 1.1. The RF field is measured by a

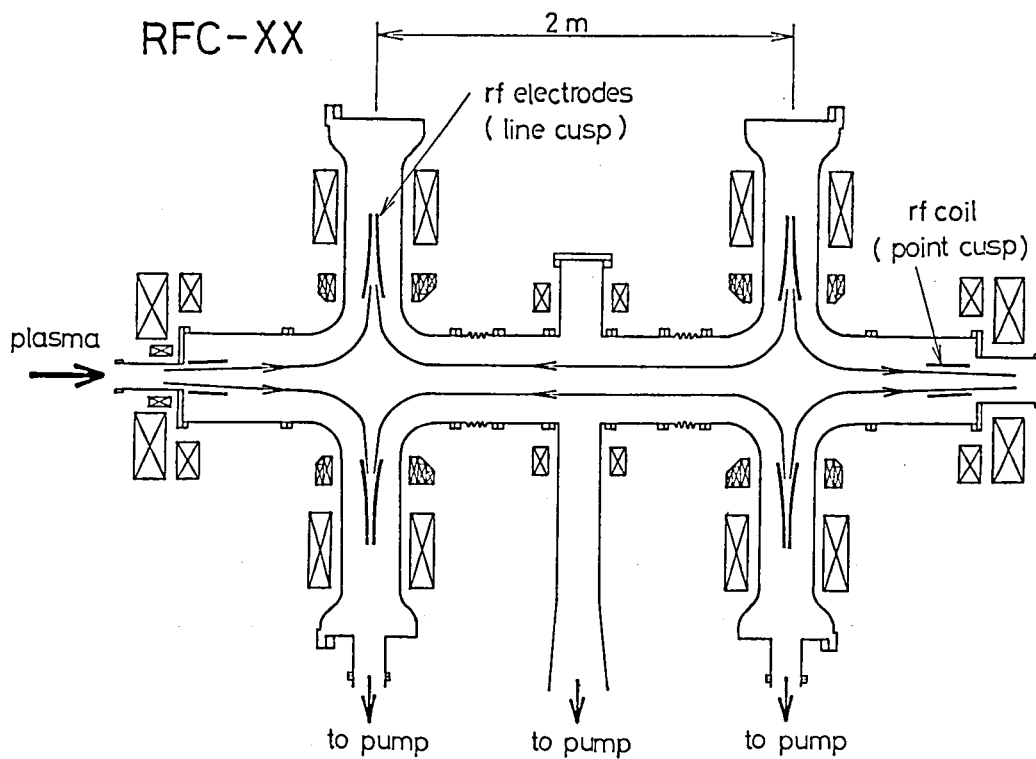
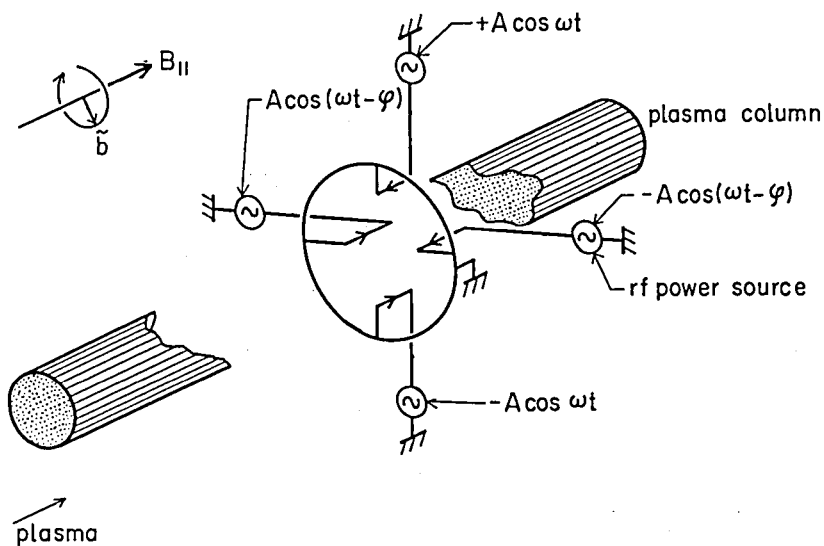


Fig.7-15 Setup for point cusp plugging in RFC-XX.



Antennas structure of Rotating rf

Fig.7-16 RF antenna for point cusp plugging.

calibrated small loop inserted 25 cm upstream from the RF coil.

B. Mode Effect on RF Plugging

We measured the RF magnetic field components, B_r and B_θ , and their phase difference by the small loop inserted onto the axis. The field amplitudes of right- and left-hand polarized components are calculated from the measured values and are plotted in Fig.7-17 as a function of the phase shift Δ of the coil currents. The point where $\Delta = -90^\circ$ ($+90^\circ$) corresponds to the $m = +1$ (-1) mode operation of the RF coil. The density was $n_0 \approx 5 \times 10^{11} \text{ cm}^{-3}$ at the cusp region of downstream side. The left- (right-) hand component of the RF field is maximum near $\Delta = -90^\circ$ ($+90^\circ$). This is in agreement with the theory described in Chap. 2. It seems that the right-hand polarized field is excited a little more effectively than the left-hand polarized field. It is thus confirmed that the polarization of the excited RF field is controllable through the phasing of the RF antenna in RFC-XX.

Figure 7-18 shows typical temporal evolutions of the output of a 24-GHz microwave interferometer, the end loss flux, and the RF voltage to the RF antenna. The line density was $2.5 \times 10^{12} \text{ cm}^{-2}$ and the antenna was operated in the $m = +1$ mode. When the RF field is applied the end loss flux through the point cusp decreases and, at the same time, the line density slightly decreases. We define the loss reduction rate $\alpha_{\perp}$ as a ratio of the end loss flux with RF to that without RF in the same way as in Sec. 3. The value of $\alpha_{\perp}$ is 0.5 in this

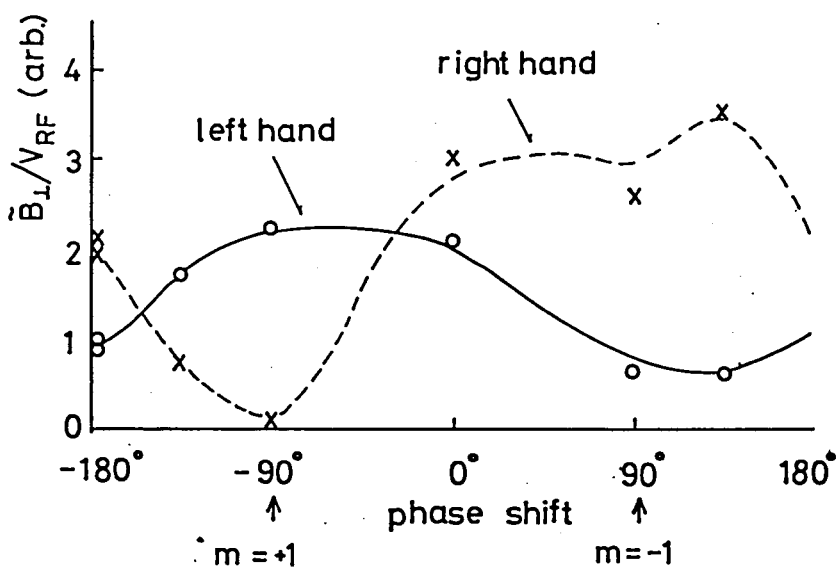


Fig.7-17 Field amplitudes of right- and left-hand polarized components versus the phase shift of RF currents.

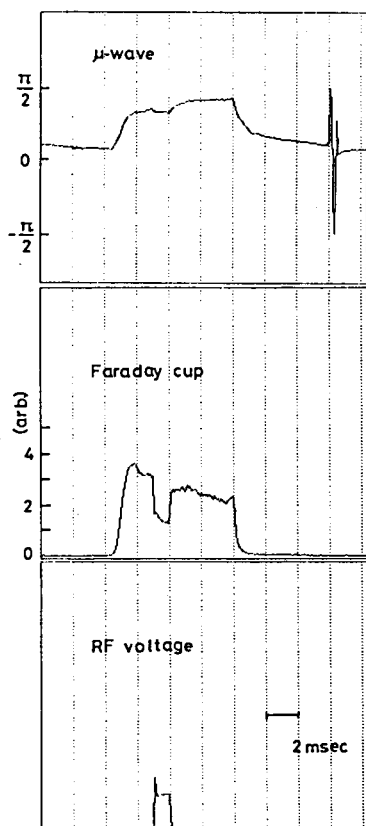


Fig.7-18 Temporal evolutions of the fringe of a microwave interferometer, the end loss flux, and the RF voltage applied to the antenna.

case. The dependence of $\alpha_{\parallel}$ on the applied RF voltage V_{RF} is measured for three cases of RF modes. As shown in Fig.7-19, the value of $\alpha_{\parallel}$ decreases with increasing the RF voltage regardless of the mode. But, for a fixed RF voltage, the reduction of the end loss flux is the maximum when the $m = +1$ RF field is applied. This is because the ponderomotive potential for ions is the largest for the left-hand circularly polarized RF field. The result in Fig.7-19 is qualitatively the same as that already obtained in HIEI. We have achieved further improvement of the plugging efficiency over that obtained in TPD-III by using the $m = +1$ RF field [2].

We plot the value of $\ln(1/\alpha_{\parallel})$ and the ratio of the density reduction associated with the RF plugging, $-\Delta n/n_0$, in Fig.7-20 versus Δ . Small changes in n_0 and V_{RF} in the measurement are corrected by normalizing the values by $(V_{RF}/n_0)^2$. The plasma parameters are the same as in Fig.7-17. It is evident that the plugging effect becomes very large for the $m = +1$ RF field. The minimum of $\ln(1/\alpha_{\parallel})$ occurs at $\Delta \approx 135^\circ$, not at 90° . Referring to Fig.7-17, we note that the polarization of the RF field at $\Delta \approx 135^\circ$ is almost completely right-handed rather than at $\Delta \approx 90^\circ$. This is the reason for the deviation of the minimum point of $\ln(1/\alpha_{\parallel})$ from $\Delta \approx 90^\circ$ in Fig.7-20. It is clearly seen from Fig.7-17 and 7-20 that the plugging effect increases with increasing the left-hand circularly polarized component of the RF field.

On the other hand, the decrease of the density at the cusp region seems to be significant when both the left- and right-hand components coexist. It is not surprising that the density

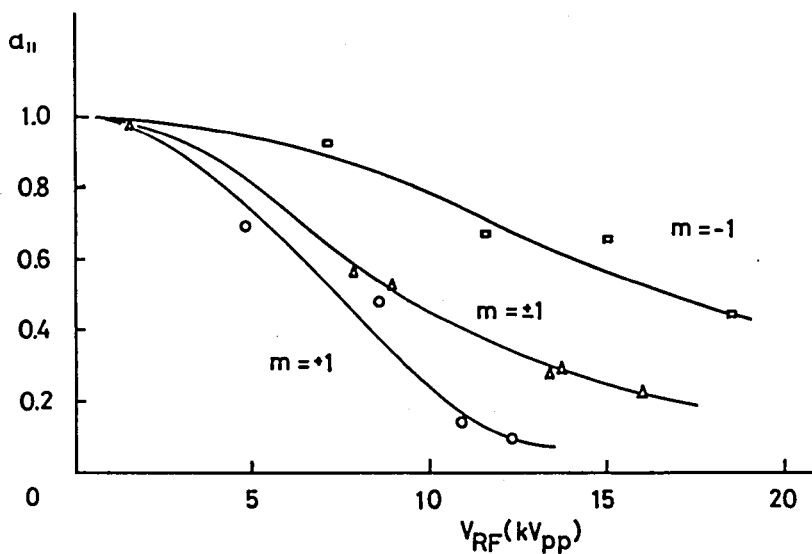


Fig.7-19 End loss reduction ratio versus the applied RF voltage for the $m = +1$, ± 1 , and -1 RF fields. The density is $0.7 \times 10^{12} \text{ cm}^{-3}$ and $\omega/\omega_{cia} = 1.43$.

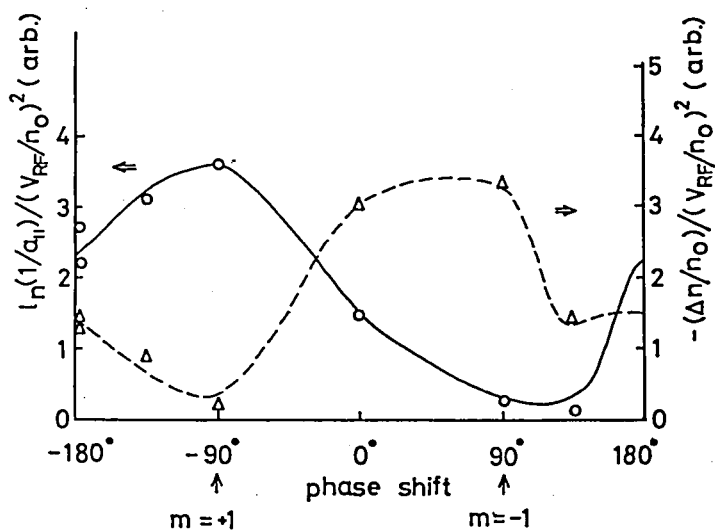


Fig.7-20 Values of $\ln(1/\alpha_{||})$ and $-\Delta n/n_0$ normalized by $(V_{RF}/n_0)^2$ as a function of the phase shift of RF currents.

in upstream side does not increase with the RF plugging. The plugging is applied only to one point cusp end, the loss flux through which is less than 10 % of the total loss flux in the absence of RF. Hence the point cusp plugging contributes little to an improvement of the total particle confinement, and rather enhance the line cusp losses through RF heating. In the case of the RF plugging with the $m = \pm 1$ mode, the significant decrease of the density was observed as in Fig. 7-20. This may be due to the convective loss peculiar to this RF mode described in Chap. 5.

C. Scaling of RF Plugging

We proceed to investigate the RF plugging of a high-temperature plasma. In addition to the RF antenna for plugging, another RF antenna is installed at the central section. This antenna excites slow waves of $m = +1$ mode at a frequency slightly lower than the local ion cyclotron frequency for ICRF heating [10]. In the experiment presented here, the value of T_i is varied from 20 eV to about 250 eV by controlling the RF power to the heating antenna. The ICRF-heated plasma passes through the field-null point before it reaches the point cusp. So, the velocity distribution function is considered to be isotropic at the RF plugging section. The plasma density is kept almost constant to be $0.7 \times 10^{12} \text{ cm}^{-3}$ at the line cusp of downstream side.

Figure 7-21 shows a typical oscillogram of (a) the diamagnetic signal at the central section, (b) the fringe shift of a 70 GHz microwave interferometer at the central section,

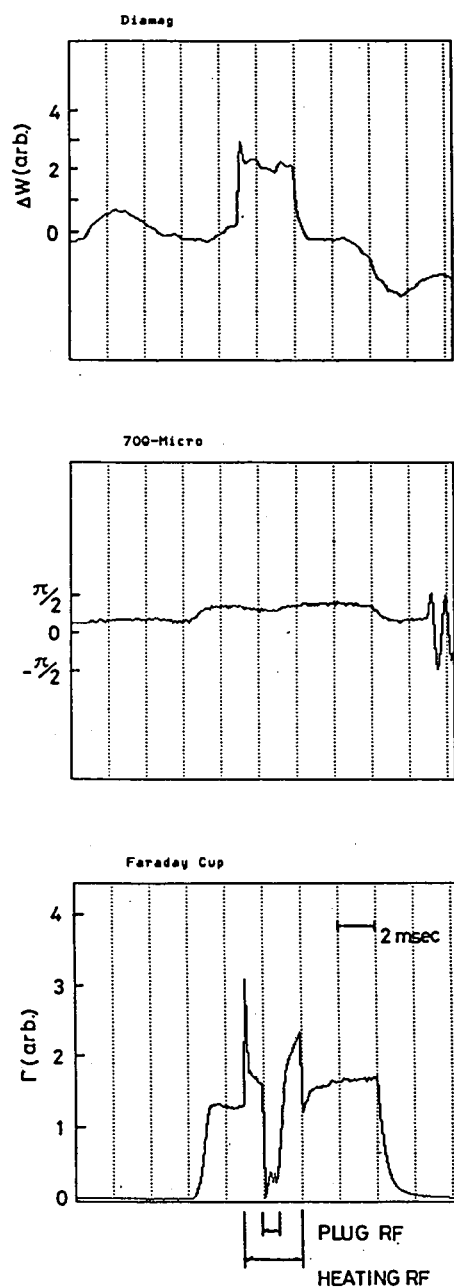


Fig.7-21 Temporal evolutions of (a) diamagnetic signal, (b) fringe shift of 70 GHz microwave interferometer, and (c) end loss flux through the point cusp. The density is $0.7 \times 10^{12} \text{ cm}^{-3}$ and $\omega/\omega_{cia} = 1.43$.

and (c) the end loss flux through the point cusp. When the heating RF power is applied, the diamagnetic signal increases with a rise time of ~ 40 μ sec followed by a gradual decrease of the density. The ion temperature was obtained from the diamagnetic signal and the density. The electron temperature has been confirmed to change little by the ICRF heating in a separate experiment. The value of T_i is about 220 eV in this case. It is seen that the end loss flux increases rapidly with the ICRF heating. By applying the plugging RF at 1 msec after the heating RF, the end loss flux is reduced to less than 10 % of the initial value. The value of ω/ω_{cia} is 1.43 at the center of the plugging antenna. We measure the value of α_n varying the RF voltage of the plugging antenna for several values of T_i . The result is displayed in Fig.7-22 for the $m = +1$ mode. The RF voltage necessary for $\alpha_n = 0.2$ is 10 kV_{pp} for $T_i = 20$ eV, while it increases to 17.5 kV_{pp} for $T_i = 220$ eV. The plot of $\ln(1/\alpha_n)$ versus V_{RF}^2/T_i is given in Fig.7-23 for $m = +1$, $m = \pm 1$, and $m = -1$ modes. It is clearly seen that α_n is proportional to $\exp(-V_{RF}^2/T_i)$ for each mode. Comparing this result with Eq.(7-4), we can obtain the absolute value of the ponderomotive potential. The value of ψ_i scales as

$$\psi_i = \eta_s V_{RF}^2, \quad (7-12)$$

where η_s is the efficiency of the RF plugging and takes the values of 1.02, 0.59, and 0.31 for the $m = +1$, ± 1 , and -1 mode, respectively, and ψ_i and V_{RF} are in the units of eV

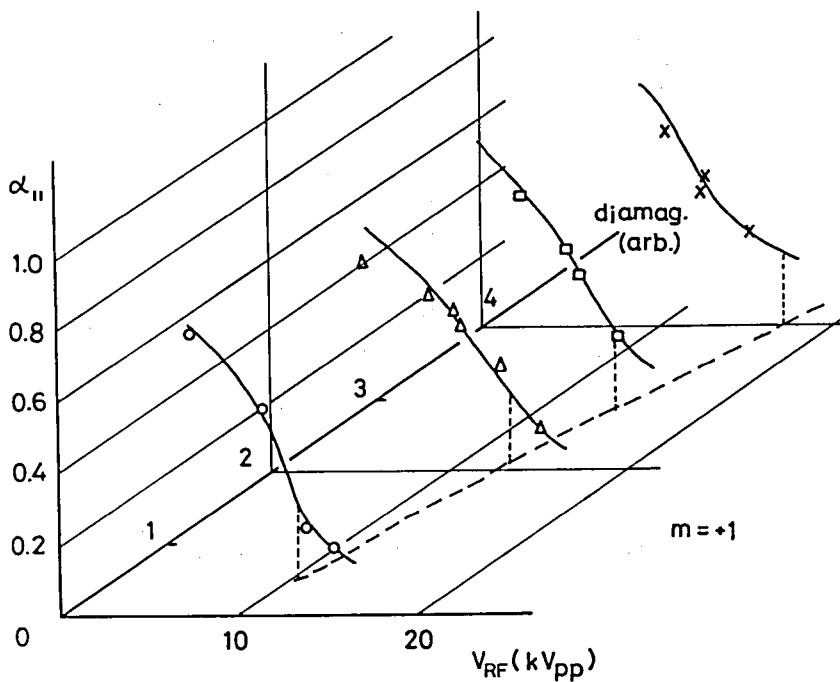


Fig.7-22 End loss reduction ratio versus the RF voltage applied to the plugging antenna for several values of T_i .

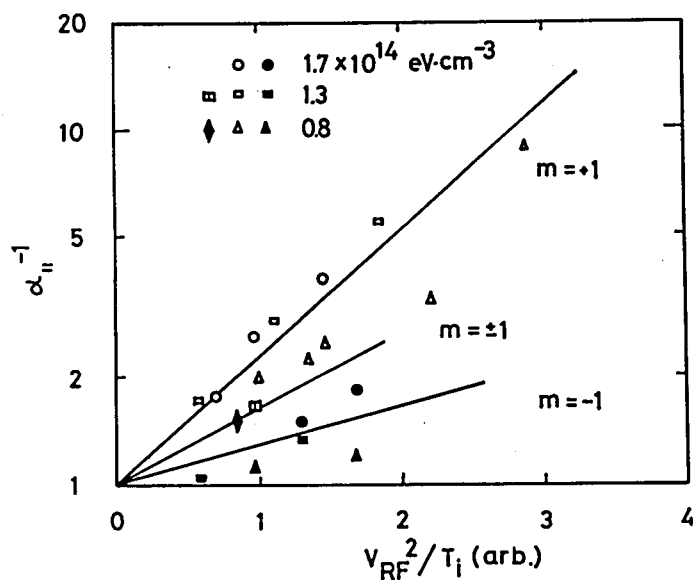


Fig.7-23 The reciprocal of $\alpha_{||}$ versus V_{RF}^2/T_i for $m = +1$, ± 1 , and -1 RF modes.

and kV_{pp} . (We have assumed that $T_i \gg T_e$.) The value of η_s for the $m = +1$ mode is larger than that for the $m = -1$ mode by a factor of 3.3. It is found from Eq.(7-6) that $(\psi_i)_{+1}/(\psi_i)_{-1} = (E_1/E_4)^2(\omega + \omega_{ci})/(\omega - \omega_{ci})$. Since the value of E_1/E_4 is 0.75 (see Fig.7-17), $(\psi_i)_{+1}/(\psi_i)_{-1}$ is calculated to be 3.2 by assuming $\omega_{ci} = \omega_{cia}$. This value is in quantitative agreement with the directly measured value in Fig.7-23. Therefore, it is proved that the ponderomotive potential for ions is essential for the RF plugging observed here. It is also confirmed that the use of the $m = +1$ RF field improves the plugging efficiency by a factor of 1.7 over the $m = \pm 1$ RF field used conventionally.

5.CONCLUSION

We have investigated the RF plugging by the rotating RF field of $m = +1$ mode in the axisymmetric mirror HIEI and the double cusp RFC-XX. The $m = +1$ RF field has been found to be more effective for end loss reduction than the $m = \pm 1$ RF field used conventionally. This result has been well explained by the theoretical model of ponderomotive potential. It has been also revealed that the plugging effect is entirely due to the ponderomotive potential on ions for both devices, since the plugging is strongly dependent on the field polarization.

In HIEI, by using the $m = +1$ RF field, the particle confinement time and the energy confinement time has been

improved by a factor of 2 and 1.7, respectively, compared to the $m = \pm 1$ RF field. The RF plugging of an RF-stabilized plasma has been investigated. It has been confirmed that both the plugging and the stabilization are achieved additively. It has been further observed that the end loss is decreased further due to the combined effect of the stabilizing RF and the plugging RF.

In RFC-XX, a scaling of the RF plugging with respect to the ion temperature has been obtained for the first time. The end loss reduction factor scales as $\exp(-\eta_s V_{RF}^2 / T_i)$ for T_i up to 250 eV. The value of η_s is the largest for the $m = \pm 1$ RF field among the three RF modes tested. This scaling formula agrees well with the theory in terms of the ponderomotive potential.

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CHAPTER 8

CONCLUSIONS

In this thesis, we have studied the plasma heating and confinement with mode-controlled ion-cyclotron RF fields both experimentally and theoretically. Summaries of the studies are given in the following.

In Chapter 2, we have considered linear and nonlinear effects of waves in the ion cyclotron frequency range, and found that the azimuthal mode of RF waves plays an important role in momentum transport and energy deposition. We have proposed the antennae for the forced excitation of RF waves with coherent azimuthal mode.

In Chapter 3, we have investigated the instability arising in the axisymmetric mirror machine HIEI. The instability observed in a wide range of density has been identified as the flute instability through the comparison of experimentally determined dispersion relation with the theoretical one.

We have demonstrated that the flute instability can be stabilized by applying the $m = \pm 2$ RF field only in the bad curvature region. The stabilizing mechanism has been revealed to be the reversal of azimuthal drifts of ions by the radial ponderomotive force. The simple calculation shows that the stabilization occurs when the ponderomotive force exceeds the effective gravity.

It has also been observed that the radial confinement of the plasma is improved with the stabilization of the flute instability. The increase of the plasma energy during the RF stabilization can be explained by the energy balance equation.

In Chapter 4, the RF stabilization of a plasma in the axisymmetric mirror HIEI has been investigated in various plasma parameters. We have obtained the empirical scaling formula of the RF stabilization that the required RF field strength increases proportionally to $n_0^{0.3 \pm 0.05} (T_i/R_c)^{0.4 \pm 0.1}$. This scaling formula agrees fairly well with the theoretical prediction.

From this scaling, a reasonable value of the RF electric field of 1.1 kV/cm is estimated for the RF stabilization of a fusion-grade plasma with $n_0 = 5 \times 10^{13} \text{ cm}^{-3}$, $T_i = 20 \text{ keV}$, and $R_c = 4.3 \text{ m}$.

From the experiments described in Chapters 3 and 4, the utility of the RF stabilization of the flute instability for high-density plasmas has been established.

In Chapter 5, we have investigated nonlinear effects in the Q-machine plasma subjected to large amplitude RF fields in ion cyclotron range of frequencies. We have observed for the first time that the convective plasma loss occurs due to the azimuthal ponderomotive force of the $m = \pm 1$ ICRF fields, which are used conventionally for ICRF heating. This result is in quantitative agreement with the theory based on the fluid model in which quasi-linear corrections are taken into account.

In order to eliminate the convective plasma loss, we have developed the RF antenna which can produce $m = +1$ circularly polarized RF field. By utilizing the $m = +1$ RF field, the convective loss is eliminated and the improvement in heating efficiency by more than 70 % is achieved.

In Chapter 6, ICRF heating by the $m = +1$ RF field has been applied to the small torus Synchromak and the simple mirror machine TPD-III. In Synchromak, it is found from the comparative measurements that the $m = +1$ RF field greatly improves the overall efficiency of ion heating over the $m = \pm 1$ RF field. The ion heating by the $m = \pm 1$ RF field is accompanied with the enhanced cross-field plasma loss. For an explanation of this result, the theory of convective plasma loss is again applied. The theory agrees quantitatively with the experimental result.

In TPD-III, although the anomalous convective loss caused by the $m = \pm 1$ RF field was not clearly observed as in HIEI-Q and Synchromak, we again confirmed that the $m = +1$ RF field could improve the heating efficiency by a factor of 1.5. The smaller improvement factor as compared with other two cases may be attributed to the much larger $E \times B$ rotation frequency in TPD-III.

We have found from the studies in Chapters 5 and 6 that the azimuthal symmetry of RF fields is indispensable for efficient heating without degradation of confinement.

In Chapter 7, we have investigated the RF plugging by the RF field of $m = +1$ mode in the axisymmetric mirror HIEI and the double cusp RFC-XX. The $m = +1$ RF field has been found

to be more effective for end loss reduction than the $m = \pm 1$ RF field used conventionally. This result has been well explained by the theoretical model of ponderomotive potential. It has been also revealed that the plugging effect is entirely due to the ponderomotive potential on ions for both devices, since the plugging is strongly dependent on the field polarization.

In HIEI, by using the $m = +1$ RF field, the particle confinement time and the energy confinement time has been improved by a factor of 2 and 1.7, respectively, compared to the $m = \pm 1$ RF field. The RF plugging of an RF-stabilized plasma has been investigated. It has been confirmed that both the plugging and the stabilization are achieved additively. It has been further observed that the end loss is decreased further due to the combined effect of the stabilizing RF and the plugging RF.

In RFC-XX, a scaling of the RF plugging with respect to the ion temperature has been obtained for the first time. The end loss reduction factor scales as $\exp(-\eta_s V_{RF}^2 / T_i)$ for T_i up to 250 eV. The value of η_s is the largest for the $m = +1$ RF field among the three RF modes tested. This scaling formula agrees well with the theory in terms of the ponderomotive potential.

Some problems are, however, left in open for future study. In Chapters 3 and 4, we used the simplified solution for the dispersion relation of the flute mode. A full solution of the dispersion relation is necessary for more detailed expla-

nation of the experimentally obtained empirical scaling of the RF stabilization. It is also necessary to extend the plasma parameters in the RF-stabilization experiments in order to prove the effectiveness of the RF stabilization in high-temperature plasmas.

Concerning the experiments in Chapters 5 and 6, the power level of ICRF heating was restricted to a few tens of kW. We must confirm the superiority of the $m = +1$ RF field in higher power levels.

The RF plugging experiment in RFC-XX described in Chapter 7 was done at one point cusp end. Plugging of all open ends is necessary for comparative measurements of overall particle and energy confinements.

As summarized above, this thesis has made some contributions to the researches of controlled nuclear fusion, especially to those of open ended magnetic confinement. We have extended the RF stabilization method so that it is applicable to a high-density, high-temperature plasma for complete axisymmetrization of a tandem mirror. It is worth emphasizing that our work has stimulated experimental, theoretical, and numerical investigations of the RF stabilization.

We have developed the ICRF heating and RF plugging with mode-controlled RF fields, which are much more efficient than those with the RF fields of conventional type. These methods have been successfully applied to present-day larger devices and are now giving further achievements.

APPENDIX A

1. WAVE FIELDS IN A CYLINDRICAL PLASMA

In plasma;

$$E_z^P = [A J_m(k_1 r) + B J_m(k_2 r)] e^{j\phi} \quad (A-1)$$

$$E_\theta^P = [A \{ \frac{j m}{r} L_2 J_m(k_1 r) + L_2 k_1 J_m'(k_1 r) \} + B \{ \frac{j m}{r} L_1 J_m(k_2 r) + L_1 k_2 J_m'(k_2 r) \}] e^{j\phi} \quad (A-2)$$

$$E_r^P = [A \{ k_1 L_2 J_m'(k_1 r) - \frac{j m}{r} L_2 J_m(k_1 r) \} + B \{ k_2 L_1 J_m'(k_2 r) - \frac{j m}{r} L_1 J_m(k_2 r) \}] e^{j\phi}, \quad (A-3)$$

where $\phi = m\theta - k_z z + \omega t$, A and B are constants, and other notations are given in Ref. 10 in Chapter 2.

In vacuum annulus between the plasma and the antenna;

$$E_z^V = [C G_{mm}(\beta r, \beta s) + D H_{mm}(\beta r, \beta w)] e^{j\phi} \quad (A-4)$$

$$E_\theta^V = [\frac{-k_{nm}}{\beta^2 r} (C G_{mm} + D H_{mm}) - \frac{j \omega \mu_0}{\beta} (E G'_{mm} + F H'_{mm})] e^{j\phi} \quad (A-5)$$

$$E_r^V = [\frac{j k_n}{\beta^2 r} (C G'_{mm} + D H'_{mm}) - \frac{\omega \mu_0 m}{\beta^2 r} (E G_{mm}' + F H_{mm}')] e^{j\phi}, \quad (A-6)$$

where $H_{mm}(\beta r, \beta w) = I_m(\beta r)K_m(\beta w) - I_m(\beta w)K_m(\beta r)$, $G_{mm}(\beta r, \beta s) = I_m(\beta r)K_m(\beta s) - I_m(\beta s)K_m(\beta r)$, and $-\beta^2 = k_0^2 - k_n^2$.

In vacuum annulus between the antenna and the wall;

$$E_z^f = V H_{mm} e^{j\phi} \quad (A-7)$$

$$E_\theta^f = \left[-\frac{k_{nm}}{\beta^2 r} V H_{mm} - \frac{j\omega\mu_0}{\beta} W H'_{mm} \right] e^{j\phi} \quad (A-8)$$

$$E_r^f = \left[\frac{jk_n}{\beta} V H'_{mm} - \frac{\omega\mu_0 m}{\beta^2 r} W H_{mm} \right] e^{j\phi}. \quad (A-9)$$

2. DEFINITION OF MATRIX $\vec{T}$

For $\vec{A}^t = (A, B, C, D, E, F, V, W)$ and $\vec{J}^t = (J_z, J_\theta, 0, 0, 0, 0, 0, 0)$, $\vec{T}$ is defined as

$$\begin{bmatrix} 0 & 0 & -\frac{j\omega\epsilon_0}{\beta} G'_{mm}(\beta s) & -\frac{j\omega\epsilon_0}{\beta} H'_{mm}(\beta s) & \frac{k_{nm}}{\beta^2 s} G_{mm}(\beta s) & \frac{k_{nm}}{\beta^2 s} H_{mm}(\beta s) & \frac{j\omega\epsilon_0}{\beta} H'_{mm}(\beta s) & -\frac{k_{nm}}{\beta^2 s} H_{mm}(\beta s) \\ 0 & 0 & 0 & 0 & G_{mm}(\beta s) & H_{mm}(\beta s) & 0 & -H'_{mm}(\beta s) \\ 0 & 0 & -\frac{k_{nm}}{\beta^2 s} G_{mm}(\beta s) & -\frac{k_{nm}}{\beta^2 s} H_{mm}(\beta s) & \frac{j\omega\mu_0}{\beta} G'_{mm}(\beta s) & \frac{j\omega\mu_0}{\beta} H'_{mm}(\beta s) & \frac{k_{nm}}{\beta^2 s} H_{mm}(\beta s) & \frac{j\omega\mu_0}{\beta} H'_{mm}(\beta s) \\ 0 & 0 & G_{mm}(\beta s) & H_{mm}(\beta s) & 0 & 0 & -H_{mm}(\beta s) & 0 \\ J_m(k_1 p) & J_m(k_2 p) & -G_{mm}(\beta p) & -H_{mm}(\beta p) & 0 & 0 & 0 & 0 \\ h_1 J_m(k_1 p) & h_2 J_m(k_2 p) & 0 & 0 & -G_{mm}(\beta p) & -H_{mm}(\beta p) & 0 & 0 \\ F_{A4} & F_{B4} & -\frac{j k_n}{\beta} G'_{mm}(\beta p) & -\frac{j k_n}{\beta} H'_{mm}(\beta p) & \frac{\omega\mu_0 m}{\beta^2 p} G_{mm}(\beta p) & \frac{\omega\mu_0 m}{\beta^2 p} H_{mm}(\beta p) & 0 & 0 \\ F_{A3} & F_{B3} & \frac{k_{nm}}{\beta^2 p} G_{mm}(\beta p) & \frac{k_{nm}}{\beta^2 p} H_{mm}(\beta p) & \frac{j\omega\mu_0}{\beta} G'_{mm}(\beta p) & \frac{j\omega\mu_0}{\beta} H'_{mm}(\beta p) & 0 & 0 \end{bmatrix} \quad (A-10)$$

where

$$F_{A,B4} = S[k_{1,2} l_{2,1} J_m'(k_{1,2} p) - \frac{j m}{p} L_{2,1} J_m(k_{1,2} p)] \\ + j D [\frac{j m}{p} l_{2,1} J_m(k_{1,2} p) + L_{2,1} k_{1,2} J_m'(k_{1,2} p)]$$

$$F_{A,B3} = \frac{j m}{p} l_{2,1} J_m(k_{1,2} p) + L_{2,1} k_{1,2} J_m'(k_{1,2} p)$$

for $(A,B)=(1,2)$,

and J_z and J_θ are axial and azimuthal current densities in the antenna.

APPENDIX B LIST OF PUBLICATIONS

CHAPTER 2

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CHAPTER 3

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R.Itatani, Y.Yasaka, and H.Imaizumi, Research Program on Mirror Confinement in Kyoto University, in Proc. of Int. Symp. on Physics in Open Ended Fusion Systems (Tsukuba, April, 1980) p.163.

K.Adati, T.Watari, T.Aoki, S.Hidekuma, S.Hiroe, M.Ichimura, T.Kawamoto, R.Kumazawa, Y.Okubo, S.Okamura, T.Sato, C.Singh, K.Takayama, M.Sugawara, R.Itatani, and Y.Yasaka, RFC-XX Plasma Confinement and Heating Experiments, in Proc. of 8th Int. Conf. on Plasma Physics and Controlled Nuclear Fusion Research (Brussels, 1980) IAEA-CN-38/F-5.