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IN THE ATC TOKAMAK

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PROPAGATION AND ABSORPTION

BY

H. TAKAHASHI

**PLASMA PHYSICS
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ICRF HEATING AND WAVE GENERATION IN THE ATC TOKAMAK(*)

Part-I: Wave Generation, Propagation and Absorption

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ABSTRACT

Experimental results and their theoretical interpretations of rf heating experiments in the ion cyclotron range of frequencies are presented. The presentation is in two parts, with this part describing the experimental facilities and characteristics of fast compressional Alfvén waves in the tokamak, and a subsequent paper which will deal with the principal heating results. It is shown that a simple coupling structure which consists essentially of a semi-circular metal strip and a pair of bumper limiters is sufficient to efficiently couple rf power to the plasma. The major part of the coupled power is radiated as the fast wave, and there is no evidence that parasitic loading, if any, accounts for a significant amount of the power. A coupling efficiency (net coupled power/incident power) in excess of 90 % is routinely obtained. The series loading resistance is 1-2 ohms, and the wave generating efficiency (wave power/net coupled

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power) is 90-95 %. The power deposition efficiency (power deposited in plasma/ wave power) is about 90 %. Overall, about 75 % of the generator output power is deposited in the plasma. General features of the fast wave propagation, manifest as occurrences of toroidal eigenmode resonances, can be described by a theoretical model based upon a cylindrical waveguide which is partially filled with a uniform plasma. However, one-to-one correspondence of the observed and predicted resonant peaks cannot be established and the model is poor in predicting the mode onset conditions. Wave damping strength at high densities, though not directly measureable, does not appear to be grossly different from predictions of "pure" cyclotron damping theories, but at lower densities the experimentally measured damping strength greatly exceeds the theoretical value, and its dependence on the resonant layer location is asymmetric with the geometrical plasma center in contradiction to the theoretical expectation. Both of these discrepancies are attributable to the two-ion hybrid resonance between the minority proton and majority deuteron species. Theoretical calculations indicate that the hybrid resonance strongly alters the E-field polarization within the cyclotron resonant layer in such a way as to enhance both the proton fundamental cyclotron and deuteron second harmonic damping. These effects are asymmetric with respect to the plasma center, and the degree of enhancement is shown to be sufficient to account for the observed discrepancies.

I. GENERAL INTRODUCTION

One promising method of additional heating for bringing a tokamak plasma to thermonuclear temperatures is to irradiate the plasma by electromagnetic fields in the Ion Cyclotron Range of Frequencies ("ICRF" heating). The impressed radio frequency (rf) field generates fast compressional Alfvén waves in the plasma whose energy is absorbed by the ions and electrons through various field-particle interactions. Experimental investigations of ICRF heating include observation of heating at power levels high enough to affect the plasma thermal energy, and studies of wave generation and propagation at low power levels. Recent theoretical investigations center on physical mechanisms of wave absorption. A historical review of these subjects with an extensive list of publications can be found elsewhere(1).

A series of ICRF heating experiments were conducted in the ATC tokamak(2), and were the first under the conditions in which the ion heating could be critically evaluated. The principal results of the experiments, showing substantial heating of the plasma core, have already been reported in a short letter(3). The present paper in two parts documents the experiments in a more complete form. In Part-I the experimental facilities are described in some detail. It was believed that ICRF heating requires an elaborate coupling structure, including a reverse-turn induction coil

and an electrostatic shield, and might be incompatible with a thermonuclear environment. However, an antenna structure is used in the ATC experiments, which is considerably simpler than antennas used in the previous experiments. Successful use of the antenna represents a step forward towards a simpler and more rugged (presumably all-metal type) structure which could be used in a thermonuclear environment. Wave radiation, propagation and damping characteristics observed in the experiments are compared with theoretical calculations. The model used for the radiation and propagation characteristics is a straight cylindrical waveguide which is partially filled with a uniform plasma. The toroidal effects are introduced by connecting the waveguide in a ring form. The wave attenuation due to cyclotron damping is computed using the theoretical results obtained by Stix. The predicted wave damping strength and loading resistance variation are compared with the experimental results. Influence of the two-ion hybrid resonance on the cyclotron damping is analyzed in a slab geometry with an inhomogeneous plasma. Attempts are made to explain the observed discrepancies between experiment and theory in terms of enhanced cyclotron damping. The enhancement factors are calculated for specific cases. In Part-II the principal heating results and methods of evaluating the ion heating are discussed. Finally, effects of ICRF heating on the plasma, other than the ion temperature change, are described. They include

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changes in the electron temperature and density profiles and occurrence of the $m=2$ MHD oscillations.

II. EXPERIMENTAL FACILITIES

1. WAVE COUPLING SYSTEM

An artist's sketch of the wave coupling structure installed on the ATC vacuum vessel is shown in Fig. 1. The output of an rf generator placed at a remote location is piped into the experimental area by a six-inch coaxial transmission line of 75 Ohm characteristic impedance. Near the tokamak there is a segment of nine-inch line that accommodates a set of two stub tuners for impedance matching (not shown in the figure). This is connected to a six-inch line and input to the tuning capacitor. The line tapers down to a smaller size near the vacuum vessel and at the vacuum feed-through the inner diameter of the outer conductor is $3 \frac{7}{8}$ in. and the outer diameter of the inner conductor is $\frac{7}{16}$ in. The vacuum seal is accomplished by conical teflon feed-throughs and O-rings. Under the vessel the line segment is terminated by another capacitor. Both capacitors are of 100 kV rating and tunable between 10 and 250 pF. The segments of the transmission line between the stub tuners and the top feed-through and between the bottom feed-through and the terminating capacitor can be pressurized to 20 lb/in^2 . During high power experiments these segments are filled with SF_6 gas to guard against

electrical breakdown in the transmission line. Inside the vacuum vessel the center conductor is connected to the antenna and the outer conductor is electrically and mechanically attached to the vessel walls. The overall electrical system and mechanical structure exterior to the tokamak vessel are similar to those used in the earlier experiments in ST tokamak(4).

The internal structure of the ATC coupler is different from those used in previous experiments in two important respects: the absence of an electrostatic shield and the presence of limiters to protect the antenna from direct contact with the dense plasma. Absence of a shield makes the ATC structure considerably simpler to fabricate and install. Fig. 2 shows the antenna and limiters installed on a mock-up structure. Indentations of the limiters are introduced to accomodate other experiments at the same port. They and components mounted on them are extraneous to the rf coupling system. The central element in the coupling structure is an antenna located on the outer side (larger major radius side) of the plasma torus. It subtends an arc-angle of 135 degrees in the poloidal plane about the plasma center which is nominally located at a major radius of 0.9 m and has an inside radius of 0.19 m. The antenna is made of a copper strip, 1 in. wide and 1/4 in. thick, and is encased in a ceramic sheath which has no metalized inner surfaces. No special precautions are made to evacuate the interior space of the sheath. The limiters lie in poloidal

planes separated by 0.20 m along the meridian plane (disregarding the indentations) and located symmetrically on either side of the antenna. The sides of the limiters facing the antenna are covered by ceramic plates. The material used for these plates and the antenna sheath is glass ceramic with good electrical and vacuum properties, machinable with relative ease and requires no firing(5). These ceramic components are incorporated in the antenna design as a precaution against electrical breakdown, but the necessity for them is not established, because no experiments are conducted without them. The inside radius of each limiter is 0.17 m and the ceramic plate is recessed by 1/8 in. The limiters are electrically grounded to the vessel and have slits cut in the radial direction (see Fig. 2) to reduce the rf current near the plasma edge. (The return current in the limiters would tend to cancel the rf field created by the antenna current.)

The terminating capacitor is first adjusted so that the voltage minimum point is located at the center of the antenna. This ensures the lowest possible voltage at the feed-throughs and minimizes the risk of electrical breakdown in the vacuum region. This arrangement also brings a current maximum of the standing wave on the transmission line to the antenna center. In order to achieve efficient coupling in the face of relatively small radiation resistance it is important to obtain a high antenna current, which is accomplished by adjusting the tuning capacitor for

a resonance. The unloaded Q of this parallel tuned circuit is about 440 and the effective antenna inductance is estimated to be about $0.52 \mu\text{H}$. The impedance (about $15 \text{ k}\Omega$ unloaded) of this circuit is matched to the line impedance by a set of two stub tuners. The stub tuner located nearest the antenna is adjustable between 0.1 and 0.15 wavelengths and the other between 0.039 and 0.086 wavelengths. The coupling structure may be regarded as an end-fed, resonant antenna, but only a small portion near the current maximum point is used to radiate the waves.

In the presence of a plasma tuning of the matching network requires iterative adjustments of the tuning capacitor and two stub tuners and is a tedious operation. A matched condition is obtained for one set of plasma parameters and is left unchanged for other conditions. As plasma reactive effects vary, the network becomes detuned and the antenna current diminishes. However, these effects are small under the heating experiment conditions and the matching remains good. The measured root-mean-square (RMS) antenna current ranges up to 450 amp but is typically 200-300 amp. The reflected power is no more than 10 % of the incident power under all plasma heating conditions, which means that the power coupling efficiency (net coupled power/incident power) of ICRF antenna in ATC is better than 90 %.

A net rf power of up to 200 kw is coupled into ATC by this system. (The upper limit on the power is not imposed by electrical breakdown, but by plasma physics considerations. See Part-II.) The maximum peak (non-RMS) feed-through voltage is estimated to be about 6 kv. No electrical breakdown is observed during any of the experiments.

2. RF INSTRUMENTATION

The locations of various rf instrumentation are shown schematically in Fig. 3. Two 50 dB directional couplers (not shown in the figure) are mounted in the transmission line between the generator and the stub tuners to measure the incident and reflected power. The rf current in the antenna is monitored by a one-turn electrostatically shielded magnetic loop which is half an inch in diameter. The center of the loop is located $3/4$ in. behind the antenna slightly below the antenna midpoint (see Fig. 4). This current transducer has a sensitivity of about 0.1 volt/ampere measured across a 50 ohm termination placed immediately at the output connector. A voltage null probe is also located $1/2$ in. behind the antenna slightly above the antenna midpoint. The probe, similar in construction to a spherical Langmuir probe ($1/16$ in. in diameter), is primarily used to locate the voltage minimum point, but it is also used to estimate the plasma density behind the antenna. Capacitive voltage dividers are housed inside the terminating and

tuning capacitor housings to monitor rf voltage on the resonant segments of the transmission line.

Arrays of rf magnetic probes measure the wave field at different locations around the torus (Fig. 3). At a location 120 degrees around the torus from the antenna there is an array of four probes distributed over a distance of 0.61 m toroidally which determine the wavelength. At the same location there is another array of probes placed at various poloidal angles in order to determine the (minor) azimuthal mode number. All probes at this location are single-turn and electrostatically shielded ("single-turn probes"). At another location 90 degrees around the torus from the antenna, there are two probes positioned almost directly above and below the plasma axis. Each of these probes consists of two single-turn electrostatically shielded loops wound in the opposite sense and connected in the push-pull configuration ("double probes"). All magnetic probes used in the experiments are oriented to couple to the toroidal component of the rf wave field. The relative phase of two probe signals is determined either by observing them directly on an oscilloscope or by means of an analog phase comparator constructed for this purpose. The comparator amplifies and clips the input signals and then beat the resultant square pulses against each other. The heterodyned signal, after a low-pass filtering, is proportional to the phase difference. Measurements (within an ambiguity of a multiple of 180 degrees) are relatively free from the

influence of rapidly varying amplitudes.

One important measure of the antenna performance is its equivalent series loading resistance. It is defined as the ratio of the net rf power (RMS) fed into the antenna to the square of its rf current (RMS), and is determined by means of an analog circuit. Signals from the directional couplers and the current transducer are rectified and squared. The difference between the incident and reflected power signals is then divided by the current signal, and the output of the analog circuit displayed directly on an oscilloscope. This device is referred to as the loading resistance (R_g -) computer in this paper. The first-generation instrument used in these experiments has a limited dynamic range and the calibration factor is somewhat dependent on the amplitudes of the input signals, but the output should be correct within 20 %. The computer is extremely useful in following rapid variation of the loading resistance caused by toroidal eigenmodes.

3. PLASMA DIAGNOSTICS

The principal ion temperature diagnostic is a three-channel charge exchange (CX) neutral particle analyzer aimed perpendicularly at the plasma. A CX spectrum is usually constructed from 14 discharges and the temperatures are determined at 15 time points at either one or two millisecond intervals. There is also a one-channel CX

analyzer viewing the plasma in a direction tangent to the torus. Temperatures determined by these analyzers will be referred to as the perpendicular and parallel CX temperatures. The locations of these instruments are indicated in Fig. 3. These analyzers are not capable of discriminating mass differences so their signals contain contributions not only from deuterons but also from minority protons that may be incidentally present in a deuterium plasma. A vacuum ultra-violet (VUV) monochrometer located nearly opposite from the antenna is used to measure the intensity and Doppler width of impurity radiation. The instrument is equipped with a vibrating LiF plate(6) and can rapidly scan the Doppler profile to yield the temperature ("Doppler temperature") of impurity species with a time-resolution of 3 msec. A ^3He neutron counter is located on the tokamak major axis 1.4 m below the machine mid-plane. It counts at the rate of one pulse for every 1.2×10^5 neutrons produced in the plasma in the compressed state. The neutron production rate is converted to the temperature assuming a Maxwellian ion energy distribution and parabolic radial variations of the ion temperature and density. The deuteron deficiency is taken to be 0.5 which takes into consideration impurities introduced during the heating pulse. The temperature deduced from the neutron counting is referred to as the "neutron temperature" for short.

In addition to these ion temperature diagnostics there are a number of others to infer the state of the plasma. There are two bolometers placed in the shadow of the limiters: one is facing the plasma along the minor radius and the other is looking in the direction tangent to the plasma torus in the meridian plane. The radial bolometer measures the integrated power loss due to charge exchange and radiation. The tangential bolometer is designed to collect principally ions which have escaped through uncontained drift ("banana") orbits and are moving in the direction of the plasma current. Each bolometer is made of thermistors mounted on a 1 mm thick oxidized nickel film and encased in a copper housing(7). The radial bolometer has a conical viewing angle of 60 degrees and the tangential one has a much narrower angle to avoid intercepting radiation and neutral particles emitted directly from the plasma. A radially scannable electrostatic probe, made of tungsten wire 2 mm long and 1 mm in diameter, is used to determine the ion density in the limiter shadow. Two soft x-radiation detectors made of Schottky diodes obtain line-integrated signals along different chords, from which localized values of 5-10 keV radiation can be derived(8). The electron temperature and density are measured by a standard Thomson scattering method(9).

4. DISCHARGE CHARACTERISTICS

ATC is a tokamak with toroidal compression capability(2). The discharge can be positioned anywhere between the major radii of 0.9 m and 0.37 m with corresponding minor radii of 0.17 m and 0.10 m. Fast compression of the major radius results in a plasma with a higher temperature and density. Although the compression capability is utilized in the course of the experiments, all ICRF heating is applied to uncompressed deuterium plasmas with a major radius between 0.84 and 0.9 m. The rf generator frequency is fixed at 25 MHz which corresponds to the second harmonic deuteron cyclotron frequency at 1.64 T. The central electron density, n_{e0} , is deduced from the microwave interferometer measurements assuming a parabolic radial profile. All attempts to heat the plasma are performed at moderate-to-high densities ($n_{e0} > 2 \times 10^{19} / \text{m}^3$) and with the cyclotron resonant layer placed near the plasma center. In this paper we refer to these conditions as "plasma heating conditions". For purposes other than heating both the field and density are varied over much wider ranges: $n_{e0} = 0.5 - 4.2 \times 10^{19} / \text{m}^3$ and $B_T = 1.2 - 1.9$ T. Some pertinent discharge characteristics are described in the remainder of this section.

The discharge typically lasts 40 ms with the rf pulse usually applied 20-25 ms after the discharge initiation lasting 10 ms. This leaves sufficient time after the pulse

to observe decay of the ion temperature. The pulse length is comparable or longer than the ion energy replacement time and is much longer than the proton-on-deuteron equipartition time. The latter inequality is important for obtaining reliable assessment of plasma heating: if the pulse length is shorter than the equipartition time, protons may be at a significantly higher temperature and the heating of a small fraction of the ions may be interpreted as the heating of the entire ion species. If the pulse is longer than the ion energy replacement time, it becomes possible to detect effects of the heating pulse on the energy confinement. The ATC experiments are the first ICRF heating experiments conducted under these conditions that permit reliable measurements of the ion heating.

The discharge is initiated in two different ways. In one way the vacuum vessel is filled with the working gas to the pressure needed to achieve a desired density level. In the other the gas is injected into the pump line or vessel shortly before the discharge initiation and the injection rate is adjusted to achieve a desired density level. In some discharges a part of the vessel interior surface is gettered by evaporating titanium vapor between discharges, which is known to reduce impurities in the plasma(10). When the gettering is used, the gas injection is the only possible way of initiating the discharge. The gettered discharges are sensitive to the timing of gas injection and are less reproducible than the ungettered discharges.

However, this mode of operation is employed for most of the higher power heating experiments, because it is found effective in retarding excitation of $m=2$ MHD oscillations. In ungettered discharges the density generally rises during most of the discharge duration, whereas in gettered discharges it falls rapidly after peaking at an early phase of the discharge. The top photographs in Fig. 5 are the microwave interferograms (taken vertically at a major radius of 0.9 m) for typical ungettered discharges with and without rf heating. The rf pulse starts at 20 ms and lasts for 10 ms.

The lower oscillograms in Fig. 5 show the corresponding one-turn loop voltage and plasma current. Because of the short discharge duration there is no period over which the plasma properties and current stay steady. Under these circumstances the loop voltage is not due solely to plasma resistance and Ohmic heating power (and other derived quantities such as the electron energy confinement time and effective- Z) cannot be determined reliably. In the presence of rf heating the loop voltage rises somewhat above the no-rf value during the pulse and dips below it after the termination of the pulse. This behavior is indicative of inductive effects and the voltage rise does not necessarily imply a change in Ohmic heating power. The plasma current is little affected by the rf pulse in high-current, high-density ("robust") discharges, but falls somewhat (up to about 5 %) in lower-current, lower-density discharges.

The plasma position is hardly changed by the pulse in robust discharges, while less robust discharges tend to be pushed radially outward. The extent of excursion is dependent on the rf power and plasma conditions, but is usually of the order of 1 cm or less.

In ungettered discharges it is usually possible to obtain steady reproducible plasma conditions. However, repeated application of rf power appears to cause gradual but secular changes in the conditions: both the density and temperature tend to lower values. (The symptoms appear similar to those under extended discharge cleaning.) In order to see effects of rf heating comparisons are often made between two consecutive discharges with and without an rf pulse. However, the initial portion of the discharge before the rf pulse initiation is often found different from the corresponding part of the discharge without rf heating. The discharge apparently possesses memory of the previous shot.

The ATC plasma contains a variety of impurity ion species. The most important among them for the considerations of wave absorption is hydrogen. An accurate measurement of its concentration is difficult, but the proton-to-deuteron ratio is estimated to be at most 5 % (11). This upper bound is obtained from the fact that no recognizable peak exists in the wing of the D_{β} spectral line at a location where the H_{β} line is expected.

Throughout this paper the proton concentration is taken to be small (<5 %) and is arbitrarily assumed to be 3 %, whenever a numerical value is needed. Other dominant light impurities are oxygen and carbon, whose concentrations depend upon the cleanliness of the vacuum vessel walls. In one set of measurements(11) the total (summed over all ionization states) carbon and oxygen densities are found to be $2.7-4.5 \times 10^{17} / \text{m}^3$ and $2.2-3.9 \times 10^{17} / \text{m}^3$ respectively. The central electron density range is $2.0-2.4 \times 10^{19} / \text{m}^3$ in these measurements. In addition, the ATC plasma contains a quantity of helium ions, which originate from the stripping cell of the perpendicular charge exchange analyzer. A rough estimate of the lower bound of their concentration is about 5 % which is based upon the measured pressure rise in the vacuum vessel when the analyzer valve is opened. There is no evidence that the ceramic material used in the coupling structure contaminated the plasma.

III. WAVE GENERATION, PROPAGATION AND ABSORPTION

1. INTRODUCTION

The earlier rf heating experiments in C-Stellarator utilized the slow branch of the low frequency cold plasma wave (also called the torsional Alfvén wave or the ion cyclotron wave) in magnetic beach configuration. In tokamaks the fast branch (the compressional Alfvén wave or the fast wave) is used at the fundamental and second

harmonic of the ion cyclotron frequency ($\Omega = \omega / \omega_{ci} = 1$ or 2) without magnetic beaches. In both cases one possible mechanism for dissipating wave energy is ion cyclotron damping. However, in the inhomogeneous B-field of a tokamak the cyclotron resonance conditions ($\Omega = 1$ or 2) are met only on a roughly cylindrical surface of a constant major radius (R_c). Finite ion thermal speed (v_i) parallel to the B-field makes the resonant interaction region ("ion cyclotron resonant layer") a layer of finite thickness given approximately by $2R_c v_i k_{\parallel} / \omega$. In one-ion species plasmas the fundamental and harmonic cyclotron resonances occur at different physical locations for a given B-field and wave frequency. It is often possible to place only one resonant layer inside the plasma and to study the associated damping mechanism independent of the others. Deuteron-proton mixtures used in recent tokamak experiments represent a more complicated situation; the proton fundamental cyclotron and the deuteron second harmonic cyclotron resonances occur at the same location, and, furthermore, a new resonance between these ion species ("two-ion hybrid resonance") comes into existence. Because of the proximity of the cyclotron and hybrid resonant layers mutual influence of these resonances plays an important role in determining dominant damping processes and division of absorbed power among different species.

The fast wave in one-ion species (say, deuterium) plasmas has predominantly right-handed (with respect to the confining B-field) polarization and its propagation characteristics are little affected by the left-handed ion motion at the second harmonic cyclotron resonance. However, the wave damping is governed by the small left-handed component (E_+) of the wave E-field in the resonant layer. Since the wave field varies twice as fast as the cyclotron period, resonant ions experience acceleration during one half of their Larmor period and deceleration during the other half. The two effects are, in general, unequal and there is a net energy transfer from the wave to the ions. The magnitude of the difference is determined by the ion Larmor radius (ρ_L) and the perpendicular gradient of the E_+ -field. The latter can be approximated by $k_\perp |E_+|$ in one-ion species plasmas, so the second harmonic damping strength is proportional to $1/2 \rho_L (k_\perp |E_+|)^2$.

Addition of a small quantity of protons in a predominantly deuterium plasma introduces two new resonances: the proton fundamental cyclotron resonance and the two-ion hybrid resonance. Both of these resonances give rise to additional damping mechanisms. The protons experience the electric field at their fundamental resonant frequency and keep accelerating in a familiar spiraling motion. If there are too many protons, their coherent motion tend to shield out the E_+ -field from the resonant layer and wave damping is consequently weak. When the

proton concentration is a few percent damping by this mechanism (commonly called "minority species damping") can be much stronger than the second harmonic damping. However, the relative damping strengths depend strongly on the parallel wave phase speed. The two-ion hybrid resonant surface lies very close to the cyclotron surface and has two principal effects on the waves: alteration of the wave E-field polarization in the ion cyclotron resonant layer and introduction of mode conversion near the hybrid resonant surface. Mode conversion is important whenever cyclotron damping at the hybrid surface is weak. This situation arises either when the proton concentration is large and the hybrid resonant surface lies well outside the cyclotron resonant layer, or when the waves are nearly perpendicularly incident on the hybrid surface. Alteration of the polarization caused by the hybrid resonance is important when the proton concentration is a few percent. Significant increase in the magnitude and gradient of the $E_{\perp}$ -field on one side of the ion cyclotron resonant layer closer to the hybrid surface results in strong enhancement of the deuteron second harmonic damping. The minority species damping is also significantly greater than in the absence of the hybrid resonance. Understanding of the damping mechanisms is important for determination not only of the plasma heating rate but also of the power division among different plasma species: the minority species damping deposits the wave energy into protons and the deuterons are heated only

indirectly through proton-deuteron collisional energy transfer, while mode conversion may also heat electrons through Landau damping of the mode-converted electrostatic waves and the second harmonic damping heats deuterons directly.

Adam and Samain(12) calculated in their pioneering work the second harmonic cyclotron damping strength in one-ion species plasmas and the fundamental cyclotron damping by the minority species in two-ion species plasmas. Stix(13), who analyzed temporal evolution of the ion distribution function under fast wave heating, obtained similar results for the damping strengths. The roles played by mode conversion at the two-ion hybrid resonance were analyzed by Swanson(14) and Perkins(15). Importance of the E-field polarization modified by the two-ion hybrid resonance in calculating the fundamental and second harmonic damping in two-ion species plasmas was recognized by Takahashi(1,16). More extensive historical review of the subject can be found elsewhere(1).

In this paper the phrase, "proton fundamental cyclotron", is abbreviated as PFC as in PFC resonance or PFC damping. The phrase, "deuteron second harmonic", is shortened as DSH as in DSH resonance or DSH heating. The phrase, "two-ion hybrid", is represented by IHB as in IHB mode conversion.

2. FAST WAVE PROPAGATION IN ATC

In ATC the wave damping length is large compared to the (minor) radial extent of the plasma. This makes the wave propagation and damping amenable to separate analyses: propagation characteristics and field structure of undamped waves are first obtained and the damping strengths are then calculated based upon the known field structure. We first consider the fast waves in a straight cylindrical waveguide partially filled with a cold uniform plasma and immersed in a uniform longitudinal magnetic field. Eigenmode solutions for such a system were obtained by Bernstein and Trehan(17) for general non-axisymmetric modes. A theory of excitation of azimuthally symmetric eigenmodes was developed by Stix(18). The theory(19) used here is an extension of the latter to more general non-axisymmetric cases using the eigen-solutions of the former. The second-order dispersion relation for the fast waves is solved numerically together with the boundary conditions to obtain eigenvalues for the wave numbers. Analytic expressions for the wave fields are then obtained in terms of Bessel functions. Other related quantities such as the wave energy density, Poynting flux and antenna radiation resistance are analytically derived from these expressions. (The term, "radiation resistance", is used for the antenna loading resistance in a straight infinitely long waveguide.) In particular, the group velocity averaged over the waveguide cross-section is computed as the ratio of the time-averaged Poynting flux

integrated over the waveguide cross-section to the wave energy within a unit length of the waveguide. Both the Poynting flux and wave energy consist of contributions from the plasma and vacuum regions within the waveguide. The relative magnitudes of the contributions depend not only on the size of the vacuum region, but also on the field structure and hence on the individual modes.

The ATC vacuum vessel is oblong in cross-section (see Fig. 4) and is difficult to model accurately in theoretical calculations. It encloses a large vacuum space on the small major radius side. When a plasma of minor radius 0.17 m is located at a major radius of 0.9 m the outer (larger major radius side) walls in the immediate vicinity of the antenna are roughly semi-circular with a radius of 0.21 m and are concentric with the plasma. At other locations the radius is 0.20 m. On one hand, the relative locations of the plasma, antenna and walls are important for calculations of the antenna radiation resistance. On the other hand, the presence of a large vacuum space has important consequences in interpreting experimental results (see below). In theoretical calculations the vacuum vessel is modeled by a torus of circular cross-section with a minor radius of 0.21 m. Were the theoretical vacuum vessel to enclose the same volume as the actual vessel, its equivalent radius would be 0.27 m. The theory is also incomplete because of its failure to account for the presence of a large port and bumper limiters in the vicinity of the antenna, which are

expected to have influence on the radiation resistance. 80 % of the actual vessel walls are made of bellows structure, but the theoretical model assumes smooth stainless-steel walls. The actual inhomogeneous plasma is represented by one having a constant density that equals the volume average of the actual density.

The fast wave dispersion relation can be conveniently represented in terms of the perpendicular ($N_{\perp}$) and parallel ($N_{\parallel}$) components of the "Alfvén refractive index(13)" ($\vec{N} = \vec{k}V_A/\omega$) as shown by the solid curve in Fig. 6. The curve is nearly of elliptic shape and can be closely approximated by the following formula:

$$N_{\perp}^2 + (1 + \Omega) N_{\parallel}^2 = 1 \quad (1)$$

This "elliptic" approximation(*) is also shown in the figure by the broken line.

The boundary conditions require discrete solutions of the dispersion relation in the waveguide. Examples of these waveguide modes are shown for a set of ATC conditions in Fig. 6. A pair of radial(l) and minor azimuthal(m)--or

(*)An approximation in which the term involving Ω is omitted ("circular" or isotropic approximation) is sometimes used. This would result in a group velocity which is too small by a factor between 3 and 4.5 (set terms involving Ω in Eq. 7 to zero).

poloidal--mode numbers, (l,m) , is associated with each waveguide solution. Note that the shortest possible wavelength in the perpendicular direction is the "Alfven wavelength" ($\lambda_A = v_A/f$, where f is the wave frequency). At sufficiently low densities (or large B_T) the fast waves in a completely filled waveguide suffer a cut-off ($k_{||}^2 < 0$) and are unable to propagate along the waveguide. The fields of a cut-off mode are evanescent and confined in the vicinity of the antenna. In a partially filled waveguide the cut-off exists for the $m=0$ and all negative- m azimuthal modes (for the field variation of the form $\exp i(k_{||}z + m\theta - \omega t)$). However, the lowest radial modes of all positive- m modes propagate at very low densities(*). (They can presumably propagate at all densities.) As the density increases, a cut-off mode may become propagating, at which point $N_{\perp}=1$ and $N_{||}=0$. We term this condition the "onset" of a new propagating mode. As the density increases further, the point representing this mode moves from the upper left-hand corner toward the lower right-hand corner along the solid curve in Fig. 6. The parallel wavelength becomes progressively shorter over the "life span" of the mode. Because of the different dependence of the PFC, DSH and IHB damping strengths on the parallel wavelength, the mode goes

(*)This was demonstrated by numerical examples by Paoloni(20) for $m=+1$ mode. Similar calculations by the present author indicate that the same is true for all positive- m modes.

through different phases in each of which a different mechanism is responsible for the wave damping. Specifying the wavelength is crucial to any statement made on the relative importance of any of these damping mechanisms. It should also be noted that the mode carries more rf power (larger loading resistance) near the onset than away from it. The exact plasma conditions under which the onset of a mode takes place are sensitively dependent on the radial density profile(21) and the presence of the vacuum space. Because of the approximate nature of the representation of the actual ATC vacuum vessel and the uniform plasma model used in the present theory, the model's prediction of the onset conditions is not expected to be accurate. The modes with parallel wavelengths shorter than $(1+\Omega)^{1/2} \lambda_A$ are evanescent in the perpendicular direction: their fields are described by the Bessel functions of an imaginary argument and are peaked at the plasma edges. These surface modes and the evanescent antenna near field are not included in the present calculations.

Toroidal effects on wave propagation are accounted for only in the sense that the tokamak vacuum vessel can be regarded as a plasma-filled waveguide connected in the form of a toroidal cavity. (The B-field inhomogeneity across the plasma cross-section is not considered in computing the wave propagation characteristics.) Waves generated by the antenna return to the antenna location each time they circumnavigate the torus and their fields produce interference with the

original antenna fields. The resultant fields depend upon the phase shift and damping the waves suffer as they go around the torus. Under such conditions processes of wave generation, propagation and damping are closely coupled together and isolated investigation of each process is difficult in practice. Nevertheless, studies of the loading resistance and wave amplitude under different plasma conditions provide evidence for wave generation and clue for the physical mechanisms of wave damping.

Suppose that an antenna carrying unit rf current generates a wave field of amplitude, s , in the absence of the toroidal effect. In a toroidal geometry with dissipation the waves experience a complex phase shift $ik_{\parallel}L$ each time they go around the torus, where L is the toroidal circumference. The total field for an rf current of I for the waves propagating around the torus many times is obtained by summing an infinite series,

$$\begin{aligned} S &= sI(1 + \exp(ik_{\parallel}L) + \exp(2ik_{\parallel}L) + \dots) \\ &= sI/(1 - \exp(ik_{\parallel}L)) \end{aligned} \quad (2)$$

When the plasma conditions are appropriate for a constructive interference of the wave fields, the cavity becomes resonant. As the conditions (principally the density) vary with time, a series of such toroidal cavity resonances should be observed. A toroidal resonance is identified by a set, (l,m,n) , of radial(l), poloidal(m) and toroidal(n) mode numbers. At a resonance and at an

anti-resonance (n is a half-integer) S is pure real and is given respectively by,

$$S_R = sI_R/(1 - \alpha) \quad (3)$$

$$S_A = sI_A/(1 + \alpha) \quad (4)$$

where $\alpha = \exp(-k_i L)$ and k_i is the imaginary part of the parallel wave number. I_R and I_A are the rf currents at the resonance and anti-resonance, respectively. These two quantities are experimentally measureable. It should be noted that the toroidal resonance can be a travelling (as opposed to a more usual standing) wave resonance: the waves travelling in only one direction can cause a resonance. Note that the spatial variation of the amplitude around the torus is determined by the imaginary part of the parallel wave number. The wave amplitude decreases monotonically from the antenna location to the opposite side of the torus rather than going through a series of nodes and loops. A travelling wave resonance has some practical consequences. Its wavelength can be determined simply by measuring the relative phases of the wave field at two or more points around the torus. Travelling waves also deposit energy more uniformly around the torus than standing waves. If the so-called mode splitting, caused by unequal influence of the poloidal B-field on the left- and right-running waves, is indeed present, all experimentally observed resonances, other than the $m=0$ modes with no mode splitting, must be travelling wave resonances. (No experimental demonstration

of this fact has been reported in the literature.)

The oscillogram in Fig. 7 shows time variation of the wave amplitude (rf envelope) measured by the "double probes" located nearly directly above(T) and below(B) the plasma axis. The different amplitudes of the two signals are caused by a slightly off-centered plasma. Similar signals are obtained from the "single-turn probes". These observations provide direct evidence of wave generation in the tokamak plasma. The peaks and troughs of the wave amplitude evidently correspond to the occurrence of toroidal resonances and anti-resonances. The theoretical calculations of the wave amplitudes for the experimentally observed density variations confirm the presence of the peaks and troughs. However, one-to-one correspondence of these peaks cannot be established in the ATC experiments, because of the approximate nature of the theoretical model. It was shown in the earlier experiments(4) in the ST tokamak, however, that the density and B-field dependences of the loading resistance peaks closely followed the fast wave dispersion relation. For these experiments the circular waveguide model was a better representation of the actual vacuum vessel and the theoretical calculations accounted for the density inhomogeneity.

3. WAVE DAMPING MEASUREMENTS

A set of R_{γ} -computer signals are shown in Fig. 8. The toroidal B-field is different for each of these oscillograms. (The plasma conditions are also slightly different.) In each case the density increases monotonically by a factor two or less over the pulse period, except the left-most oscillogram for which the density reaches a maximum during the pulse. The locations of the DSH (deuteron second harmonic) resonance surface are indicated in the figure. The data are taken at density levels sufficiently low ($n_{e0} = 7.5-9.6 \times 10^{18} / \text{m}^3$ at the mid-point of the pulse) so that the $m=0$ and all negative m -number modes are cut-off. The prominent peaks in these oscillograms are determined to be the $m=+1$ mode by phase measurements of magnetic probe signals. (However, small-amplitude, higher azimuthal modes ($m > +1$) may also be simultaneously present in the signals.) When the DSH resonant surface is within the plasma, the toroidal resonance is broader, indicating the presence of damping processes on or near the resonant surface.

Quantitative evaluation of the damping strength can be made from magnetic probe signals such as those shown in Fig. 7. The quantity, s , in Eqs. 3 and 4, which is proportional to the radiation resistance, is, in general, not directly measureable. However, this can be eliminated by taking the ratio of the wave amplitude at a resonance to

that at an adjacent anti-resonance, if the variation of the radiation resistance for a small change in plasma parameters is neglected. The ratio is related to the wave damping length (L_D) through (22),

$$L_D/L = [\ln(\Upsilon + 1)/(\Upsilon - 1)]^{-1} \quad (5)$$

where $\Upsilon = I_A S_R / I_R S_A$. The normalized damping length (L_D/L) is the number of times the waves go around the torus before their amplitudes e-fold. The wave damping length is related to the resonant cavity quality factor (Q) through,

$$L_D \approx 1/k_i = 2Qu_3/\omega \quad (6)$$

where u_3 is the parallel group velocity. It should be noted that the group velocity varies widely depending upon the mode and that a large Q mode does not necessarily imply a large L_D mode. An approximation to the group velocity within the plasma can be obtained from Eq. 1,

$$u_3 = N_{||} V_A (1 + \Omega) / (1 - \Omega N_{||}^2 / 2) \quad (7)$$

The average group velocity in a partially filled waveguide is greater than this value because of the vacuum contributions. This can be a source of ambiguity, if either L_D or Q is measured and the other is computed using Eq. 6. This is especially true for the ATC experiments, because the circular cylindrical model does not properly account for the large vacuum space. It should also be noted that an error in computing u_3 (see footnote to Eq. 1) would result in

values of L_D too small by a large factor, if Q is experimentally determined and L_D calculated from it.

The damping length determined using Eq. 5 is shown(23) in Fig. 9(a) for the conditions corresponding to the oscillograms shown in Fig. 8. The abscissa is the major radius location of the DSH resonance surface. The observed modes are the $m=+1$ azimuthal modes, but may include different toroidal mode numbers. When the resonant surface is outside the plasma on the high field side ($R(\Omega=2)=70$ cm), the damping is weak and the waves go around the torus some 13 times before their amplitudes e-fold. This damping is interpreted to be caused by vessel wall resistance. The measured damping length is smaller by a factor of 20 than theoretical wall resistive damping length computed for these modes in a circular cylindrical waveguide made of perfectly smooth stainless-steel walls. The observed damping length appears reasonable, however, because of the increased path lengths of the wall currents due to (1)the elongated vacuum vessel cross-section (a 1.6-fold increase in poloidal circumference), (2)the partial bellows structure (an average of five-fold increase in toroidal circumference), (3)the presence of ports and other obstructions to the wall current paths and (4)the surface roughness ($0.8-13.0 \mu m$)(24) comparable to the skin depth ($8.5 \mu m$). In the following calculations the wall damping length is taken to be 20 times shorter than the value for the perfect cylindrical surface. When the resonant layer is

outside the plasma on the low field side ($R(\Omega = 2) = 110$ cm), the damping is stronger than on the high field side indicating the presence of some additional damping mechanisms. Nature of these damping mechanisms are not well understood. When the resonant layer is within the plasma, the normalized damping length is reduced to about 2. The ratio, $13/2 = 6.5$, of the damping lengths with and without the resonance is a measure of the effectiveness of energy deposition by the fast wave: about 90 % ($6.5/(1+6.5) = 0.87$) of the power is deposited within the plasma and about 10 % in the walls. These observations lead to an optimistic view for fast wave heating of tokamak plasmas.

The widths of the resonant peaks are also related to the cavity Q-factors, but the quantitative determination may be complicated because the resonant curves are traced out by varying electron density rather than the wave frequency as is customary in Q measurements. However, if the waveguide is completely filled with a uniform plasma, these two procedures can be shown to be approximately equivalent. Q can then be calculated by,

$$Q \equiv \frac{\omega_0}{\Delta\omega} = \frac{\omega_{pi0}}{\Delta\omega_{pi}} = \frac{2n_{e0}}{\Delta n_e} \quad (8)$$

where ω_0 , ω_{pi0} and n_{e0} are the wave angular frequency, ion plasma frequency and electron density at the resonance. $\Delta\omega$, $\Delta\omega_{pi}$ and Δn_e are the differences in these quantities at the two half-power points. This equivalence exists in

the case of the completely filled waveguide because dielectric properties of the uniform plasma alone determine the parallel wavelength and ω and $1/\omega_{pi}$ play approximately the same role. The situation is different in the case of a partially filled waveguide. A change in ω_{pi} affects only the plasma region, while a change in ω affects both the plasma and vacuum regions. In the limit of a very small plasma a large change in the plasma density is required to produce an appreciable change in the parallel wavelength and the resonant peaks traced out by the density variation would appear to have a Q value much lower than the true value. This effect would be significant in the case of a large vacuum space or a low plasma density. The ratio of the total vacuum vessel volume to the plasma volume is 2.5 for ATC and 1.7 for TFR400. The large discrepancies reported earlier for these devices(1,25) between the Q-factors determined from Eq. 8 and theoretical predictions arose, in part, from the inaccuracy of Eq. 8 because of large vacuum spaces. Determination of the Q-factor by varying the wave frequency may be more desirable under these circumstances. In this case, however, the location of the resonant surface moves as the frequency is swept. Since the fractional frequency variation corresponding to the change of the resonant surface from one plasma edge to the other is approximately twice the inverse aspect ratio (ϵ), Q must be at least as great as $10/\epsilon$ in order to obtain a spatial resolution of 5 % (of the plasma diameter) or better.

A notable feature of the damping curve in Fig. 9(a) is that the region of maximum damping is not symmetric with respect to the plasma axis at $R=87$ cm, and is shifted toward the larger major radius side. This shift was also observed in TFR experiments(25). Another manifestation of the asymmetry(26) is shown in Fig. 9(b). The abscissa is identical to that of Fig. 9(a). The ordinate is the output voltage of the bolometer facing the plasma and is proportional to the power loss due to charge exchange and radiation integrated over the period of rf pulse. The net difference between discharges with and without the heating pulse is plotted. A broad peak is located on the larger major radius side of the plasma axis where the center of the strong damping region is also located. Under these conditions the rf power is evidently more efficiently absorbed by the ions, some of which become lost through charge exchange and are intercepted by the bolometer.

Another set of output signals(23) from the R_g -computer is shown in Fig. 10. The resonant surface ($R(\Omega=2)=88$ cm) is located in the vicinity of the plasma center ($R_0=87-89$ cm) for all cases. The electron density levels are different for each case and are shown in the figure. The lowest density case (#1) is taken under conditions nominally identical to those of the center oscillogram of Fig. 8. Toroidal resonance peaks of the loading resistance become progressively broader at higher densities and are difficult to recognize at the highest density level. The

latter is a typical condition under which most heating experiments are undertaken. In the absence of isolated resonant peaks in these high density cases experimental determination of the damping strength based upon Eq. 5 is not possible.

4. CYCLOTRON DAMPING

We first compare the experimentally observed time variations of the loading resistance and wave damping strength with predictions based upon the theory of cyclotron damping. One low density case, for which direct experimental measurements of the damping strength are made, and one high density case, for which the measurements are not available, are examined in some detail. The electron Landau damping and transit time magnetic pumping (TTMP) (*) are also included in the calculations, but their effects on the fast waves are weak under the ATC conditions. The damping strengths needed in Eq. 2-4 are derived from Stix's results(13) given in terms of the quality factor of a resonant cavity. They are for the minority (Q_I), second harmonic (Q_{II}) and Landau/TTMP (Q_e) damping given respectively by,

$$Q_I = 2 \left(\frac{a}{R_0} \right) \left(\frac{n_D}{n_H} \right) \left[1 + \frac{9\pi}{16} \left(\frac{n_H}{n_D} \frac{u_{II}}{v_H} \right)^2 \right] \quad (9)$$

(*)Stix(13) showed that these two effects are coupled and must be treated together.

$$Q_{II} = \frac{2\left(\frac{a}{R_0}\right)}{(k_1 \rho_D)^2 \left[1 - \left(\frac{r_{res}}{a}\right)^2\right]^{1/2}} \quad (10)$$

$$Q_e = \frac{2}{\pi^{1/2} \beta_e \left(\frac{u_{||}}{v_e}\right) \exp\left(-\left(\frac{u_{||}}{v_e}\right)^2\right)} \quad (11)$$

Here, a and r_{res} are respectively the plasma minor radius and the distance between the cyclotron resonant layer and the plasma axis. n_H and n_D are respectively the proton and deuteron densities. $u_{||}$, v_H and v_D are respectively the wave parallel phase speed, proton and deuteron thermal speeds. β_e is the ratio of the electron thermal pressure to the magnetic pressure.

In deriving these results Stix accounted for the reduction of the E_+ -field at the cyclotron surface due to the presence of resonant minority protons. However, E_+ and its gradient were evaluated at the resonant surface and were taken to be constant across the thin resonant layer. These are reasonable approximations for one-ion species plasmas, but neither enhancement of cyclotron damping, nor mode conversion, both caused by the presence of the two-ion hybrid resonance, were included. Further approximations introduced in the analysis were the assumptions that E_+ was also constant along the resonant surface and that $N_{\perp}^2 \approx 1$. The latter approximation may result in an over-estimate of the second harmonic damping strength by a factor up to four

for those modes that are not near their onset.

It is sometimes stated that minority damping is much stronger than the second harmonic damping when the minority concentration is a few percent. However, the strong wavelength dependence of $Q_{\perp}$ in Eq. 9 should be noted. As the wavelength of a new propagating mode becomes progressively shorter with rising plasma density, the dominant damping mechanism rapidly changes from the DSH (Deuteron Second Harmonic) to PFC (Proton Fundamental Cyclotron) damping. Higher radial modes with a large $k_{\perp}$ and small $k_{\parallel}$ are more strongly damped by the DSH damping. Near the onset both damping strengths are proportional to the ion temperature. Away from the onset and at sufficiently high temperatures $Q_{\perp}$ becomes independent of temperature, which may have interesting consequences because the DSH damping becomes dominant for all wavelengths in high temperature, reactor-like plasmas. Note the symmetry of the DSH damping with respect to the relative locations of the resonant layer and plasma axis (Eq. 10). This is in contrast to the experimental observations discussed earlier (Fig. 9(a)).

Theoretical loading resistance variation for the low density case is shown in Fig. 11(a) and should be compared with the oscillogram at the center of Fig. 8. The radiation resistance is first computed based upon the theoretical model outlined in Sec. 2, and the toroidal effect is introduced through Eq. 2. The imaginary part of the

parallel wave number needed in this equation is obtained from Q-factors given in Eq. 9-11 by converting them to the damping length through Eq. 6. The average group velocity is computed assuming a waveguide radius of 0.21 m, but the true average group velocity is expected to be greater, because of the large vacuum space in ATC. An average density equal to one half the measured central density is used and the proton concentration is assumed to be 3 %. Ion and electron temperature variations are taken from a typical discharge in this density range. There are only two poloidal modes ($m=+1$ and $+2$) that carry appreciable power. The loading resistances of the individual modes up to $m=+3$ as well as the total resistance are shown in the figure. The damping lengths predicted from theory are too large, so that the theoretical toroidal resonant peaks which can reach 6.3 ohms are much larger than the largest experimental values of 0.5 ohms. Note the ambiguities introduced by converting the Q-factors to the damping length: the discrepancies caused by too large Q-factors predicted by the theory are partly offset by a too small group velocity used in the conversion. Without this offsetting effect the difference between experiment and theory would be bigger.

The theoretical values of the minority (I), second harmonic (II), electron Landau/TTMP (E), wall resistive (W) and total (T) damping lengths are tabulated in Table 1 for each toroidal mode shown in Fig. 11(a). The figures for the wall damping length are 20 times smaller than the value

computed for perfectly smooth walls. The damping lengths predicted by this theory are comparable for the PFC and DSH damping for these toroidal modes, and are much greater than the wall resistive damping length. If these predictions were true, then most of the power would be dissipated in the walls and there would be no significant change in the measured damping length as the resonant layer moved into the plasma. The previous section shows that this is not the case: the measured damping length is smaller than the tabulated values by factors between 10 and 120, depending upon the modes. The electron Landau/TMP damping is completely negligible ($L_D/L > 10^4$).

Fig. 12(a) shows the computed loading resistance variation for the high density case which should be compared with the top right oscillogram (#5) of Fig. 10. The proton concentration is again 3 %. There are five propagating modes that carry appreciable power (the first radial modes of $m=-1$ through $m=+2$ and the second radial mode of $m=+1$). The loading resistance of each mode and the sum over all modes are shown in the figure. (The $m=+1$ curve is the sum of its first and second radial modes.) The theoretical radiation resistance of a mode is, in general, highest at its onset, and is further enhanced by the fact that the long wavelength $n=1$ toroidal mode comes into resonance immediately following the onset. The first radial mode of $m=-1$ poloidal mode and the second radial mode of $m=+1$ mode have the onset at nearly identical times (at about 23 ms).

The $n=1$ toroidal modes of these poloidal modes create a high peak in the loading resistance. Such conspicuous peaks, characterized by the asymmetry in their shape, are not observed in present experiments. This discrepancy is considered to be due to the deficiencies of the uniform plasma model. When the density inhomogeneity is taken into consideration in the theory, the onset of each of these modes may take place less abruptly, and also at different times, and the peak may appear less conspicuous. (Another possible explanation for this discrepancy is that some damping mechanisms which are especially effective at the onset may be present. Both the IHB mode conversion and enhanced DSH damping are possibilities.) The other resonant peaks in Fig. 12(a) are more prominent than the experimentally observed ones, indicating actual damping strengths stronger than the theoretical predictions. However, the differences are not as pronounced as in the low density case (Fig. 11(a)).

Theoretical values of the minority (I), second harmonic (II), electron Landau/TTMP (E), wall resistive (W) and total (T) damping strengths are tabulated in Table II for each toroidal mode shown in Fig. 12(a). The figures for the wall damping length are 20 times smaller than the values computed for perfectly smooth walls. The DSH damping is the strongest mechanism for all low toroidal number modes ($n < 3$). The $m=-1$ poloidal modes and the $l=2$ radial modes appear only at high densities and are much more strongly

absorbed by the DSH damping than the PFC damping. These modes would deposit most of their energy directly into deuterons and make fast wave heating less susceptible to "thermal runaway" of the minority species(28). Unlike the low density case the theoretically predicted total damping is now dominated by power dissipation in the plasma rather than in the walls. The theoretical power deposition efficiency ranges from 74 % to 97 % for these modes. The efficiency is generally better for those modes that carry appreciable power. Although no direct measurements are available, comparisons of the theoretical loading curve with the experimental one indicate that the computed damping strengths are still an underestimate of the actual values. The actual efficiency is therefore expected to be substantially better than these figures. Electron Landau/TTMP damping is completely negligible ($L_D/L > 10^4$) except where entries are made in the table.

The theory in this section is based upon a uniform plasma model with the cyclotron and Landau/TTMP damping strengths obtained by Stix. The uniform plasma model is inaccurate in predicting the mode onset, and one-to-one correspondence of the toroidal resonant peaks cannot be obtained. This discussion has shown that the predicted damping strengths at low densities ($n_{e0} = 0.5 - 1.0 \times 10^{19} / \text{m}^3$), where the experimental measurements for individual modes can be made, are too small by a large factor (18-120). The dependence of the damping strength on the resonant layer

location also fails to predict the observed asymmetry about the plasma center. Although the damping strength of individual modes cannot be measured at higher densities, where many overlapping modes are present, the predicted time variation of the loading resistance indicates that the discrepancies are much smaller.

5. ENHANCED CYCLOTRON DAMPING

We next consider whether better agreement can be found between experiment and theory, if enhancement of the cyclotron damping caused by the presence of the two-ion hybrid resonance is included. In Fig. 13 the geometry of the two-ion hybrid and cyclotron resonance surfaces are shown. The plasma parameters are chosen to correspond to the typical ATC plasma heating experiments, but the proton concentration is assumed to be much larger (10 %) than the estimated actual value for clearer graphical representation of the surfaces. The radial density variation is assumed parabolic and the B-field is taken to be inversely proportional to the major radius. In this inhomogeneous plasma model the fast wave propagation ($n_{\perp}^2 > 0$) and evanescence ($n_{\perp}^2 < 0$) depend on the local plasma properties. The surface designated by $F=0$ is the "fast wave cut-off surface", outside of which the fast wave is evanescent. Within small crescent-shaped areas inbedded in the evanescent region the slow (ion cyclotron) waves can propagate. On the cut-off surface the wave field is purely

right-hand polarized and E_+ is zero. Thus, there is a region on and outside the cut-off surface in which different physics governs the wave propagation and damping. Little attention has been paid, so far, to this region in the literature. We will also ignore this peripheral region which contains a tenuous plasma.

For small proton concentrations ($C_H \sim 3\%$) the proton-deuteron hybrid resonant surface lies only a small distance (~ 1 cm) away from the cyclotron resonant surface. The cyclotron and hybrid resonant layers overlap, and their mutual influence must be taken into account. The hybrid resonance (an electrostatic resonance) tends to produce a linear E-field polarization (i.e., $E_+ = E_-$) and thus to enhance the magnitude of E_+ -field. The proton cyclotron resonance tends to make the polarization circular and right-handed and thus to reduce the magnitude of E_+ . (The theories by Adam and Samain(12) and by Stix(13) include only the latter effect.) Within the region between the two resonant surfaces both the magnitude and gradient of the E_+ -field are greater than in the absence of the hybrid resonance. The magnetic field inhomogeneity across the plasma cross-section must be included in the analysis, and a treatment in cylindrical geometry becomes difficult. However, because of the relatively thin and straight geometry of the cyclotron and hybrid resonant layers in the important central plasma region, the situation may be amenable to an approximate treatment in a simpler slab

geometry.

Boundary value problems in slab geometry with an inhomogeneous B-field and a non-uniform plasma were solved numerically by Takahashi(1) to obtain the E-field structure. The theory was based upon a second-order differential equation (Eq. 23 of Ref. 1) that treated the PFC (Proton Fundamental Cyclotron) damping in a self-consistent manner, but the DSH (Deuteron Second Harmonic) damping was not included. Thus, the theory was valid only when the latter had small effects. Because of the proximity of the cyclotron resonant layer, the hybrid resonance is heavily damped: the real and imaginary parts of the perpendicular wave number computed for the local plasma properties were comparable and the real part of the coefficient of the second order derivative term in the differential equation did not vanish at the hybrid layer. A criterion for such weak resonances was obtained earlier(1). The resonance is weak when the following inequality is satisfied:

$$C_H \equiv \frac{n_H}{n_e} \leq \frac{2}{1.082} C_D \left(n_H \frac{v_H}{c} \right) \left[\frac{1}{1 - \Omega_D^2} + \frac{n_H^2}{\Omega_{PD}^2} \right] \quad (12)$$

Here, $C_D = n_D/n_e$, $\Omega_D = \omega_{tD}/\omega$, $p_D = \omega_{pD}/\omega$ and c is the speed of light in vacuum. In the present paper we are also concerned with such weak resonance cases.

In Fig. 14 the E-field structure for specific cases are reproduced from Ref. 1. The x-coordinate is in the direction of inhomogeneity with its origin at the slab center. The abscissa of the figure is $R_0 - x$, where $R_0 = 0.9$ m is the scale length of the B-field inhomogeneity ("major radius") at the center. The tokamak axis is to the right of the figure. The slab half-width ("minor radius") is $a = 0.17$ m, and $n_{e0} = 2.5 \times 10^{19} / \text{m}^3$. The top figures are variations of E_x across the slab for three different values of magnetic field: one for which the PFC is located at the center and two others for which the resonance is displaced by 0.06 m on either side of the center. The lower figures are variations of the E_+ -field for the corresponding three cases. (Note the ten-to-one difference in the ordinate between the two groups of figures.) Locations of the ion cyclotron (ICR) and two-ion hybrid (HBR) resonant surfaces are indicated in these figures.

Effects of the weak hybrid resonance appear as small bumps on E_x -curves (E_y is little affected) in Fig. 14. The effects are stronger on the low field side of the axis, evidently as a result of the fact that the field and density gradients are in the same direction there. However, influence of the hybrid resonance on the E_+ -field is significant. Compare $C_H = 0$ and 3 % cases shown in the lower central figure. The large difference is a consequence of the fact that the field polarization is strongly altered by the hybrid resonance. Particularly noteworthy here is the

appearance of a very steep gradient of the E_+ -field over a narrow region between the cyclotron and hybrid resonances. In the case shown in the lower left figure the gradient in this region is enhanced by one order of magnitude over the value obtained in the absence of the hybrid resonance. It is then expected from the physical mechanism discussed earlier that the DSH damping should be enhanced by two orders of magnitude. Since the magnitude of the E_+ -field is also greater than without the hybrid resonance, the PFC damping should also be enhanced. Furthermore, the dependence of both of these damping strengths on the resonant layer location is asymmetric with respect to the plasma center. A more quantitative analysis is presented in the remainder of this section.

Ions that move across the cyclotron resonant layer receive a kick in their perpendicular energies due to cyclotron acceleration. The power absorbed by the ions can be calculated by taking a statistical average of these incremental energy gains. For a slab of unit height and depth the power absorbed by the protons and deuterons are given respectively by (16),

$$P_I = 4\pi R_c \left(\frac{\omega_{PH}^2}{\omega} \right) \frac{\epsilon_0}{4} |E_+(x)|_{x=x_c}^2 J_I \quad (13)$$

$$P_{II} = 4\pi R_c \left(\frac{\omega_{PD}^2}{\omega} \right) \frac{\epsilon_0}{4} |E_+(x)|_{x=x_c}^2 \left(\frac{\rho_D}{2a} \right)^2 J_{II} \quad (14)$$

Here, x_c and R_c are the x-coordinate and the "major radius" of the cyclotron resonant surface. ω_{pH} , ω_{pD} and ρ_D are respectively the proton and deuteron plasma frequencies and the deuteron Larmor radius. The integrals, J_I and J_{II} , are given by,

$$J_I = \frac{2}{\pi^{1/2}} \int_0^\infty d\zeta_H |\mathcal{E}(\zeta_H)|^2 \exp - \zeta_H^2 \quad (15)$$

$$J_{II} = \frac{2}{\pi^{1/2}} \int_0^\infty d\zeta_D |\mathcal{F}(\zeta_D)|^2 \exp - \zeta_D^2 \quad (16)$$

where

$$\zeta_M = (u_{||}/v_M) (x - x_c) / R_c \quad (17)$$

$$\mathcal{E}(x) = E_+(x) / E_+(x_c) \quad (18)$$

$$\mathcal{F}(x) = a [dE_+(x) / dx] / E_+(x_c) \quad (19)$$

Dimensionless quantities, $\mathcal{E}(x)$ and $\mathcal{F}(x)$, are the E_+ -field and its gradient normalized by the E_+ -field at the cyclotron resonant surface. The subscript, M, in Eq. 17 refers to either proton(H) or deuteron(D) and v_M is the most probable speed of the M-species ions. In this analysis the wave number, k_y , (y axis perpendicular to the x-coordinate and B-field) is assumed to be zero.

If the E_+ -field is evaluated at the resonant surface and assumed constant across the resonant layer, then $\xi(x)=1$ and J_I integrates to unity. Thus, the approximation used by Stix would correspond to $J_I=1$ and the actual value of J_I represents the "enhancement factor" of the PFC power absorption. Similarly, if the E_+ -field gradient is approximated by $k_\perp |E_+|$ and evaluated at the resonant surface, J_{II} equals to $(k_\perp a)^2$. Stix further assumed $N_I^2=1$, or equivalently, $k_\perp=k_A$, where k_A is the Alfvén wave number (ω/v_A). $J_{II}/(k_A a)^2$ is therefore the enhancement factor for the DSH power absorption. The results of calculations for some specific cases are presented in the remainder of this section. For these calculations the integrals in Eqs. 15 and 16 are evaluated numerically using the field variation obtained from the solutions of the boundary value problem discussed in an earlier paragraph. The parameters common for all cases are: $a=0.17$ m, $R_0=0.9$ m, $C_H=3$ % and the on-axis ion temperature, $T_{i0}=200$ eV.

Several computed quantities for four different locations of the cyclotron surface are presented in Table III. P_I and P_{II} are, respectively, the power absorbed by the PFC and DSH damping within a slab of unit height and depth. These values are for waves travelling in one direction only, and are normalized in such a way that the Poynting flux is 1 MW per unit height of the slab. Power absorption by both mechanisms is strong when the resonant

surface is located at the low field side of the center, which is consistent with the observed asymmetry in the wave damping (Fig. 9(a)). The asymmetry is particularly pronounced for DSH damping. If the deuterons absorb much more power when the resonant surface is in the low field side, power loss from the plasma through deuteron charge exchange should also be much greater under these conditions. This expectation is not inconsistent with the observed asymmetry of the bolometer signals discussed earlier (Fig. 9(b)). The eigenvalue for the parallel refractive index ($n_{||}$) is about 30 for all cases and thus would correspond to a toroidal mode number of 15. This is considerably greater than the toroidal mode numbers shown in Figs. 11 and 12. (The slab geometry leads to a parallel wavelength considerably shorter than that obtained in a cylindrical geometry with the same peak density. This is partly due to the fact that k_y is taken to be zero.) Thus, the PFC damping is about one order of magnitude stronger than the DSH damping in these cases, and the assumption underlying the theory is valid. When the resonant surface is located at the plasma center, the enhancement factors are respectively 1.8 and 5.2 for the PFC and DSH damping, but if the resonant surface is located 0.04 m away from the plasma center on the low field side, the enhancement factors are respectively 2.8 and 13. (When the resonant surface is further away from the center and comes too close to the fast wave cut-off surface, these enhancement factors become very

large, but the normalizing factor (E_+ -field at the resonant surface) tends to zero. The enhancement factors lose their significance in these cases, since Stix's results were obtained based upon the plasma properties of the central core.)

In Fig. 12(b) variation of the loading resistance is shown for the same case as (a), but with the PFC and DSH damping strengths multiplied respectively by 1.8 and 5.2. Since the enhancement factors should be different for different modes, the results given here are only approximations. However, they indicate that the modest enhancement factors expected from the theory of the two-ion hybrid resonance are sufficient to suppress the prominent toroidal resonance peaks and to make the appearance of the loading curve reasonably similar to experimental observations. The sharp rise of the loading resistance at the onset of the two new modes and subsequent gradual fall-off (instead of gradual rise) of the resistance are the principal discrepancies that are still evident. The similarity should be better when a more realistic plasma model is used and the mode onset is more accurately predicted.

In Table IV the results of computations for four different densities are summarized. At high densities the enhancement factors are close to unity. This shows that the damping strengths at high densities predicted by the present

theory tend toward those of Stix's theory for "pure" PFC and DSH damping. For a given amount of Poynting flux, greater power is dissipated at lower densities by both mechanisms. This tendency is particularly strong for the DSH damping. The power dissipated by this mechanism reaches nearly 15 % at the lowest tabulated density. The enhancement factors are 15 and 590 for the PFC and DSH damping for this case. However, the latter is likely to be an overestimate, because the significant effect of the DSH damping is not included in the E-field calculations and the validity of the theory is questionable.

In Fig. 11(b) the time variation of theoretical loading with enhanced damping is shown. The enhancement factor for the PFC damping is taken to be 15 from Table III, but the factor for the DSH damping is arbitrarily taken to be 100 rather than 590 as indicated in the table. Although one-to-one correspondence of individual peaks is again impossible to obtain, the general appearance and the level of the loading resistance are much closer to the experimental results (Fig. 8). This shows that the experimentally observed damping strengths are not incompatible with enhanced cyclotron damping.

The lack of mode structure in the loading curve was interpreted by some researchers(27) as an indication of anomalous power deposition. Such "parasitic" loading was thought to be especially severe for antennas without

electrostatic shields. However, interpretation of the present experiments is different. At low plasma densities $m=+1$ is the only mode that can carry significant rf energy, whereas at higher densities more propagating poloidal modes become available to carry away rf power. Higher radial modes, which may also become propagating, are especially efficient in this respect. Thus, as the density level increases, the level of the loading resistance also rises as seen in Fig. 10. In the highest density case (#5) the apparent lack of structure is caused by overlapping of strongly damped modes. When the damping is sufficiently strong, the loading resistance should approach the radiation resistance. The observed loading resistance of 0.8-2 ohms is in good agreement with theoretically predicted radiation resistance. This indicates that most of the rf power is indeed radiated as the fast wave, and that it is not necessary in the ATC experiments to consider anomalous phenomena to account for the bulk of the power coupled into the tokamak vessel.

6. TWO-ION HYBRID MODE CONVERSION

The analysis in the previous section is based upon a simple model for the plasma dielectric tensor that results in a second order differential equation. In the limit of the cold plasma approximation the coefficient of the highest order derivative in the equation vanishes at the hybrid resonance, and higher order terms must be retained to

describe the wave propagation in the vicinity of the resonance. The analysis in the previous section deals with physical conditions under which thermal effects, primarily PFC (Proton Fundamental Cyclotron) and DSH (Deuteron Second Harmonic) damping, are strong enough so that the second order term remains more important than the higher order terms. Such physical situations would arise when the proton concentration is small, ion temperature is high, or the parallel wavelength is short. Eq. 12 expresses a quantitative criterion for the presence of the weak hybrid resonance. The ATC plasmas are believed to belong to this category under most plasma heating conditions, and the wave damping data are interpreted in this light.

For conditions under which the higher order terms become more important, relatively undamped short-wavelength waves come into existence, and some of the fast wave power may be channeled into them through linear mode conversion processes. The E-field structure would then be different from the one shown in Fig. 14. (Short-wavelength undulations would appear in the vicinity of the hybrid resonance.) The division of deposited power between the protons and deuterons would also be altered. In fact, the mode converted waves may be damped by electron Landau damping and their power deposited in the electrons. The process should be important when the parallel wavelength is long and the waves propagate nearly perpendicularly to the hybrid surface.

There are theoretical treatments that calculate the power absorption based upon mode conversion. Some (14,15,29) are based upon dispersion relations that are considerably more elaborate than the one used in this paper. However, none of the theoretical treatments solved boundary value problems to obtain the E-field structure in a bounded plasma, and so were only applicable when the wave absorption was so strong that no waves reflected from the boundaries need be considered. The ATC experimental conditions clearly do not correspond to such cases: even at the highest densities for which the wave damping is found strong, the waves propagate around the torus several times (and bounced back and forth many more times from the minor radial boundaries) before their amplitudes e-fold. When the damping due to the mode conversion, crudely derived by Takahashi(30) from Budden's theory, is used in calculating the loading resistance response at low densities (such as one in Fig. 11), it is found that the toroidal resonance peaks are completely annihilated, which disagrees with experimental observations (Fig. 8). It is thus believed that in the ATC experiments the primary effects of the IHB (two-Ion HyBrid) resonance are enhancement of both the PFC and DSH damping rather than channeling of the fast wave power into electrostatic waves through mode conversion. Because of the lack of more elaborate theories applicable to the ATC conditions, however, detailed comparisons of the experimental results and theoretical predictions cannot be

made.

In a plasma with changing density each propagating mode with given poloidal and radial mode numbers goes through different phases of wave propagation characteristics. Near the onset the parallel wavelength is very long and the waves propagate nearly perpendicularly to the the confining B-field. At higher densities the same mode has a shorter wavelength. In each of these phases a different physical mechanism is responsible for the wave damping. Swanson(14) showed that waves propagating perpendicularly to the IHB resonant surface are strongly damped by mode conversion. It is conceivable, also, in the ATC experiments that modes near the onset are indeed mainly damped by mode conversion and some rf power is deposited in the electrons. However, such phases may be passed through relatively quickly, preventing the characteristics of mode conversion damping to be observed clearly. It requires more detailed experimental studies under carefully controlled density variation to ascertain the importance of the mode conversion damping. It is shown here that the gross features of the loading resistance behavior can be accounted for by a simpler interpretation in terms of enhanced cyclotron damping.

7. CONCLUSIONS

In the ATC ICRF experiments a wave coupling structure is employed which is considerably simpler than those used in the previous experiments: it consists of essentially a semi-circular metal strip and a pair of bumper limiters. In particular these experiments demonstrate that satisfactory antenna performance can be obtained without an electrostatic shield commonly used in the previous experiments. The success of the ATC antenna represents a step forward towards a simpler and more rugged (presumably all-metal type) structure which could be used in a thermonuclear environment. (Although the antenna and limiter faces are covered by ceramic components, they may be redundant. However, this cannot be ascertained from the present experiments.)

It is shown that the major part of the power coupled into the tokamak vessel is radiated as fast waves and that there is no evidence that significant power is channeled through so-called "parasitic loading". This is not to say that no rf power is channeled into plasma phenomena other than the fast wave generation. The conclusion is based upon comparisons of the experimental results with predictions using a theoretical model which has limited capability in reproducing some aspects of the experimental observations. If there are any "parasitic" effects, however, they are not of major consequence in carrying away the rf power.

A coupling efficiency (net coupled power/incident power) in excess of 90 % is routinely attained. The total loading resistance is 1-2 ohms under typical plasma heating conditions compared to the loading resistance in the absence of plasma of about 0.1 Ohm. Thus, the circuit loss accounts for 5-10 % and a wave generating efficiency (wave power/net coupled power) of 90-95 % is achieved. Wave damping measurements with the cyclotron resonance layer inside and outside the plasma show that about 10 % of the wave power is dissipated in the vessel walls and that the power deposition efficiency (power deposited in plasma/wave power) is about 90 %. The overall efficiency of delivering the power to the plasma is the product of all these efficiencies: approximately 75 % of the generator output power is actually deposited in the plasma.

Direct measurements of the wave damping strength is possible only at low plasma densities at which individual modes are well separated. The measured damping strength under such conditions is found to be much stronger than that predicted by the "pure" PFC (Proton Fundamental Cyclotron) and DSH (Deuteron Second Harmonic) damping theories. The range of major radius locations of the resonant surface over which damping is strong is found to be asymmetric with respect to the geometrical plasma center and shifted toward the low field side. This is in contradiction to the predictions of the "pure" cyclotron damping theories. At higher densities the presence of many strongly damped modes

makes quantitative determination of the damping strength impossible. Under these conditions the experiments are checked against the theories by comparing the general features of the time variation of the loading resistance. The "pure" cyclotron damping theories appear to predict damping which is also too small at high densities, but the discrepancies are smaller than at low densities.

Both of these discrepancies are attributed to the IHB (two-Ion HyBrid) resonance between the minority protons and majority deuterons. The hybrid resonance strongly alters the E-field polarization within the cyclotron resonant layer in such a way as to enhance both the PFC and DSH damping. Enhancement of wave damping strength is calculated in a slab geometry using a second-order differential equation. The enhancement factors tend to unity at high densities, which shows that the "pure" cyclotron damping theories of Adam and Samain(12) and of Stix(13) would give essentially correct damping strength at high densities. The latter theory was developed specifically for reactor regimes. However, at low densities the enhancement factors may become very large. Large discrepancies reported in earlier literature between the experimental damping strength and the "pure" cyclotron damping theories thus stem from inapplicability of these theories at low densities at which damping strength measurements are possible. The enhancement factors are also found asymmetric with respect to the plasma center in agreement with the experiments. This asymmetry is

particularly large for the DSH damping.

The ATC experimental results are interpreted in terms of the enhanced cyclotron damping rather than the IHB mode conversion damping, which would deposit rf power into the electrons. A detailed examination of the experimental data in the light of mode conversion is not made, primarily because of the lack of adequate theoretical calculations applicable to ATC plasma conditions. In order to ascertain the importance of the mode conversion process, both theory and experiments must be refined considerably.

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TABLE CAPTIONS

Table I

The theoretical wave damping length is shown normalized to the toroidal circumference, for the toroidal eigenmodes shown in Fig. 11(a). The normalized damping length is the number of times the waves go around the torus before their amplitudes e-fold. A set of radial(l), poloidal(m) and toroidal(n) eigenmode numbers identifies the resonant mode. The damping lengths are computed based upon "pure" cyclotron damping theories for the proton fundamental cyclotron(I) and deuteron second harmonic(II) damping. The electron Landau/TMP damping (E) is completely negligible. The wall resistive damping strengths are assumed to be 20 times theoretical values for a cylindrical waveguide made of smooth stainless steel walls.

Table II

The theoretical wave damping length is shown normalized to the toroidal circumference, for the toroidal eigenmodes shown in Fig. 12(a). See caption to Table-I for symbol explanations. The electron Landau/TMP damping (E) is completely negligible except where entries are made.

Table-III

These are some computed quantities showing the influence of the two-ion hybrid resonance on the proton fundamental cyclotron (subscript I) and deuteron second harmonic

(subscript II) damping. The damping length is tabulated for four different values of the on-axis plasma density (n_{e0}), together with corresponding values of the parallel refractive index ($n_{||}$).

Table-IV

These are some computed quantities showing the influence of the two-ion hybrid resonance on the proton fundamental cyclotron (subscript I) and deuteron second harmonic (subscript II) damping. The damping length is tabulated for four different locations of the cyclotron resonant layer (R_c), together with corresponding values of the parallel refractive index ($n_{||}$).

TABLE I

(l, m, n)	I	II	E	W	T
(1, 1, 3)	170	230	-	13	12
(1, 1, 4)	56	100	-	11	9
(1, 2, 3)	240	410	-	17	16
(1, 2, 4)	79	150	-	13	11
(1, 3, 4)	95	220	-	17	14

TABLE II

(1, m, n)	I	II	E	W	T
(1, -1, 1)	76	1.7	-	55	1.6
(1, -1, 2)	29	2.4	-	27	2.0
(1, 0, 6)	5.0	5.4	160	9.2	2.0
(1, 1, 10)	3.1	24	80	17	2.3
(2, 1, 1)	81	1.7	-	54	1.7
(2, 1, 2)	33	2.7	-	23	2.3
(1, 2, 8)	6.5	20	140	8.6	3.1
(1, 3, 8)	7.0	20	160	9.6	3.3

TABLE III

R_c	m	0.98	0.94	0.90	0.86
n_{II}		28	29	30	31
P_I	kw/m ²	25	78	55	6.0
P_{II}	kw/m ²	1.6	4.0	2.0	0.11
J_I		7.2	2.8	1.8	1.2
$J_{II}/(k_A a)^2$		64	13	5.2	2.1
Q_I		560	180	230	2100
Q_{II}		8500	3400	6500	120000
(L_D/L)		14	4.5	6.4	58
(L_D/L)		220	89	180	3100

TABLE IV

n_{eo}	$10^{19}/m^3$	0.8	1.5	2.5	3.5
n_{II}		14	22	30	36
P_I	kw/m^2	300	95	55	40
P_{II}	kw/m^2	44	12	2.0	0.55
J_I		15	3.1	1.8	1.3
$J_{II}/(k_A a)^2$		590	54	5.2	1.1
Q_I		570	150	230	360
Q_{II}		3900	1200	6500	26000
(L_D/L)		1.2	3.7	6.4	8.8
(L_D/L)		7.1	30	180	640

FIGURE CAPTIONS

Fig. 1

Wave coupling structure shown with a cut-away view of the ATC vacuum vessel and plasma.

Fig. 2

Antenna and limiters installed on a mock-up structure. The indentations of the limiters are introduced to accommodate other experiments. They and components mounted on them are extraneous to the rf coupling system.

Fig. 3

Locations of various instrumentation around the torus.

Fig. 4

Schematic of the wave coupling structure (shown in part).

Fig. 5

Oscillograms of microwave interferometer, plasma current and one-turn loop voltage signals with and without rf heating. The rf pulse is from 20 to 30 ms.

Fig. 6

Fast wave dispersion relation in terms of Alfvén refractive indices. Solid curve is the exact solution and broken curve the "elliptic" approximation. ●'s indicate discrete waveguide modes for a set of plasma conditions in ATC.

Fig. 7

Time variation of the wave amplitude (rf envelope) observed

by two magnetic probes located nearly directly above (T) and below (B) the plasma axis.

Fig. 8

Oscillograms of the loading resistance computer outputs. The toroidal magnetic field is different for each oscillogram. The locations of the deuteron second harmonic cyclotron resonant surface are indicated.

Fig. 9

(a) Variation of the wave damping length normalized to the toroidal circumference as the toroidal field strength is varied. The abscissa is the major radius location of the deuteron second harmonic resonant layer. The region of minimum damping length is asymmetric with respect to the nominal plasma center at 87 cm. (b) Variation of bolometer signal which registers power loss through charge exchange and radiation. The net difference between discharges with and without rf pulse is plotted. The nominal extent of the plasma region is indicated. The outer bumper limiter is positioned at 107 cm.

Fig. 10

Oscillograms of the R_g -computer output. The cyclotron resonant surface is located near the plasma center in all cases. Variation of the central electron density for each case is also shown.

Fig. 11

Theoretical loading resistance corresponding to the experimental conditions of the center oscillogram of Fig. 8.

(a) Damping strengths based upon minority and second harmonic damping obtained by Stix (b) Minority damping strength is multiplied by an enhancement factor of 15 and the second harmonic damping strength by 100.

Fig. 12

Theoretical loading resistance for the experimental conditions of the top-right oscillogram of Fig. 10.

(a) Damping strengths based upon minority and second harmonic damping obtained by Stix (b) Minority damping strength is multiplied by an enhancement factor of 1.8 and the second harmonic damping strength by 5.2.

Fig. 13

Geometries of the cyclotron ($\Omega_H=1$) and two-ion hybrid ($H=0$) resonant surfaces and of the fast wave cut-off ($F=0$) surface. The proton concentration ($C_H=10\%$) is taken to be much higher than the expected actual value for clearer graphical presentation of the surfaces. The tokamak center is to the right of the figure.

Fig. 14

(a) Variation of the E_x -component of the wave electric field across the slab for three different values of B_z : one for which the proton cyclotron resonance falls at the plasma center (0.9 m) and two others for which the resonance is

displaced by 0.06 m on either side of the center. The slab half-width is 0.17 m. The hybrid resonance has only small effects on E_x , but they are asymmetric with respect to the center.

(b) Variation of E_+ -field for the same three cases. The hybrid resonance has strong effects on E_+ . Note the appearance of a very steep gradient between the hybrid and cyclotron resonant layers.

Fig. 1. 766223

Fig. 3. 786130

Fig. 4. 793279

RF OFF

RF ON

V_L
 I_V
0

I_p
-60 kA

V_L
 I_V
0

I_p
-60 kA

0

FIG. 5. 780-06

Fig. 6. 793207

Fig. 7. 773469

Fig. 8. 772590

Fig. 9. 793278

Fig. 10. 766046

Fig. 11. 793211

Fig. 12. 793210

Fig. 13. 793277

Fig. 14. 773179