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A RESONANT-CAVITY ICRF COUPLER FOR LARGE TOKAMAKS

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ABSTRACT

A new resonant-cavity ICRF coupler is proposed for large tokamaks. The design features a novel resonant cavity, an rf magnetic field orientation that effectively radiates fast Alfvén waves, matching to 40 Ω transmission lines, and an electric field orientation so that the strongest rf electric fields are orthogonal to the main toroidal magnetic field thereby benefitting from magnetic insulation. As a result, the power handling capability is excellent. For the case of the "Rig-Dee" Doublet III tokamak, a single 35 cm $\times$ 50 cm coupler can launch 20 MW of fast Alfvén waves. Extrapolation to fusion reactor parameters is straightforward.

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I. INTRODUCTION

Plasma heating in tokamaks by means of fast Alfvén waves in the Ion-Cyclotron Range of Frequencies (ICRF) has achieved notable experimental successes¹⁻³ which are understood in terms of theory.⁴⁻⁸ As a result, ICRF plasma heating has become the preferred option for heating first-generation tokamak reactors to ignition.^{9,10}

In all the experiments to date, the couplers (i.e., the antennas) which radiate the fast Alfvén waves into the plasma have been induction loops located within the vacuum vessel. In addition, the coupling loops in the PLT tokamak are covered with a ceramic insulator.¹¹ Such antennas are unsatisfactory for use in a fusion reactor, where engineering considerations require a modular, easily replaced antenna. The potential radiation damage to insulators from fusion neutrons calls for an all-metal coupler design. Furthermore, a high power flux ($\sim 10 \text{ kW/cm}^2$) is desirable in a reactor to reduce the fraction of the wall area devoted to heating.

But a reactor coupler design must be proved on research devices such as the "Big-Dee" Doublet III,¹² JET,¹⁰ and TFTR¹³ tokamaks. These tokamaks have magnetic field strengths that are the same as (or smaller than) the fields envisioned for a reactor. Since most reactor ICRF heating schemes utilize second harmonic heating, the impressed frequencies are similar in reactor and research tokamaks, while the actual size of a research tokamak is roughly a factor-of-two smaller than that of a fusion reactor. The design of resonant cavities must be able to accommodate the smaller research tokamaks with no change in their basic configurations.

This paper reports the electrical properties of a new resonant-cavity ICRF coupler based on the principles of small-gap cavities used in microwave tubes. The design takes advantage of the large port size that will be

available on the next generation of research tokamaks. The coupler introduces power through a large area Faraday screen coincident with the vacuum vessel wall. Impressive power capabilities result.

Our calculation is based on dimensions appropriate to the "Big-Dee" Doublet III tokamak and can be readily repeated for TFTR or reactor installations. The motivation for the design is fourfold: (1) It is an all-metal, modular design, utilizing a large area Faraday screen, and does not protrude beyond the vacuum vessel wall. (2) The rf magnetic field orientation is similar to coil couplers and effectively radiates fast Alfvén waves. (3) The strong electric fields associated with the reactive energy are contained entirely within the cavity, and oriented perpendicular to the main toroidal field, thereby benefitting from magnetic insulation. (4) The vacuum interfaces are located in a pair of 40 Ω rectangular coaxial transmission lines which are matched to the cavity and carry only forward power. Features (3) and (4) provide the excellent power handling capability. A single resonant cavity coupler should be able to introduce 20 MW of ICRF power into Doublet III, with electrical fields significantly below measured breakdown levels. This corresponds to a power density of 10 kW/cm² - a density entirely adequate for reactor applications. With its all-metal construction, the Doublet-III cavity coupler will serve as a demonstration of a reactor ICRF coupler.

In Sec. II the propagation of the fast Alfvén wave in the periphery of the plasma is discussed and calculated. This calculation provides the basis for computing the radiation resistance of the coupler. Section III provides a description of the coupler configuration and computations of its electrical properties. The paper concludes with a discussion.

II. FAST ALFVEN WAVE PROPAGATION

The surface region of the plasma provides the transition from the low-density "scrape-off" region between the vacuum vessel and the limiter to the high density interior plasma where the wavelength of the fast Alfvén wave is short and geometrical optics considerations apply. This section addresses the questions: What is the appropriate definition of the plasma edge in the context of fast Alfvén wave propagation? What is the effective surface impedance of the plasma? We shall find that a procedure analogous to the way the skin effect is treated in metallic resonant cavities provides an accurate method for computing the power carried by fast Alfvén waves into the plasma interior. In particular, a surface (to be called plasma boundary) is defined at which the tangential electric field vanishes in the first approximation, but where the tangential magnetic field has a finite value. The (small) surface impedance of the plasma then gives the tangential electric field and consequently the fast Alfvén wave power flux into the plasma in terms of the tangential magnetic field at the boundary.

The starting point is a slab model of the plasma periphery with a straight, uniform magnetic field in the z -direction. The parallel electric field is considered to be negligible everywhere because of a Faraday screen located in the plane of the vacuum vessel wall. The consequences of an imperfect Faraday screen¹⁴ are beyond the scope of this paper. Figure 1 shows the overall configuration and the coordinate system used in the calculations. Both y and z are ignorable coordinates, which allows the use of Fourier transforms in these directions. As a result, the equations governing propagation of the fast Alfvén wave can be written simply as

$$\left(\frac{\omega^2 \mu_0^2}{c^2 (\omega^2 - \Omega^2)} + k^2 + k_{\parallel}^2 \right) E_x + i \left[\frac{\omega^2 \mu_0^3}{c^2 \Omega (\omega^2 - \Omega^2)} E_y + k \frac{\partial}{\partial x} E_y \right] = 0 \quad , \quad (2.1)$$

$$-i \left(\frac{\omega^2 \omega^3}{c^2 \Omega (\omega^2 - \Omega^2)} E_x - k \frac{\partial}{\partial x} E_x \right) + \left[\frac{\omega^2 \omega^2}{c^2 (\omega^2 - \Omega^2)} + k_{\parallel}^2 - \frac{\partial^2}{\partial x^2} \right] E_y = 0 \quad (2.2)$$

where the ion plasma frequency is related to the electron density via

$$\omega_{pi}^2 = \frac{4\pi e^2 Z}{AM_p} n(x) \quad (2.3)$$

and is regarded as a function of x . The Fourier transforms are defined by

$$\vec{E}(x, k, k_{\parallel}) e^{-i\omega t} \equiv \int dz dy e^{-i(ky + k_{\parallel}z + \omega t)} \vec{E}_0(x, y, z) \quad (2.4)$$

and the angular frequency will be taken to be $\omega = 2\Omega$ as is appropriate for the second harmonic heating favored by reactor applications studies.^{9,10}

In the periphery of the main plasma (but still inside the limiter) the electron density profile can be well-modeled by a linear function

$$n(x) = \frac{2n_0 x}{a} \quad (2.5)$$

For the particular case of a parabolic density profile, n_0 and a represent the central density and plasma minor radius, respectively. If we further set $k = k_{\parallel} = 0$, then Eqs. (2.1) and (2.2) combine into the simple form

$$\frac{\partial^2}{\partial x^2} E_y + \frac{2\omega^2 \omega^2}{ac^2 \Omega^2} x E_y = 0 \quad (2.5)$$

where $\omega_{pi0}^2 = (4\pi e^2 Z n_0 / AM_p)$. By introducing the nondimensional variable

$$v = x/x_0 \quad , \quad (2.6)$$

$$x_0 = \left(\frac{ac^2 \omega^2}{8\omega_{p0}^2} \right)^{1/3} = \left(\frac{-ac^2}{8\omega_{p0}^2} \right)^{1/3} \quad , \quad (2.7)$$

one can cast Eq. (2.5) into the form of an Airy Equation:

$$\frac{\partial^2}{\partial v^2} E_y + v E_y = 0 \quad . \quad (2.8)$$

The appropriate solution to Eq. (2.8) corresponds to waves propagating away from the antenna and is

$$E_y = \text{Bi}(-v) + i \text{Ai}(-v) \quad (2.9)$$

in terms of the functions defined in Abramowitz and Stegun.¹⁵

Equations (2.5) - (2.9) play an important role in the physics of fast wave propagation in the plasma periphery. Equation (2.8) shows that the transition to the geometrical optics regime occurs in the region $v \gg 1$ and that a characteristic scale length x_0 and density $\bar{n}$ are involved:

$$x_0 = \left(\frac{ac^2}{8\omega_{p0}^2} \right)^{1/3} = 3.5 \text{ cm} \left(\frac{a}{70\text{cm}} \right)^{1/3} \left(\frac{10^{14} \text{ cm}^{-3}}{n_0} \right)^{1/3} \left(\frac{A}{Z} \right)^{1/3} \quad , \quad (2.10)$$

$$\bar{n} = \frac{2n_0 x_0}{a} = 1.0 \cdot 10^{13} \text{ cm}^{-3} \left(\frac{n_0}{10^{14} \text{ cm}^{-3}} \right)^{2/3} \left(\frac{70\text{cm}}{a} \right)^{2/3} \left(\frac{A}{Z} \right)^{1/3} \quad . \quad (2.11)$$

The values for these quantities make two physics points clear: First, transition to the geometrical optics propagation occurs at a density well below the central density but still substantially above the plasma densities

expected in the scrape-off region. The linear density profile model should be good in this regime. Secondly, the characteristic radial scale length x_0 is quite short, so that the parallel and poloidal wavelengths, which greatly exceed x_0 for large tokamaks, should not play a major role in the transition to geometrical optics propagations. This point will be made quantitative below.

Let us now introduce a model density profile which will both be a reasonable profile and provide a mathematically tractable expression for analysis. This profile is

$$n(x) \approx \begin{cases} \frac{2n_0}{a} x = \bar{n} v & x > x_0 \\ \bar{n} \exp\left(-\frac{x-x_0}{x_0}\right) \approx \bar{n} e^u & x < x_0 \end{cases} \quad (2.12)$$

where

$$u = v - 1 = \frac{x-x_0}{x_0} \quad (2.13)$$

and $x = x_0$ is the point where $n(x) = \bar{n}$. In the low density scrape-off regime, this model produces an experimental fall-off of density with a characteristic length $x_0 \sim 4$ cm.

We can next use the definitions

$$\beta = kx_0 \quad q = k_{\parallel}x_0 \quad (2.14)$$

and the quantities $\bar{n}$, x_0 to cast Eqs. (2.1) and (2.2) into nondimensional form. The key feature of this representation is that for large tokamaks the quantities β and q will be small for all values k , $k_\parallel$ where the power spectrum of the coupler has appreciable values. This will permit useful solutions in closed form.

First, consider the high density region where $v > 1$. The nondimensional equations are

$$E_x = -i \frac{vyE + \beta(y^2 - 1)(\partial E / \partial v)}{v + (y^2 - 1)(q^2 + \beta^2)} \quad (2.15)$$

$$-i \left(\frac{vy}{y^2 - 1} E_x - \beta \frac{\partial}{\partial v} E_x \right) + \left(\frac{v}{y^2 - 1} + q^2 - \frac{\partial^2}{\partial v^2} \right) E = 0 \quad (2.16)$$

where

$$y = \frac{\omega}{Q} \quad \text{and} \quad E = E_y \quad (2.17)$$

Equations (2.15) and (2.16) can be combined to give

$$\begin{aligned} \frac{\partial^2 E}{\partial v^2} + \frac{(y^2 - 1)\beta^2}{[v + (y^2 - 1)(q^2 + \beta^2)][v + (y^2 - 1)q^2]} \frac{\partial E}{\partial v} \\ + \left(\frac{v^2 - 2vq^2 - q^4(y^2 - 1)}{v + q^2(y^2 - 1)} - \frac{\beta y(\beta^2 + q^2)(y^2 - 1)}{[v + (y^2 - 1)(q^2 + \beta^2)][v + (y^2 - 1)q^2]} - \beta^2 \right) E = 0 \end{aligned} \quad (2.18)$$

It is clear that for $v \gg 1$, terms involving q and β can be ignored when $q, \beta \ll 1$, and Eq. (2.8) is an excellent approximation to Eq. (2.18).

In the low density regime ($v < 1$, $u < 0$) similar arguments yield the

equations

$$E_x = -i \frac{y e^u E + \beta (y^2 - 1) (\partial E / \partial u)}{e^u + (y^2 - 1)(q^2 + \beta^2)} \quad (2.19)$$

and

$$-i \left(\frac{e^u y}{y^2 - 1} E_x - \beta \frac{\partial}{\partial u} E_x \right) + \left(\frac{e^u}{y^2 - 1} + q^2 - \frac{\partial^2}{\partial u^2} \right) E = 0 \quad (2.20)$$

These equations can be combined in two useful ways. First, a straightforward substitution yields

$$\begin{aligned} \frac{\partial^2 E}{\partial u^2} + \frac{\beta^2 (y^2 - 1) e^u}{[e^u + (q^2 + \beta^2)(y^2 - 1)][e^u + q^2 (y^2 - 1)]} \frac{\partial E}{\partial u} \\ + \left(\frac{e^{2u} - 2q^2 e^u - q^4 (y^2 - 1)}{e^u + q^2 (y^2 - 1)} - \frac{\beta y (q^2 + \beta^2) (y^2 - 1) e^u}{[e^u + (q^2 + \beta^2)(y^2 - 1)][e^u + q^2 (y^2 - 1)]} - \beta^2 \right) E = 0. \end{aligned} \quad (2.21)$$

This equation shows when $e^u \gg q^2, \beta^2$, Eq. (2.21) can be simplified to

$$\frac{\partial^2 E}{\partial u^2} + e^u E = 0 \quad (2.22)$$

which can be solved in terms of Bessel functions¹⁵

$$E = A J_0(\xi) + B Y_0(\xi) \quad (2.23)$$

where

$$\xi = \omega e^{u/2} \quad . \quad (2.24)$$

The coefficients A and B are determined by the requirement of continuity of E and $\partial E / \partial u$ with the solution to Eq. (2.9) at $u = 0$ ($v = 1$). This yields the equations

$$A J_0(2) + B Y_0(2) = B i(-1) + i A i(-1) \quad (2.25)$$

$$-A J_1(2) - B Y_1(2) = -B i'(-1) - i A i'(-1)$$

whose solution is

$$\begin{aligned} A &= 0.9848 + i(0.1638) & B &= -0.2283 + i(0.9775) \\ A/B &= -0.0642 - i(0.9921) \end{aligned} \quad (2.26)$$

Equation (2.21) can also be easily solved in the limit $e^u \ll \alpha^2, \beta^2$:

$$E = D_1 e^{-\kappa u} + D_2 e^{\kappa u} \quad (2.27)$$

where

$$\kappa = (\alpha^2 + \beta^2)^{1/2} \quad . \quad (2.28)$$

It remains to find the solutions to Eqs. (2.19) - (2.20) in the region $e^u \sim \alpha^2, \beta^2 \ll 1$, and to match the solutions of Eqs. (2.23) and (2.27). This is accomplished by introducing ϕ via the definition

$$[e^u + (q^2 + \beta^2)(y^2 - 1)]\phi = [e^u + q^2(y^2 - 1)] \frac{\partial E}{\partial u} - \beta y e^u E \quad (2.29)$$

The second-order differential equation system is then represented by two coupled first-order equations

$$\frac{\partial \phi}{\partial u} + \frac{\beta y e^u}{e^u + q^2(y^2 - 1)} \phi + \frac{(e^{2u} - 2q^2 e^u - q^4(y^2 - 1))}{e^u + q^2(y^2 - 1)} E = 0 \quad (2.30)$$

$$\frac{\partial E}{\partial u} - \frac{\beta y e^u}{e^u + q^2(y^2 - 1)} E - \frac{e^u + (q^2 + \beta^2)(y^2 - 1)}{e^u + q^2(y^2 - 1)} \phi = 0 \quad (2.31)$$

The further substitutions

$$\phi = \bar{\phi} [e^u + q^2(y^2 - 1)]^{-\beta y} \quad E = \bar{E} [e^u + q^2(y^2 - 1)]^{\beta y} \quad (2.32)$$

generate the simple equations

$$\frac{\partial \bar{\phi}}{\partial u} = - \frac{[e^{2u} + 2q^2 e^u - q^4(y^2 - 1)]}{[e^u + q^2(y^2 - 1)]^{1+2\beta y}} \bar{E} \quad (2.33)$$

$$\frac{\partial \bar{E}}{\partial u} = \frac{e^u + (q^2 + \beta^2)(y^2 - 1)}{[e^u + q^2(y^2 - 1)]^{1+2\beta y}} \bar{\phi} \quad (2.34)$$

The approximate solution is obtained by ignoring the right-hand-side of Eq. (2.33) so that

$$\bar{\phi} = \phi_0 \quad (2.35)$$

and

$$E = -\phi_0 [e^u + q^2(y^2 - 1)]^{\beta y} \int_0^u \frac{e^{u'} + (q^2 + \beta^2)(y^2 - 1)}{[e^{u'} + q^2(y^2 - 1)]^{1+2\beta y}} du' + C \quad (2.36)$$

with ϕ_0 and C being constants of integration which are determined matching Eq. (2.36) to the expansion of Eqs. (2.23) and (2.24) for small values of ξ in the region $1 \gg e^u \gg q^2, \beta^2$. The approximate expansions are

$$E = A + \frac{B}{\pi}(u + 2\gamma) = \frac{\phi_0}{\beta y} \sinh(\beta y u) + C = \phi_0 u + C \quad (2.37)$$

and the matching yields

$$\phi_0 = \frac{B}{\pi} \quad C = A + \frac{2\gamma B}{\pi} \quad (2.38)$$

The integration of Eq. (2.36) for values of u such that $e^u \ll q^2, \beta^2$ can be effected by the substitution

$$e^{u'} = q^2(y^2 - 1)\eta \quad (2.39)$$

which transforms Eq. (2.36) into

$$E = \frac{-\phi_0}{[q^2(y^2 - 1)]^{\beta y}} \left(\frac{\beta^2 + q^2}{q^2} \int_{\eta_0}^{\eta_{\max}} \frac{d\eta}{\eta(1+\eta)^{1+2\beta y}} + \int_{\eta_0}^{\eta_{\max}} \frac{d\eta}{(1+\eta)^{1+2\beta y}} \right) + C \quad (2.40)$$

where

$$\eta_0 = e^u [q^2(y^2 - 1)]^{-1} \ll 1$$

$$\eta_{\max} = [q^2(y^2 - 1)]^{-1} \gg 1 \quad . \quad (2.41)$$

Let us concentrate on the integral

$$I = \int_{\eta_0}^{\eta_{\max}} \frac{d\eta}{\eta(1+\eta)^{1+2\beta y}} = \int_{\eta_0}^{\eta_{\max}} \frac{d\eta}{\eta(1+\eta)} + \int_0^{\infty} \frac{d\eta}{\eta(1+\eta)} \left(\frac{1}{(1+\eta)^{2\beta y}} - 1 \right) \quad (2.42)$$

where in the second integral the limit of integrations can be extended because the integral converges. The substitution $1 + \eta = e^s$ transforms the second integral

$$\int_0^{\infty} \frac{d\eta}{\eta(1+\eta)} \left(\frac{1}{(1+\eta)^{2\beta y}} - 1 \right) = -2 \int_0^{\infty} \frac{ds e^{-s\beta y}}{(e^s - 1)} \sinh(\beta y s) = -\frac{\pi^2}{3} \beta y \quad (2.43)$$

where we have assumed $\beta y \ll 1$. Therefore we obtain

$$I = \ln \left(\frac{\eta_{\max} (1+\eta_0)}{\eta_0 (1+\eta_{\max})} \right) - \frac{\pi^2}{3} \beta y \quad . \quad (2.44)$$

We can readily perform the remaining integrals and obtain the solution

$$E = \phi_0 \left(-\frac{\sinh(u_0 \beta y)}{\beta y} + \left(\frac{\beta^2 + q^2}{q^2} \right) e^{-u_0 \beta y} \left(u - u_0 + \beta y \frac{\pi^2}{3} \right) \right) + C \quad (2.45)$$

where

$$u_0 = \ln [q^2 (y^2 - 1)] \quad . \quad (2.46)$$

Let us again make use of the inequality $\beta y \ll 1$ so that the solution for E can

be simplified to

$$E = \phi_0 e^{-u_0 \beta y} \left[u + \frac{\beta^2}{q^2} (u - u_0) + \beta y \left(\frac{\pi^2 (\beta^2 + q^2)}{3 q^2} + u_c^2 \right) \right] \quad (2.47)$$

The key quantity in determining the plasma response is

$$\frac{1}{E} \frac{\partial E}{\partial u} = \frac{1}{\left\{ u - u_0 \left(\frac{\beta^2}{\beta^2 + q^2} \right) + \beta y \left[\frac{\pi^2}{3} + \frac{u_0^2 q^2}{(\beta^2 + q^2)} \right] + \frac{u_0 \beta y}{(q^2 + \beta^2)} \left(2\gamma + \frac{\pi A}{B} \right) \right\}} \quad (2.48)$$

Let us rewrite the vacuum solution Eq. (2.27) as

$$E = (1 + i\kappa \Delta_1) \exp[-\kappa(u - u_1)] - (1 - i\kappa \Delta_1) \exp[\kappa(u - u_1)] \quad (2.49)$$

where $\kappa = (q^2 + \beta^2)^{1/2} \ll 1$. Equation (2.49) can be expanded to give

$$\frac{1}{E} \frac{\partial E}{\partial u} = \frac{1}{u - u_1 - i\Delta_1} \quad (2.50)$$

Identification of expansions Eqs. (2.48) and (2.50), yields

$$\Delta_1 = - \frac{q^2}{(q^2 + \beta^2)} \operatorname{Im} \left(\frac{\pi A}{B} \right) e^{u_0 \beta y} = 3.117 \frac{q^2}{q^2 + \beta^2} e^{u_0 \beta y} \quad (2.51)$$

$$u_1 = - \left(\frac{q^2}{q^2 + \beta^2} \right) e^{u_0 \beta y} \left[2\gamma + \pi \operatorname{Re} \left(\frac{A}{B} \right) \right] + u_0 \frac{\beta^2}{q^2 + \beta^2} - \beta y \left(\frac{\pi^2}{3} + \frac{u_0^2 q^2}{q^2 + \beta^2} \right) \quad (2.52)$$

Equations (2.51) and (2.52) are the key results. From Eq. (2.49), it is clear that u_1 defines the plasma boundary — the point where the tangential electric almost vanishes. Although this definition depends on β^2 and q^2 , it is not

particularly sensitive to their ratio since the value $u_0 \approx -1.5$ characterizes many proposed tokamak installations. Therefore, to a good approximation, we can take

$$u_1 \approx -1.2 + 0(\beta y) \quad (2.53)$$

and the plasma boundary is effectively the place where the plasma density is about $n = 3 \cdot 10^{12} \text{ cm}^{-3}$ (for a proton plasma). At that position the tangential magnetic H_z is given by

$$i\omega\mu_0 H_z = \frac{1}{x_0} \frac{\partial E}{\partial u} \frac{q^2}{q^2 + \beta^2} = \frac{1}{x_0} \frac{E}{(u - u_1 - i\beta_1)} \frac{q^2}{(q^2 + \beta^2)} \quad (2.54)$$

and the Poynting Flux is

$$P_x = \text{Re}(E H_z^*) = 3.117 (\omega\mu_0 x_0) |H_z|^2 e^{u_0 \beta y} \quad (2.55)$$

The real part of the plasma surface impedance Z_p follows directly from Eq. (2.55) when $\beta y \ll 1$ and is

$$Z_p = 3.117 (\omega\mu_0 x_0) = \frac{v}{z} \eta_0 \quad (2.56)$$

where $\eta_0 = 377\Omega$, $v = \omega/2\pi$ is the impressed frequency, and

$$\frac{v}{z} = \frac{c}{(3.117)2\pi x_0} = 428 \text{ MHz} \left(\frac{z}{\text{\AA}}\right)^{1/3} \left(\frac{n_0}{10^{14} \text{ cm}^{-3}}\right)^{1/3} \left(\frac{70 \text{ cm}}{a}\right)^{1/3} \quad (2.57)$$

Discussion: Equations (2.52), (2.53), and (2.56) are the main results of this section. The surface impedance Z_p permits us to calculate the energy

flux into the plasma, given the tangential magnetic field at the plasma surface. Let us note that ν_0 depends on x_0 , which is the scale length for exponential decay of the density in the region where the plasma density is close to the value $\bar{n}$ [c.f. Eq. (2.11)]. We should emphasize that the results for the real part of the surface impedance are independent of the nondimensional wave numbers q and β when $\beta y \ll 1$. Even when βy is not particularly small, the power spectrum of the coupler is an even function β and the average value of the exponential factor in Eq. (2.55) remains close to unity.

A skin-effect method for computing the power radiated into the plasma requires a determination of the "plasma surface" where (to a first approximation) the tangential electric field vanishes. The power radiated into the plasma will depend on the location of this surface, because the closer it is to the coupler, the higher the value of H_z at the surface, and consequently, the higher the power radiated into the plasma [c.f. Eq. (2.55)]. Equation (2.52) defines the location of the plasma surface. This definition depends on the nondimensional wave numbers and creates an uncertainty of the order of x_0 in the location of the plasma surface. This uncertainty is modest compared to the coil-plasma spacings $d_1 \approx 10 - 15$ cm envisioned (see Fig. 3). Consequently, defining the plasma surface as the point where $n \approx 4 \cdot 10^{12} \text{ cm}^{-3}$ should not introduce significant uncertainties into the performance of the resonant cavity coupler beyond that associated with the use of reasonable model plasma density profile.

It is straightforward to redo these calculations for a strictly exponential density profile $n = \bar{n} \exp(x/x_0)$. Definition (2.14) and Eqs. (2.52), (2.53), and (2.55) remain valid, but the definition of the electron density $\bar{n}$ changes to

$$\bar{n} = (5.2 \cdot 10^{12} \text{ cm}^{-3}) \left(\frac{5 \text{ cm}}{x_j} \right)^2 \left(\frac{\Lambda}{2} \right) \quad (2.58)$$

The plasma surface is again defined as the point where $n = 0.3\bar{n}$. The reason why our previous results remain valid rests on the observation that $A/B \approx -i$ [c.f. Eq. (2.26)]. Consequently, expression (2.23) is linearly proportional to the Hankel function $H_0^{(1)}(\xi)$ which is the desired solution of Eq. (2.22). Formula (2.58) follows from the relations

$$\frac{x_o^2 \omega_o^2 \omega_{pi}^2}{c^2 \Omega^2} = 1, \quad \omega_{pi}^2 = \frac{4\pi n_e e^2}{A M_p}$$

required to obtain Eq. (2.22).

III. A RESONANT-CAVITY COUPLER

A. Cavity geometry

The geometrical constraints of the "Biq-Dee" set the size of the opening for an ICRF coupler at 35 cm toroidally and 50 cm poloidally. Figures 1 and 2 depict the geometry of the coupler cavity. It is a rectangular box containing a central parallel-plate capacitor composed of two vertical plates attached to the bottom surrounding a single vertical plate attached to the top. The two bottom plates are movable so that the gap spacing q can be varied to tune the cavity. The capacitance is

$$C = 2\epsilon_o \frac{wh}{q} \quad (3.1)$$

The inductance can be determined from two-dimensional vector potential calculations which yield the inductance-per-unit-length. The key parameter

here is the distance δ between the Faraday screen and the plasma surface (defined in Sec. II). Section II demonstrated that the surface impedance of the plasma is low, permitting us to model the plasma boundary as a perfect conductor for the purposes of computing the cavity inductance. The inductance calculation is based on a two-dimensional approximation where the rf magnetic field $\vec{H}$ is given by

$$\vec{H} = \vec{\nabla}\psi \times \hat{y} \quad (3.2)$$

and the rf flux function (vector potential) ψ satisfies

$$\nabla^2\psi = 0$$

with the boundary conditions being $\psi = \psi_0$ on the central conductor and $\psi = 0$ on the outside boundary. The inductance-per-unit-length is then

$$\frac{\partial L}{\partial y} = \mu_0 \alpha \quad (3.3)$$

where

$$\alpha = \frac{\psi_0}{\oint d\ell (\partial\psi/\partial n)} \quad (3.4)$$

and the line integral encloses the central conductor. The cavity inductance is simply

$$L = \mu_0 \alpha h \quad (3.5)$$

leading to a resonant frequency given by

$$\omega_o^2 = \frac{1}{LC} = \frac{c^2}{2h^2 \alpha} \left(\frac{q}{w} \right)^2. \quad (3.6)$$

Clearly, the cavity resonant frequency depends on the plasma boundary distance δ through its effect on α . We have investigated the sensitivity through three models for δ , which are portrayed in Fig. 3. The results of computations are reported in Table I, based on the values of Table II. For the 60 MHz resonant frequency desired for second harmonic proton heating in Doublet-III and the dimensions of Tables I and II, Eq. (3.6) can be cast as

$$q = (2.3 \text{ cm}) \left(\frac{\alpha}{0.2} \right) \left(\frac{\nu_o}{60 \text{ MHz}} \right)^2 \quad (3.7)$$

We see that gap spacings between 1 - 2 cm are sufficient for tuning the cavity to accommodate various plasma-boundary positions.

B. Cavity - Q

The principal simplifying feature of the cavity-plasma interaction is the low value of the plasma surface impedance Z_p , valid when the parallel and poloidal wavelengths of the impressed H_z field are sufficiently long. Equation (2.56) gives the expression for Z_p . The simple results hold when $q, \beta \ll 1$. Figure 4 shows the parallel and poloidal power spectrum of the Doublet III coupler. Most of the power spectrum does lie at small values of q and β . This will be even more true for reactor couplers which will be larger than the Doublet III coupler.

The low value of the surface impedance means that the radiation from the cavity can be computed via usual skin-effect procedures. First, the plasma

surface can be regarded as a good conductor from the point of view of calculating the cavity inductance. This step has already been taken in the calculations reported in Table I. Secondly, the power radiated by the cavity can be computed via

$$P = h Z_p \int_L H_z^2 dz \quad (3.8)$$

where the integral is along the plasma boundary of Fig. 2c. The numerical inductance calculations have also computed the nondimensional quantity Λ

$$\Lambda = \frac{2\ell}{\psi_0^2} \int_L \left(\frac{\partial\psi}{\partial x}\right)^2 dz = \frac{2\ell \mu_0^2 h^2}{\phi^2} \int H_z^2 dz \quad (3.9)$$

where

$$\phi = \mu_0 a h I \quad (3.10)$$

is the oscillating rf flux. The total stored energy is

$$U = \phi I. \quad (3.11)$$

Combining Eqs. (3.8) - (3.11), one can obtain the following expression for Q

$$\begin{aligned} Q &= \frac{2\omega\ell\eta}{\alpha\Lambda Z_p c} = \left(\frac{6.3}{\alpha\Lambda}\right) \left(\frac{\ell}{35 \text{ cm}}\right) \left(\frac{n_0}{10^{14} \text{ cm}^{-3}}\right)^{1/3} \left(\frac{Z}{\Lambda}\right)^{1/3} \left(\frac{70 \text{ cm}}{a}\right)^{1/3} \\ &= \frac{2}{3.117} \left(\frac{\ell}{x_0}\right) \frac{1}{\alpha\Lambda} \quad (3.12) \end{aligned}$$

Table III presents the results. It is clear that the cavity Q depends

critically on the plasma coil separation d_1 . An excellent approximation is

$$Q = 3.9 \left(\frac{d_1}{10 \text{ cm}} \right)^2 \left(\frac{\ell}{35 \text{ cm}} \right) \left(\frac{n_0}{10^{14} \text{ cm}^{-3}} \right)^{1/3}. \quad (3.13)$$

The increase of Q with d_1 results from the evanescent decoupling that occurs in the vacuum region between the Faraday shield and the plasma. The precise definition of the plasma surface is discussed in Sec. II.

The rms voltage in the cavity is

$$V = \omega \Phi. \quad (3.14)$$

We can express the relation between power and cavity voltage in terms of a parallel resistance $R_{\parallel}$ by

$$P = \frac{V^2}{R_{\parallel}} \quad (3.15)$$

where

$$R_1 = \eta_0^2 \left(\frac{2\ell}{c} \frac{h\omega^2}{2} \right) \frac{1}{Z_p \Lambda}, \quad (3.16)$$

$$R_1 = (1490 \Omega) \frac{1}{\Lambda} \left(\frac{n_0}{10^{14} \text{ cm}^{-3}} \right)^{1/3} \left(\frac{B}{2T} \right). \quad (3.17)$$

The loading for ICRF antennas has customarily been expressed in terms of a series resistance R_s which is defined by

$$P = R_s I^2 \quad (3.18)$$

where I is the current flowing in the cavity. Equations (3.8) - (3.10) combine to give

$$R_s = \frac{\Lambda a^2 Z_h}{2l} \frac{P}{P} \quad (3.19)$$

$$= (37.7 \, \Omega) \Lambda a^2 \left(\frac{10^{14} \text{ cm}^{-3}}{n_0} \right)^{1/3} \left(\frac{R}{2T} \right). \quad (3.20)$$

In formulas (3.17) and (3.20), we have assumed second harmonic heating, so that the frequency scales linearly with the magnetic field.

The loading of the cavity is quite sensitive to the plasma-cavity separation. Moderate Q (10-30) cavities will permit us to match to transmission lines, while not requiring unacceptably large voltages.

Table III shows that with a 10 - cm thick scrape-off plasma, the series resistance remains sufficiently high and the cavity- Q moderate. This thick scrape-off plasma is quite desirable in a reactor to protect the coupler from heat outflow.

C. Matching to transmission lines

The resonant cavity is excited by magnetic loop coupling from a pair of 40 Ω transmission lines. The transmission lines are of relatively low impedance to permit high power capability at acceptable voltages. The rectangular geometry is chosen both to result in low self-inductance for the coupling loop [see Eq. (3.25)] and to provide magnetic insulation across the smaller conductor separation at the vacuum-gas interface which is presumed to occur in a region of significant toroidal field.

The equivalent circuit for the matching network is shown in Fig. 5. This

leads to the equations

$$0 = (-i\omega L_C + \frac{i}{\omega C} + R_S)I_C - i\omega M I_L \quad , \quad (3.21)$$

$$2V_L = -i\omega L_L I_L - i\omega M I_C \quad , \quad (3.22)$$

and the well-known formula

$$Q = \frac{\omega L_C}{R_S} = \frac{R_{||}}{\omega L_C} \quad . \quad (3.23)$$

Equations (3.21) - (3.23) employ the following definitions which use the dimensions given in Figs. 2 and 3.

$$L_C = \mu_0 \alpha h = \text{cavity inductance} \quad , \quad (3.24)$$

$$L_L = \frac{\mu_0 s \Delta}{z} = \text{self inductance of coupling loop} \quad , \quad (3.25)$$

$$M = \frac{4\mu_0 \alpha s \Delta}{\pi \sinh(\pi d_2/l)} = \text{mutual inductance between cavity and loop} \quad , \quad (3.26)$$

$$C = \frac{2\epsilon_0 wh}{q} = \frac{1}{\omega_0^2 L_C} = \text{capacitance of cavity} \quad , \quad (3.27)$$

$$I_C = \text{current in cavity} \quad ,$$

$$I_L = \text{current in transmission line} \quad ,$$

$$V_L = \text{voltage in a transmission line} \quad ,$$

ω_0 = resonant frequency of cavity ,

$R_L = 40 \Omega$ = impedance of transmission line .

Let us comment on the expressions for L_L and M . The self-inductance of the coupling coil is simply that of a solenoid of area $s\Delta$ and length z . Formula (3.25) probably underestimates L_L because of a contribution from volume in the transmission lines. The mutual inductance is based on the result

$$B_M = \frac{4\Phi}{\ell h \sinh(\pi d_2/\ell)} = -\frac{4\mu_0 I_C}{\ell \sinh(\pi d_2/\ell)} = \begin{array}{l} \text{magnetic field near backwall of} \\ \text{cavity} \end{array} \quad (3.28)$$

which, in turn, rests on the approximation that the cavity rf flux Φ must pass through a relatively narrow gap between the central conductor and the cavity sidewall. The mutual flux Φ_M coupled by the transmission line is then

$$\Phi_M = B_M s\Delta \equiv M I_C \quad (3.29)$$

from which the mutual inductance follows. It is also useful to note that $L_L = \beta M$, where

$$\beta = \frac{\ell}{4az} \sinh(\pi d_2/\ell) \approx 3.4 . \quad (3.30)$$

The matching calculation is performed by solving Eqs. (3.21) and (3.22) for the ratio

$$\frac{V_L}{I_L} = \left(\frac{1}{2}\right) \left[\frac{\omega M^2}{Q L_c \left[\left(1 - \omega_o^2 / \omega^2\right)^2 + 1/Q^2 \right]} \right. \\ \left. - i \left[\omega L_L - \frac{\omega M^2 \left(1 - \omega_o^2 / \omega^2\right)}{L_c \left[\left(1 - \omega_o^2 / \omega^2\right)^2 + 1/Q^2 \right]} \right] \right] = Z_L$$

A perfect match to the transmission line requires $Z_L = R_L$. This is a complex equation which generates two matching equations

$$2R_L = \frac{\omega M^2}{Q L_c \left[\left(1 - \omega_o^2 / \omega^2\right)^2 + 1/Q^2 \right]} \quad (3.31)$$

$$\omega L_L - \frac{\omega M^2 \left(1 - \omega_o^2 / \omega^2\right)}{L_c \left[\left(1 - \omega_o^2 / \omega^2\right)^2 + 1/Q^2 \right]} = 0 \quad (3.32)$$

There are two adjustable parameters which permit these equations to be satisfied. They are ω_o , the cavity resonant frequency, and L_c , the distance the coupling loop penetrates into the back of the cavity. The cavity resonant frequency is varied by changing the capacitor gaps.

Equations (3.31) and (3.32) can be manipulated to give

$$\frac{3\omega M}{2RLQ} = \left[\frac{\beta^2 \omega_L^2}{2R_L \left(Q - \beta^2 \omega_L^2 / 2R_L \right)} \right]^{1/2} \frac{1}{Q} = 1 - \frac{\omega_o^2}{\omega^2} \quad (3.33)$$

$$M^2 = \frac{2R_L L_c}{\omega \left(Q - \beta^2 \omega_L^2 / 2R_L \right)} \quad (3.34)$$

For meaningful solutions to exist, Q must be sufficiently large:

$$Q > Q_0 = \frac{\beta^2 \omega_L c}{2R_L} = 5.1 \quad . \quad (3.35)$$

By combining Eqs. (3.24), (3.26), (3.34), and (3.35) we can obtain an expression for Δ

$$\frac{\Delta}{z} = \frac{2R_L}{\omega \mu_0} \left(\frac{Q_0}{Q - Q_0} \right)^{1/2} = (0.50) \left(\frac{Q_0}{Q - Q_0} \right)^{1/2} \quad . \quad (3.36)$$

Cavity Q values above 10 are required to make Δ less than d_2 , which is implied in the configuration of Fig. 2. Thus the resonant coil coupler will work best when the cavity loading is not too high, i.e., when the plasma-wall separation exceeds 10 cm. Under these circumstances, a perfect match to the transmission line can be obtained and the required shift of the resonant frequency

$$\frac{\Delta\omega}{\omega_0} \approx \frac{1}{2Q} \left(\frac{Q_0}{Q - Q_0} \right)^{1/2} \quad (3.37)$$

is small. Clearly, the mechanical design of the coupling system will have to permit a variation in distance Δ to achieve a match to the transmission line. The range $5 < \Delta < 10$ cm should suffice.

D. Power handling capability

The power handling capability is governed by the electric fields that can be maintained without breakdown. There are two different criteria: the first criterion governs the electric field E_1 that can be sustained along the insulator at the vacuum-gas interface which occurs in the transmission lines. It is assumed that this transition takes place in a region of significant toroidal magnetic field so the relevant electric field is that

between the central conductor and the side wall of the rectangular coaxial line. The second criterion involves the electric fields which can be withstood across a vacuum gap. There are two classes here, E_{21} , the electric field perpendicular to the main toroidal field and $E_{2\parallel}$, the electric field parallel to the magnetic field.

The paper of Hwang, Grotz, and Hosea¹¹ provides data for both criteria. Their measurements for vacuum breakdown are for a magnetic-field-free case. Table IV summarizes their measurements and the criteria we adopt for the measurements. We have assumed that magnetic insulation will be sufficiently effective so that the breakdown electric field across the toroidal field is significantly higher than the parallel breakdown field. The criteria are all significantly less than the measured values and thus should be conservative. Indeed, experience with all-metal ICRF antennas in the TFR tokamak has shown that 50 kV/cm across the toroidal field can be withstood even when the antenna is inside the vacuum vessel and immersed in the scrape-off plasma.

Table V shows the voltages and electric fields which are associated with the Doublet III coupler operated at 20 MW. The electric fields are all at or below the criteria of Table IV. In particular, the electric fields in the transmission lines are below the criteria, so that a significant reflected power could be tolerated without inducing breakdown. A power of 20 MW corresponds to a power flux of 10 kW/cm^2 - a value that is very attractive for heating fusion plasmas.

The keys to excellent power handling capability are matching the transmission line to reduce voltages at the vacuum-gas interface insulator, and orienting the strong electric fields in the capacitor orthogonal to the toroidal field.

IV. DISCUSSION

ICRF antennas in the next generation of tokamaks will differ significantly from the induction loops used presently. They will be recessed behind the vacuum vessel wall, and the power will be transmitted through a large area Faraday screen. This work shows that these couplers will bring about an impressive improvement in performance - power fluxes of 10 kW/cm^2 can be obtained without exceeding conservative electric field criteria. The large area Faraday screen not only shields out coupling to surface electrostatic modes, but also prevents plasma particles and radiation from entering the resonant cavity. In addition, it permits toroidal currents to flow uninterrupted in the vacuum vessel wall, and thus does not degrade the passive stabilization properties of the vacuum vessel. The resonant-cavity is all metal and of a modular nature which can be simply bolted onto a rectangular port opening. The insulators associated with the vacuum gas transition in the coaxial transmission could be located behind neutron shielding. Altogether, the resonant-cavity coupler is an attractive reactor prototype.

A key finding of this work is that gradual decay of density in the scrape-off plasmas, which we have modeled by an exponential profile, increases the plasma surface impedance by roughly a factor-of-three over previous results¹⁶ for a strictly linear plasma density model. Straightforward calculations for the linear model yield

$$Z_p = 1.19 \left(\frac{\omega_p}{c} \right) \eta_0 \quad (4.1)$$

This expression can be easily obtained by evaluating Eq. (2.9) at $v = 0$, and using Eqs. (2.7), (2.54), and (2.55). Comparison of Eq. (4.1) with Eq. (2.56) points out the increase in surface impedance associated with exponential

profile.

The plasma surface impedance was determined by a model which used outgoing wave boundary conditions. This is an accurate representation, provided not much wave energy is reflected back to the region of the antenna. Many second harmonic heating calculations should have significant single-pass absorption for the fast Alfvén wave, thereby justifying the outgoing wave boundary conditions. For fast-wave current drive where the single-pass absorption is expected to be weak, the plasma surface impedance must be reexamined.

Our calculation has been based on rather simple models for inductances, capacitance, etc. The first step experimentally is to build a 1/3-size model and at three times the frequency, provide a high dielectric constant medium (e.g., water) into which the coupler can launch waves. The cavity-dielectric medium spacing could be varied. Thus, a rather complete engineering verification of the resonant cavity coupler could be obtained before the first full-scale tokamak tests. If the model coupler were placed in a region of high magnetic field, the benefits of magnetic insulation could also be verified.

At this stage, the resonant cavity ICRF coupler appears to possess many attractive properties. These properties must be proved by engineering tests and actual experience on tokamaks.

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REFERENCES

- ¹ J. Hosea et al., "High Power ICRF Heating on PLT and Extrapolation to Future Devices," in Proceedings of the Third Joint Varenna-Grenoble International Symposium on Heating in Toroidal Plasmas, Grenoble, 1982, Vol. I, pp. 213-224.
- ² Equipe TFR, "ICRF Heating and Antenna Coupling Studies in TFR" in Proceedings of the Third Joint Varenna-Grenoble International Symposium on Heating in Toroidal Plasmas, Grenoble, 1982, Vol. I, pp. 225-242.
- ³ JFT-2 Group, "ICRF Heating Experiment in JFT-2" in Proc. Third Joint Varenna-Grenoble International Symposium on Heating in Toroidal Plasmas, Grenoble, 1982, Vol. I, pp. 259-271.
- ⁴ F. W. Perkins, "Heating Tokamaks via the Ion-Cyclotron and Ion-Ion Hybrid Resonances," Nucl. Fusion 17, 1197 (1977).
- ⁵ J. Jacquinet, B. D. McVey, and J. E. Scharer, "Mode Conversion of the Fast Magnetosonic Wave in a Deuterium-Hydrogen Tokamak Plasma," Phys. Rev. Lett. 39, 88 (1977).
- ⁶ R. R. Weynants, "Ion-Heating at Twice the Ion-Cyclotron Frequency in Reactor-Oriented Machines," Phys. Rev. Lett. 33, 78 (1974).
- ⁷ T. H. Stix, "Fast-Wave Heating of a Two-Component Plasma," Nucl. Fusion 15, 737 (1975).
- ⁸ TFR Group, A. Truc, and D. Gresillon, "Observations of Fast and Slow Modes Associated With ICRF Heating in the TFR Tokamak by CO₂ Laser Scattering," Nucl. Fusion 22, 1577 (1982).
- ⁹ W. M. Stacey et al., "U.S. FED-INTOR Activity and U.S. Contribution to the International Tokamak Reactor Phase-2A Workshop," Georgia Institute of Technology Report USA FED-INTOR/82-1 (1982).

- ¹⁰ P. H. Rebut, "JET Joint Undertaking: March 1982," in Proceedings of the Third Joint Varenna-Grenoble International Symposium Heating in Toroidal Plasmas, Grenoble, 1982, Vol. III, pp. 989-998.
- ¹¹ D. Q. Hwang, G. Grotz, and J. C. Hosea, "Surface Physics Problems During ICRF Heating of Tokamak Plasmas," J. Vac. Sci. Technol. 20, 1273 (1982).
- ¹² L. G. Davis and J. M. Rawls, "Modification of Doublet III to a Large Dee Facility," Proceedings of Ninth Symposium on Engineering Problems in Research (IEEE, New York, 1981), Vol. II, pp. 1939-1942.
- ¹³ H. P. Furth, "Present and Future Toroidal Confinement Experiments at the Princeton Plasma Physics Laboratory," Physica Scr. (to be published).
- ¹⁴ M. P. Evrard and R. R. Weynants, "ICRH Antenna Coupling to Slow Waves due to the Poloidal Field," in Proceedings of Third Joint Varenna-Grenoble International Symposium Heating in Toroidal Plasmas, Grenoble, 1982, Vol. I, pp. 339-346.
- ¹⁵ M. Abramowitz and I. Stegun, Handbook of Mathematical Functions (U.S. Government Printing Office, 1955).
- ¹⁶ F. W. Perkins, "ICRF Couplers," Bull. Am. Phys. Soc. 26, 929, (1981).

Table I. Resonant Cavity Parameters

Distance to Plasma Boundary δ (cm)	d_1 (cm)	α	q (cm)
10	15	0.164	1.88
5	10	0.152	1.74
0	5	0.121	1.39

Table II. Parameters for Doublet-III Cavity Couplers

$\nu = 60 \text{ MHz} = 2^{\text{nd}}$ harmonic proton gyrofrequency at 2T

$h = 50 \text{ cm} = \text{poloidal cavity height}$

$l = 35 \text{ cm} = \text{toroidal cavity length}$

$w = 15 \text{ cm} = \text{width of center conductor}$

$d_2 = 10 \text{ cm} = \text{center conductor - backwall spacing}$

$a = 70 \text{ cm} = \text{plasma minor radius}$

$t = 6 \text{ cm} = \text{overall thickness of center conductor}$

$z = 17 \text{ cm} = \text{width of coaxial conductor}$

$s = 35 \text{ cm} = \text{length of cavity coupling loop}$

$\Delta = 5\text{-}10 \text{ cm} = \text{adjustable distance between cavity coupling loop and cavity backwall}$

Table III. Cavity Q, Resistances.

δ (cm)	a_1 (cm)	α	Λ	Q	R_{f} (Ω)	R_{s} (Ω)
10	15	0.164	3.78	10.0	395	3.84
5	10	0.152	10.74	3.9	139	9.35
0	5	0.121	49.52	1.1	30	27.4

Table IV. Electric Field Criteria

Case	Ref. 11	Criterion
Insulator at vacuum-gas interface	10 kV/cm	5 kV/cm
Vacuum gap across toroidal field	120 kV/cm*	50 kV/cm
Vacuum gap along toroidal field	120 kV/cm*	20 kV/cm

*Measurement at zero magnetic field

Table V. Voltages For 20 MW Coupler.*

Element	Power (MW)	Gap Spacing (cm)	$R_{ }$ (Ω)	V (kV)	E (kV/cm)	Criterion (kV/cm)
Single Transmission Line	10	9	40	20	2.2	5
Capacitor gap (across toroidal field)	20	1.8	395	90	50 kV/cm	50 kV/cm
Capacitor-side- wall (along toroidal field)	20	10	395	90	9 kV/cm	20 kV/cm

* $\delta = 10$ cm, $\phi = 10$ assumed (see Table III).

FIGURE CAPTIONS

FIG. 1. a) Sketch of resonant cavity and transmission line installation. b) Orientation of the rf electric $\vec{E}$ and magnetic fields $\vec{B}$ transmitted through Faraday screen. The electric and magnetic fields are identical in each slot. $\vec{B}_0$ is the steady toroidal field. c) Coordinate system and density profile for wave propagation calculations in the plasma periphery. At density $\bar{n}$ [see Eq. (2.11)], the profile switches from an exponential to a linear dependence on x .

FIG. 2. Geometry of the Big-Dee Doublet III ICRF Coupler. a) Overall view: the large area Faraday screen coincides with the vacuum vessel wall. b) Cross section of rectangular coaxial line. c) Cross section of resonant cavity with parallel-plate central conductor. d) Side view of cavity with loop coupling from coaxial lines.

FIG. 3. Geometry for inductance calculations. a) Side view of Cavity. b) Cross section of cavity. c) Two-dimensional boundary value problem for rf flux function ψ . The results of the calculations are very insensitive to the precise location of surfaces A and B provided they are separated from the cavity by a distance of at least δ .

FIG. 4. Wave number spectra of H_z at the plasma boundary of Fig. 2c. The results are almost independent of cavity-plasma separation δ . The dashed curve is the q-spectrum and the solid curve is the β -spectrum.

FIG. 5. Equivalent circuit for transmission line-cavity coupling.

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Fig. 2

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Fig. 3

Fig. 4

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Fig. 5

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