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**ORNL TNS Program: Microwave Start-up of
Tokamak Plasmas Near Electron Cyclotron and
Upper Hybrid Resonances**

Y-K. M. Peng
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OAK RIDGE NATIONAL LABORATORY
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FUSION ENERGY DIVISION

ORNL TNS PROGRAM:

MICROWAVE START-UP OF TOKAMAK PLASMAS NEAR
ELECTRON CYCLOTRON AND UPPER HYBRID RESONANCES

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ABSTRACT

The scenario of toroidal plasma start-up with microwave initiation and heating near the electron cyclotron frequency is suggested and examined here. We assume microwave irradiation from the high field side and an anomalously large absorption of the extraordinary waves near the upper hybrid resonance. The dominant electron energy losses are assumed to be due to magnetic field curvature and parallel drifts, ionization of neutrals, cooling by ions, and radiation by low Z impurities. It is shown by particle and energy balance considerations that electron temperatures around 250 eV and densities of 10^{12} - 10^{13} cm⁻³ can be maintained, at least in a narrow region near the upper hybrid resonance, with modest microwave powers in the Impurity Study Experiment (ISX) (120 kW at 28 GHz) and The Next Step (TNS) (0.57 MW at 120 GHz). The loop voltages required for start-up from these initial plasmas are also estimated. It is shown that the loop voltage can be reduced by a factor of five to ten from that for unassisted start-up without an increase in the resistive loss in volt-seconds. If this reduction in loop voltage is verified in the ISX experiments, substantial savings in the cost of power supplies for the ohmic heating (OH) and equilibrium field (EF) coils can be realized in future large tokamaks.

1. INTRODUCTION

Recent tokamak start-up studies [1, 2] have indicated that the loop voltage, V_ℓ , for full-bore plasma initiation and current buildup scales linearly with the major radius. For tokamak fusion reactors, V_ℓ is estimated to be 300-500 V [1]. Because of the simultaneous high peak currents required in the OH primary windings, this implies large pulsed power supplies of a few gigawatts for a small fraction of a second. The estimated cost of the generators, switches, and energy storage and transfer systems is currently estimated to be around \$150 million [3]. Major research efforts are likely to be initiated toward the development of such power supplies and energy storage systems [4].

There is, therefore, a strong economic incentive to explore methods of reducing the power supply requirements. Various auxiliary start-up approaches have been discussed in the past. These include the small-radius start-up schemes with magnetic limiters [5] against a poloidal field divertor [6] and against a mechanical limiter [7]. These approaches have the potential advantage of reducing the impurity influx and avoiding current skin effects in large tokamaks, but do not substantially alter the required peak V_ℓ . Initiation at the radio frequencies has been widely used to obtain initial plasmas of a few electron volts. However, it has been shown to have relatively small effects on the required V_ℓ [2].

Recent modeling of start-up processes in the Oak Ridge Tokamak (ORMAK) [1] and the Princeton Large Torus (PLT) [2] has also shown that radiation losses due to minute low Z impurities dominate the plasma energy balance when the electron temperature is below 50-100 eV during start-up. To reach such a temperature, the resistive power input to the low current

plasma has to be roughly twice that during the high current, steady-state phase. (A summary of the understandings on plasma start-up from a recent workshop on OH power supply requirements in tokamak reactors is included in the Appendix.) It is seen that in order to substantially reduce the loop voltage, some auxiliary plasma heating mechanism is needed to "burn through" the low Z impurity radiation barrier. A reduced V_ℓ will then be sufficient to build up the plasma current to its steady-state value. Such a method, based on microwave irradiation near the upper hybrid resonance (UHR) and electron cyclotron resonance (ECR) frequencies, is the subject of this study.

The microwave breakdown of low pressure hydrogen and the anomalous absorption by the resulting plasma in the toroidal magnetic field have been studied experimentally by Anisimov et al. [8]. In these experiments, powers up to 50 kW at 9.7 GHz were deposited from the high field side, primarily as extraordinary waves, into a toroidal chamber with major and minor radii of 22 cm and 6 cm, respectively. Electron energies around a few hundred electron volts were observed at powers above several kilowatts. In the ELMO Bumpy Torus (EBT) experiments [9], steady-state electron temperatures of several hundred electron volts have been maintained with roughly 50 kW of microwave power. In ISX, continuous wave sources at 200 kW and 28 GHz will be made available for ECR heating experiments. There is immediate interest in studying initiation and heating processes of a currentless toroidal plasma near the ECR.

In the following, the interaction between the extraordinary wave and the plasma near the ECR is briefly summarized in Section 2. A global model for the particle and energy balance in a currentless toroidal plasma is

described in Section 3. Its application to ISX and a tokamak ignition reactor (e.g., TNS) to obtain high electron temperatures ($T_e \sim 250$ eV) is examined in Section 4. The required loop voltages for start-up from such initial plasmas are estimated in Section 5. Our results are summarized in Section 6, where the possible effects of an ambipolar potential on the plasma equilibrium and the influences of macroscopic instabilities on confinement are also discussed.

2. INTERACTIONS BETWEEN EXTRAORDINARY WAVE AND PLASMA — A BRIEF SUMMARY

We assume that a steady-state plasma coexists with an input extraordinary wave from the inside wall of the toroidal chamber. Due to the relatively poor electron confinement time, the plasma is assumed to be a partially ionized gas except possibly where significant wave absorption occurs.

As indicated in Ref. 8, two types of wave-plasma interactions are observed, depending on whether the electron density satisfies the condition

$$\frac{\omega_{pe}^2}{\omega_o^2} \gtrsim 2 \left(\frac{\omega_c}{\omega_o} - 1 \right)_r \quad (1)$$

Here, $\omega_{pe} (= \sqrt{4\pi n_e e^2 / m_e} \approx 5.64 \times 10^4 \sqrt{n_e})$ is the electron plasma frequency, ω_o is the applied wave frequency, and $\omega_c (= eB / m_e c \approx 1.76 \times 10^7 B)$ is the electron cyclotron frequency.* The subscript r represents the boundary of strong ECR interaction. For filling pressures less than 10^{-4} torr and

* Except where otherwise noted, cgs units are used and temperatures are given in electron volts.

$T_e \sim 10$ eV, relativistic detuning [10, 11] of the resonant electrons determines this boundary. For a root-mean-squared (rms) vacuum wave field $\tilde{E}_0$ of roughly 10^3 V/cm, the right-hand side of Eq. (1) is roughly 0.04 [11]. For sufficiently high T_e ($\gtrsim 200$ eV), the resonance boundary is then determined by the Doppler shift due to the thermal electron velocity. In this case, one has [12]

$$\left(\frac{\omega_c}{\omega_0} - 1 \right)_r \approx \left(\frac{N_{\parallel}^2 T_e}{m_e c^2} \right)^{1/2} \quad (2)$$

where $N_{\parallel}$ is the parallel refractive index having values near unity.

2.1 LOW DENSITY REGIME

Equation (1) represents the condition under which the quasi-extraordinary mode no longer has the correct polarization to accelerate electrons at the cyclotron resonance [13]. When Eq. (1) is violated at sufficiently low densities, strong ECR interaction occurs around $\omega_0 = \omega_c$ giving an electron energy of [10]

$$W_{LR} = 2m_e c^2 \left(\frac{e\tilde{E}_0}{m_e \omega_0 c} \right)^{2/3} \approx 1.54 \times 10^9 \left[\frac{\tilde{E}_0 \text{ (V/cm)}}{\omega_0} \right]^{2/3} \text{ (eV)} \quad (3)$$

Away from the resonance region, the electron energy associated with the "sloshing" by the oscillating electric field is then

$$W_{LNR} = e^2 \tilde{E}_0^2 / m_e \omega_0^2 \left(\frac{\omega_c}{\omega_0} - 1 \right)^2 \approx 1.76 \times 10^{15} \tilde{E}_0^2 \left(\frac{\text{V}}{\text{cm}} \right)^2 / \omega_0^2 \left(\frac{\omega_c}{\omega_0} - 1 \right)^2 \text{ (eV)} \quad (4)$$

Typically, with $\tilde{E}_0 \sim 10^3$ V/cm at $\omega_0 \approx 1.76 \times 10^{11} \text{ sec}^{-1}$ (28 GHz), one has $W_{LR} \sim 5$ keV while $W_{LNR} \sim 12$ eV at $|(\omega_c/\omega_0) - 1| \approx 0.07$.

2.2 HIGH DENSITY REGIME

When Eq. (1) is satisfied at sufficiently high densities, the absorption of the extraordinary wave at the ECR region becomes negligible. The wave then linearly accumulates near the UHR region where

$$\omega_o = \omega_h \equiv (\omega_{pe}^2 + \omega_c^2)^{1/2} \quad (5)$$

with a concomitant linear amplification of the wave $\tilde{E}$ [14]. For small $\tilde{E}$ such that the linear wave theory applies, the extraordinary mode is expected to go through linear transformation into the electron Bernstein mode [15].

On the other hand, if $\tilde{E}$ near the UHR becomes larger than $\sim 10^2$ - 10^3 V/cm, parametric decay into ion-acoustic waves and upper hybrid (electron Bernstein) waves is likely to occur [16], primarily heating the electrons. Since various ECR heating experiments at low density and temperature [8, 17, 18] have resulted in bulk electron heating near the UHR, we shall assume that this is the dominant wave absorption process for the high density regime to be considered here.

2.3 REGION OF MICROWAVE POWER DEPOSITION

The spatial region in which this anomalous absorption process occurs can be estimated according to the rate of increase of $\tilde{E}$ as the wave approaches the UHR zone. It is well known that $\tilde{E}$ of the quasi-extraordinary mode satisfies [19]

$$\frac{i\tilde{E}_R}{\tilde{E}_Z} = \frac{N_L^2 - S}{\tilde{D}}, \quad N^2 = N_L^2 + N_{\parallel}^2, \quad N_L^2 = \frac{\tilde{R}L}{S},$$

$$S = \frac{\omega_o^2 - \omega_h^2}{\omega_o^2 - \omega_c^2}, \quad \tilde{D} = -\frac{\omega_c \omega_{pe}^2}{\omega_o(\omega_o^2 - \omega_c^2)}, \quad (6)$$

$$\tilde{R}L = 1 - \frac{\omega_{pe}^2(2\omega_o^2 - \omega_{pe}^2)}{\omega_o^2(\omega_o^2 - \omega_c^2)}$$

where N_L is the perpendicular refractive index and $\tilde{E}_R$ and $\tilde{E}_Z$ are the components of $\tilde{E}$ in the axisymmetric coordinates (R, ϕ, Z) . Note that $\tilde{E}_Z$ remains roughly constant in R , and is similar in size to the vacuum wave field. For $\omega_o^2 \sim \omega_c^2 \sim \omega_h^2 \gg \omega_{pe}^2$ and according to Eq. (6), the location at which E_R is linearly enhanced by a factor of three occurs where $\omega_h^2 - \omega_o^2 \approx \omega_{pe}^2/8$. This means that substantial increase in E_R occurs only at a distance $\Delta_{UH} \sim (R_o/16) \times (\omega_{pe}/\omega_o)^2$ away from the UHR surface. Since the distance between the ECR and the UHR is $8\Delta_{UH}$, one can assume that the anomalous absorption of the microwave power p_{μ} in general occurs in a narrow region of width Δ_{UH} near the UHR surface. This is apparently supported by experimental observations [8].

The assumed interactions between the input extraordinary wave and the currentless toroidal plasma in high and low density regimes are depicted in Fig. 1. Note that in either regime, the wave is in general absorbed in one pass through the chamber. Particle and energy balance considerations are used in the next two sections to estimate n_e and T_e near and in the region where strong wave absorptions occur.

3. MODEL OF PARTICLE AND ENERGY BALANCE

3.1 PARTICLE BALANCE

The principal particle loss mechanisms are assumed to be the magnetic curvature drift, v_D , and the parallel drift, v_δ , when a small poloidal field "error," δB_p , exists. With the electron ionization as the primary particle source, we have

$$\frac{dn_e}{dt} = n_e n_o S_H - \frac{n_e}{\tau_D} - \frac{n_e}{\tau_\delta} = 0 \quad (7)$$

where n_o is the atomic hydrogen density and S_H is the net ionization rate coefficient [1, 2]. The particle drift velocities and times are given by

$$v_D = \frac{cm_e}{eBR} \left(v_\parallel^2 + \frac{v_\perp^2}{2} \right) \approx \frac{2cT}{eBR} \quad (8)$$

$$\tau_D = \frac{b}{v_D} \approx \frac{bRB}{2 \times 10^8 T}$$

$$v_\delta = v_\parallel \frac{\delta B_p}{B}, \quad \tau_\delta = \frac{b}{v_\delta} \quad (9)$$

Here, b is the chamber minor radius in the vertical direction, B is the toroidal magnetic field, and $v_\parallel^2 \approx v_\perp^2 \approx v^2/3$ has been assumed. From Eqs (8) and (9), one sees that if δB_p satisfies

$$\delta B_p > \sqrt{3} \frac{m_e cv}{eR} = \frac{c}{eR} \sqrt{6m_e T_e} \approx 5.84 \frac{\sqrt{T_e}}{R} \quad (10)$$

loss due to v_δ cannot be neglected. For $R \approx 90$ cm and $T_e \sim 100$ eV, Eq. (10) for electrons becomes $\delta B_p > 0.65$ G, which is comparable to the earth's field. Therefore, it could be necessary to provide field-cancelling coils to reduce v_δ and ensure an electron confinement time near τ_D .

As will be seen below, the ion temperature is comparable to the neutral temperature. The ions will then have a much longer confinement time than the electrons, implying the existence of an ambipolar electric field. The effects of this field will be discussed in Section 6.

With a given T_e , only n_o is determined by Eq. (7) in a steady state. The electron density, n_e , is determined indirectly by assuming some given filling pressure.

3.2 ENERGY BALANCE

The ion temperature is determined by the balance between heating by electrons and charge-exchange loss,

$$\frac{3}{2} n_e (T_i - T_o) n_o S_{cx} = Q_{ie} = \frac{3m_e n_e (T_e - T_i)}{m_i \tau_{ie}} \quad (11)$$

where T_o is the neutral temperature and τ_{ie} is the electron-ion collision time. We have [20]

$$\begin{aligned} \frac{T_i - T_o}{T_e - T_i} &= \frac{16}{3} \frac{\sqrt{2\pi m_e} \lambda_{ei} e^4}{m_i S_{cx} T_e^{3/2}} \left(\frac{n_e}{n_o} \right) \\ &\approx \begin{cases} 5.74 \\ 2.02 \end{cases} T_e^{-3/2} \left(\frac{n_e}{n_o} \right) \begin{cases} \text{for } T_e \sim 10 \text{ eV} \\ \text{for } T_e \sim 100 \text{ eV} \end{cases} \quad (12) \end{aligned}$$

If (n_e/n_o) is roughly unity for $T \sim 10$ eV and is roughly 10^2 for $T \sim 100$ eV, the right-hand side of Eq. (12) has the values approximately 0.15 and 0.20, respectively. We will thus assume that $T_i \approx T_o$ in the following.

The dominant electron energy losses are assumed to result from ionizations of the neutrals; collisions with ions, Q_{ie} ; VB and δB_p drifts; and impurity radiations, p_{rad} , and ionizations, p_{ion} . We have,

$$\frac{d}{dt} \left(\frac{3}{2} n_e T_e \right) = 0$$

$$= p_\mu - W_{ion} n_e n_o S_H - Q_{ie} - \frac{3}{2} n_e T_e \left(\frac{1}{\tau_D} + \frac{1}{\tau_\delta} \right) - p_{rad} - p_{ion} \quad (13)$$

Here, W_{ion} is assumed to include the energy losses by hydrogen ionization, dissociation, and radiation. It is set to be 30 eV [21]. The quantities p_{rad} and p_{ion} are assumed to be due to oxygen and have been discussed in Ref. 22.

The magnitude of p_μ in Eq. (13) should be discussed here. When the wave is anomalously absorbed near the UHR, p_μ can be approximated by the wave energy flux divided by Δ_{UH} . This then determines T_e according to Eq. (13) for a given n_e . Away from UHR, however, electrons reach energies given by Eqs (3) and (4) within a time scale much shorter than the energy loss time scales shown in Eq. (13). So T_e should be assumed according to Eqs (3) and (4). With n_e and n_o determined from Eq. (7), p_μ is then given by Eq. (13). This p_μ should be interpreted as the absorbed wave power density, i.e., the local wave damping rate.

In ISX ($a = b = 25$ cm, $R = 92$ cm, $B = 10$ kG) with $n_e \sim 0.5 \times 10^{12}$ cm $^{-3}$ and $\tilde{E}_o \sim 500$ V/cm, Eq. (4) gives $W_{1NR} \sim T_e \sim 3$ eV at $|\omega_c/\omega_o - 1| \sim 0.07$. From Eqs (8) and (9), we see that $\tau_\delta \sim 8.4$ msec ($\delta B_p \sim 0.5$ G) and $\tau_D \sim 38$ msec. One finds that $(3n_e T_e / 2\tau_\delta) \sim 4 \times 10^{-5}$ W/cm 3 , which is less than $(p_{rad} + p_{ion}) \sim 2.2 \times 10^{-4}$ W/cm 3 with 2% oxygen. The total wave damping rate is negligible. At the ECR, however, Eqs (3), (8), and (9) give $W_{1R} \sim T_e \sim 3.1$ keV, $\tau_D \sim 37$ μ sec and $\tau_\delta \sim 260$ μ sec. We have $(2n_e T_e / 2\tau_D) \sim 10$ W/cm 3 . This is a substantial wave damping rate. However, when compared with the vacuum wave flux of $(c\tilde{E}_o^2/8\pi) \sim 330$ W/cm 2 , it indicates a damping

length (~ 33 cm) much larger than the width of the ECR region [$\Delta_{EC} \sim 4$ cm according to Eq. (1)]. As a result, most of the wave energy will reach the UHR region in a sufficiently large toroidal chamber, where linear or non-linear conversions to electron Bernstein modes occur.

It should be noted that for fixed $\tilde{E}_0$ and n_e in the low density regime, Eqs (3) and (8) show that the wave energy damping rate in the ECR is proportional to $B^{1/3} (bR)^{-1}$. Reducing b to 6 cm, R to 22 cm, and B to 3 kG from the ISX values results in a wave energy absorption rate of 120 W/cm^3 , making it necessary to avoid the low density regime in the experiments by Anisimov et al. [8]. Increasing the tokamak linear size beyond ISX by a factor of five and the toroidal field to above 40 kG would make the wave energy absorption negligible in the ECR region even in the low density regime.

4. APPLICATIONS TO ISX AND TNS

To estimate an upper boundary on the required microwave power, we first assume an input power density such that a high energy electron at ECR on the average ionizes several neutrals before escaping. With $\delta B_p \sim 0.5$ G and $T_e > 1.0$ keV, τ_δ is much larger than τ_D . Combining Eqs (3) and (7), one obtains

$$\tilde{E}_0 (\text{V/cm}) < 5.85 \times 10^{-27} \omega_0 (bR n_0 S_H)^{3/2} \quad (15)$$

where the value of the ionization rate coefficient S_H is around $2 \times 10^{-8} \text{ cm}^3/\text{sec}$ at $T_e \sim 1$ keV [20]. So for a filling pressure of 5×10^{-5} torr ($n_0 = 1.7 \times 10^{12} \text{ cm}^{-3}$ before substantial ionization) in ISX, one finds an upper bound of $\tilde{E}_0 = 670 \text{ V/cm}$, beyond which the ECR electrons become uncoupled from the neutrals. Here we shall use $\tilde{E}_0 \sim 300 \text{ V/cm}$

to avoid generating many runaway electrons. Assuming that the microwave shines on an area of 10^3 cm^2 near the ECR region, we find that the total wave power input is roughly 120 kW. Note that in bigger tokamaks, runaway electrons become progressively easier to avoid.

With this power input, consider the region where $|(\omega_c/\omega_o) - 1| \approx 0.03$ so that $W_{\text{INR}} \sim 10 \text{ eV}$. We have $\tau_D \sim 12 \text{ msec}$ and $\tau_\delta \sim 4 \text{ msec}$. With $S_H \sim 6 \times 10^{-9} \text{ cm}^3/\text{sec}$, the steady-state neutral density determined from Eq. (7) is $n_o \sim 5 \times 10^{10} \text{ cm}^{-3}$. For a filling pressure of molecular hydrogen of $5 \times 10^{-5} \text{ torr}$ with 1% oxygen, the electron density is $3.2 \times 10^{12} \text{ cm}^{-3}$, putting the plasma in the high density regime according to Eq. (1). The dominant energy loss rate in Eq. (13) is $Q_{ie} \sim 0.064 \text{ W/cm}^3$ (assuming $T_i \approx 0$), while the total energy loss rate is less than 0.071 W/cm^3 . This would amount to a total loss of roughly 77 kW if T_e and n_e are uniform throughout the ISX chamber. However, Eq. (4) shows that $T_e > 10 \text{ eV}$ should be confined within a region of roughly 6 cm around the ECR region. The corresponding power loss is then only about 12 kW. The remaining wave energy is assumed to reach the UHR region.

Assuming that n_e is more or less uniform in the chamber (according to the indications of Ref. 8), the UHR surface is located near $R_{\text{UH}} \approx 107 \text{ cm}$ in ISX and is roughly 2 cm ($\approx \Delta_{\text{UH}}$) from the location where $\tilde{E} \sim 3E_o$. We shall assume that anomalous wave absorption occurs in this region. The plasma volume in which the electrons are nonlinearly heated by the wave is then $(2b\Delta_{\text{UH}} \cdot 2\pi R_{\text{UH}}) \sim 7.1 \times 10^4 \text{ cm}^3$, giving $p_\mu \approx 1.6 \text{ W/cm}^3$. Since T_e is expected to be of the order of 100 eV, the dominant loss is due to particle drifts. We have from Eq. (13)

$$5.18(30 + \frac{3}{2} T_e)(86.4 \sqrt{T_e} + 15.5 T_e) \approx 1.6 \times 10^7 \left(\frac{\text{erg}}{\text{sec} \cdot \text{cm}^3} \right) \quad (16)$$

where T_e is in electron volts. This gives an estimate of $T_e \gtrsim 250$ eV near the UHR zone.

For a D-shaped reactor (e.g., TNS) with $B = 43$ kG, $R_0 = 500$ cm, $a = 125$ cm, $b = 200$ cm, and $\omega_0 = 7.56 \times 10^{11} \text{ sec}^{-1}$ (120 GHz), we assume a typical δB_p of 2 G. With $n_e \approx 10^{13} \text{ cm}^{-3}$ (at 1.5×10^{-4} torr), one finds $R_{UH} \approx 515$ cm and $\Delta_{UH} \approx 2$ cm. Assuming $T_e \sim 250$ eV, we have $\tau_\delta \approx 8.0 \text{ msec} \ll \tau_D$. With 2% oxygen, the maximum $(p_{\text{rad}} + p_{\text{ion}})$ occurs at $T_e \sim 20$ eV and is of the order of 0.1 W/cm^3 . It is negligible when $T_e > 100$ eV. Energy loss to the ions can be exaggerated by assuming $T_i \approx 0$ eV, giving $Q_{ie} \sim 0.14 \text{ W/cm}^3$. Drift losses in Eq. (13) are estimated to be 0.08 W/cm^3 . The volume in which strong electron heating occurs is roughly $2.6 \times 10^6 \text{ cm}^3$, giving an estimate of 0.57 MW for the power required. To ensure anomalous wave absorption due to a sufficiently large $\tilde{E}$ near the UHR region, the total area of wave irradiation should be around 10^4 cm^2 , giving $\tilde{E}_0 \sim 220 \text{ V/cm}$.

It is seen that for large tokamaks with fixed $\tilde{E}_0$ and T_e (~ 250 eV), the electron energy loss is dominated by the parallel drift. According to Eq. (9) and the relation $\Delta_{UH} \propto R n_e B^{-2}$, the total wave power then scales like $(n_e / \tau_\delta) b R \Delta_{UH} \propto n_e^2 R^2 B^{-3}$. If one assumes that roughly $n_e \propto B \propto R$ in going to larger tokamaks, he finds that $P_\mu \propto R$. In the case of unassisted start-up in full bore, we have $V_\ell \propto R$ [1, 2] and the peak OH current roughly satisfies $I_{OH} \propto I_p \propto a^2 B R^{-1} \propto R B$ for fixed aspect ratio. The required peak OH power then scales like $P_{OH} \propto R^2 B$ or R^3 if $B \propto R$ is assumed. It is clear that the microwave power for start-up has a more favorable scaling with the tokamak size than the peak OH power for the unassisted start-up in full bore.

5. REQUIRED LOOP VOLTAGE FOR START-UP

We assume that a steady-state plasma is produced in the toroidal chamber by the microwave irradiation, heating the plasma to $T_e \approx 250$ eV in a relatively narrow region near the UHR surface (see Fig. 1). When a constant loop voltage, V_ℓ , is applied, the induced plasma current will be predominantly in the region of high T_e . The ensuing poloidal magnetic field then "smears out" the microwave power input to the electrons and reduces T_e dramatically except in a region of radius $a_o \geq \Delta_{UH}$ near the magnetic axis. In the following, we shall make the pessimistic assumption that T_e stays near 250 eV until the resistive OH power to the plasma exceeds the microwave power. Also, the initial current channel is assumed to have a minor radius of a_o .

We further assume that the plasma current density stays uniform during start-up with the plasma minor radius, a , satisfying

$$a^2 = \frac{\mu_o R q}{2\pi B} I = a_o^2 \left(\frac{I}{I_o} \right) \quad (17)$$

where q is the safety factor, I is the plasma current (which equals I_o when $a = a_o$), and mks units are used for convenience. The plasma inductance can be written as

$$L = L_o - \frac{1}{2} \mu_o R \ln \left(\frac{I}{I_o} \right), \quad L_o = \mu_o R \left(\ln \frac{8R}{a_o} - 1.75 \right) \quad (18)$$

With Eqs (17) and (18), one can show that the behavior of I satisfies

$$\gamma = (\alpha - \frac{1}{2} \ln y) \frac{dy}{dt},$$

$$y = \frac{I}{I_0}, \quad \gamma = \frac{V_\ell - \mathcal{R}_0 I_0}{I_0 \mu_0 R}, \quad \alpha = \frac{L_0}{\mu_0 R} - 0.5, \quad (19)$$

$$\mathcal{R}_0 = \frac{2R}{a_0^2} \eta, \quad \eta(\Omega\text{-cm}) \approx 5.19 \times 10^{-3} \frac{Z_{\text{eff}}^{\lambda} \lambda_{\text{ei}}}{T^{3/2}(\text{eV})}$$

Here, $V_\ell = V_{\text{ind}} + V_{\text{res}}$ and $V_{\text{res}} = \mathcal{R}I = \mathcal{R}_0 I_0$ are constants in time. The solution of Eq. (19) is then

$$\gamma(t - t_0) = (\alpha + \frac{1}{2})(y - 1) - \frac{1}{2} y \ln y \quad (20)$$

where t_0 equals $(-L_0/\mathcal{R}_0) \ln[1 - (V_{\text{res}}/V_\ell)]$ and is the time required to increase I to I_0 when a is fixed at a_0 . For large values of (R/a_0) , t_0 is negligible. The dependence of y on $\gamma(t - t_0)$ is plotted in Fig. 2 for the case where $\lambda_{\text{ei}} \approx 15$, $Z_{\text{eff}} \approx 2$, and $(R/a_0) \approx 53.5$ and 28.3 .

There is considerable uncertainty in choosing a proper a_0 . When a_0 is increased by a factor of k for a fixed V_ℓ , Eqs (17) and (19) show that I_0 is multiplied by k^2 , γ is multiplied by k^{-2} , and $I_0 \mathcal{R}_0$ is unchanged. Figure 2 (points A and B) then shows that the times required for start-up to the same final I with $(R/a_0) \approx 53.5$ and 28.3 are nearly identical. This shows that the start-up process considered here is insensitive to variations in a_0 .

For ISX with $V_\ell = 40$ V, $a_0 = 2$ cm, $R = R_{\text{LM}} \approx 107$ cm, $B \approx 8.60$ kG, $I_0 \approx 0.80$ kA, $\mathcal{R}_0 I_0 = 1.7$ V, and $q = 2$, we have $\gamma = 3.56 \times 10^4 \text{ sec}^{-1}$ and $\alpha = 3.90$. Figure 2 (point A) shows that I reaches 200 kA when $(t - t_0) \approx 11.1$ msec. The corresponding resistive loss in volt-seconds is less than 0.02 Wb. If a resistive loss in volt-seconds of 0.1 Wb is permitted, a

start-up time of 59 msec is acceptable. Then γ should be $6.68 \times 10^3 \text{ sec}^{-1}$ and the required V_ℓ becomes only 8.9 V, which is a factor of nearly five below the loop voltage designed for ISX.

For TNS with $a_o = 9 \text{ cm}$, $R = R_{UH} = 515 \text{ cm}$, $B_T = 43 \text{ kG}$, $I_o = 16.9 \text{ kA}$, $R_o I_o = 8.47 \text{ V}$, and $q = 2$, we have $\gamma = 1.03 \times 10^3 \text{ sec}^{-1}$ and $\alpha = 3.86$.

Figure 2 (point A) then shows that $I = 4.2 \text{ MA}$ when $(t - t_o) = 0.38 \text{ sec}$ with $V_\ell = 122 \text{ V}$. The corresponding resistive loss in volt-seconds is 3.2 Wb. If a resistive loss of 15 Wb is permitted, the start-up time can be 1.8 sec. We then have $\gamma = 2.21 \times 10^2 \text{ sec}^{-1}$ and $V_\ell = 33 \text{ V}$. This is an order of magnitude less than the 300-500 V estimated for full-bore start-ups in reactors [1].

Finally, since V_{res} is now much less than V_{ind} , doubling Z_{eff} to 4 in Eq. (19) will double the resistive loss in volt-seconds during start-up while increasing V_ℓ by only a small amount.

6. SUMMARY AND DISCUSSION

We have demonstrated that microwave irradiation in the extraordinary mode at modest power, p_μ , from the high field side of currentless toroidal plasmas near the electron cyclotron resonance can substantially reduce the loop voltage, V_ℓ , required for start-up. It is shown that, in this fashion, V_ℓ can be reduced by a factor of roughly five in ISX at $p_\mu \sim 120 \text{ kW}$ and by a factor of roughly ten in TNS at $p_\mu \sim 0.57 \text{ MW}$. This will increase the duration for start-up by a factor of two to five from the unassisted case without increasing the resistive losses in volt-seconds. Upcoming start-up experiments in ISX will afford an opportunity to verify the estimates obtained here. If they are verified by the experiments, significant savings

in the power supplies for the OH and EF coils could be realized in future large tokamak reactors by this method. With the available high V_ℓ in existing tokamaks, this scheme could substantially increase the pulse length by minimizing the resistive loss in volt-seconds during start-up.

Several assumptions are used in making these estimates. They include the maintenance of a high density plasma ($n_e \sim 10^{12}-10^{13} \text{ cm}^{-3}$) to avoid producing many high energy runaway electrons at the electron cyclotron resonance, the parametric decay of the input wave into slightly oblique electron Bernstein waves and oblique ion-acoustic waves, the strong damping of the electron Bernstein waves to preferentially heat the electrons in a narrow region near the upper hybrid resonance, and a relatively low stray poloidal field ($\partial B_p / B$) $\sim 0.5 \times 10^{-4}$. We have assumed an electron particle and energy balance model based on electron impact ionization of the neutrals, losses by curvature and parallel drifts, collisional cooling by ions, and radiation by low Z impurities. To retain some pessimism in the estimates, the ions are assumed to be cold. It is further assumed that this localized wave heating process of the electrons promotes a small-radius channel of plasma current during the initial phase of start-up. As a result, this scheme could be an effective method of small-radius start-up.

Although our estimates of particle and energy balance tend to be pessimistic whenever large uncertainties are involved, we have not considered two very important questions in proposing this auxiliary start-up scheme: the effects of ambipolar electric field on the plasma equilibrium and the effects of macroscopic instabilities on the plasma confinement. While these effects may have strong bearings on the proposed start-up scheme, it

is appropriate to offer some discussions here and leave more detailed studies for the future.

6.1 AMBIPOLAR ELECTRIC FIELD AND PLASMA EQUILIBRIUM [23]

Since the ion temperature is much lower than the electron temperature, the parallel diffusion coefficient for the electrons, D_e , is more than a factor of $\sqrt{m_i/m_e}$ larger than D_i . The ambipolar electric field, E_{amb} , will then be such that

$$\Gamma_e = -D_e (\nabla n_e + \frac{en_e}{T_e} E_{amb}) \approx 0 \quad (21)$$

along the field lines. This implies an ambipolar potential, $V_{amb} \sim T_e/e$, which is sufficient to electrostatically confine the bulk of the electrons against losses due to parallel drift.

As a result, the curvature drift, v_D , and the $\underline{E} \times \underline{B}$ drift, v_E , become dominant. Since the electrons and the ions have opposing v_D , it is natural to ask whether the curvature drift would result in a vertical electric field, E_z , which propels the plasma outward in the direction of the major radius. The following demonstrates that the magnitude of E_z can be made negligibly small by using a metallic vacuum vessel with moderate conductivity.

In ISX, the top-to-bottom resistance of the vacuum vessel is of the order of 1 m Ω . At $T_e = 250$ eV and $n_e = 3.2 \times 10^{12}$ cm $^{-3}$, the current density, $J_z (= n_e e v_D)$, due to the electron curvature drift is roughly 0.028 A/cm 2 . The total current can be overestimated by assuming that T_e and n_e are uniform in the plasma. One obtains a total vertical current of less than 900 A in the ISX plasma. The corresponding voltage in the vertical direction

can be no larger than 1 V, giving $E_z \sim 0.02$ V/cm. The corresponding $\underline{E} \times \underline{B}$ drift in the direction of the major radius is only 2 m/sec, which is much smaller than v_D . Therefore, the ambipolar electric field is directed away from the plasma center.

The debye length for an unmagnetized plasma with $T_e = 250$ eV and $n_e = 3.2 \times 10^{12}$ cm $^{-3}$ is less than 10^{-2} cm. However, because of the toroidal magnetic field, the depth of the sheath region can be much larger than the debye length, with the electric field extending into the plasma [24]. The resulting ambipolar electric field, E_{amb} , is then roughly $(V_{amb}/a) \sim (T_e/ea)$. The corresponding $\underline{E} \times \underline{B}$ drift is roughly (cT_e/eBa) , which is larger than v_D as long as $(R/a) > 2$. The $\underline{E} \times \underline{B}$ drift introduces an effective rotational transform of the particle drift motion, which should result in an electron confinement time much larger than τ_D , barring losses enhanced by macroscopic instabilities. As a result, the spatial profiles of T_e and n_e are likely to be much broader than Δ_{UH} , the region of anomalous wave absorption. As discussed in Section 5, this increased size of the initial plasma is not expected to substantially alter the estimates of the required loop voltage and volt-seconds.

In addition, the ambipolar potential tends to electrostatically screen impurity ions away from the plasma [9]. The microwave-assisted start-up could then serve as an effective method of impurity control during start-up. The effects of the ambipolar electric field on the plasma equilibrium are expected to be beneficial in general. Our estimate of T_e is thus likely to be pessimistic.

6.2 MACROSCOPIC INSTABILITIES

In the experiments of Anisimov et al., plasmas in the high density regime [as defined by Eq. (1)] were apparently stable for a duration of roughly 30 μ sec. It is of interest to discuss briefly typical macroscopic instabilities and their possible effects on the plasma.

The most dangerous instabilities are expected to be the interchanges [25, 26]. For closed field lines and negligible electric field, the stability condition simplifies to

$$W\delta p\delta W + \gamma p(\delta W)^2 > 0 \quad (22)$$

where p is the plasma pressure, $W = -\oint dl/B = -(\pi cR^2/I_w)$, and I_w is the total current in the toroidal field coils. Equation (22) can be rewritten as

$$-\frac{R}{p} \frac{dp}{dR} \sim \frac{R}{L_p} < 2\gamma = \frac{10}{3}, \quad (23)$$

where p is the plasma pressure and L_p is the pressure gradient scale length in the direction of the major radius. Thus, for a value of (R/L_p) less than 3.3, the plasma can be internally stable against interchanges. Equation (23) shows that the equilibrium profiles of T_e and n_e may have a strong bearing on the stability of the plasma. It can also be argued that the interchange instabilities tend to nonlinearly broaden the spatial profile of p until Eq. (23) is nearly satisfied. However, Eq. (23) will be violated near the plasma edge where p approaches zero. In this case, a small loop voltage or some helical stellarator field may be applied in conjunction with the microwave irradiation to provide an average magnetic well and help stabilize the interchanges.

In the event that the interchanges can be marginally stabilized by proper pressure profiles, drift instabilities driven by ∇T_e and ∇n_e may become important. These instabilities, at their worst, could introduce particle and energy losses according to the Bohm diffusion coefficient

$$\frac{a^2}{\tau_B} \sim D_B \approx \frac{cT_e}{16eB} \quad (24)$$

With $T_e \approx 250$ eV, τ_B equals 4 msec in ISX ($a \approx 25$ cm) and 0.1 sec in TNE ($a \approx 125$ cm), which are larger than or comparable to τ_D according to Eq. (18). It is seen that the enhanced diffusion by drift instabilities may limit T_e to low values only away from the region where the anomalous absorption of the microwave occurs.

These discussions indicate the need of more detailed studies on the interactions between a strong extraordinary wave and an electrostatically confined axisymmetric currentless plasma, the equilibrium profiles and stability properties of such a wave, and its impact on the start-up process in tokamaks. If our estimates of the microwave-assisted start-up are verified by future experiments, they would contribute to promoting the development of high power microwave supplies around 120 GHz.

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APPENDIX: SUMMARY OF WORKSHOP DISCUSSIONS ON PLASMA START-UP

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July 19-20, 1977

Los Alamos Scientific Laboratory,

Los Alamos, New Mexico

1. CURRENT STATUS OF START-UP CODES

- Zero-dimensional (0-D) codes are adequate now for estimating the volt-second and plasma loop voltage requirements for full-bore, low density start-up.
- Plasma initiation has been modeled at the low Z atomic physics levels, which are important when the plasma temperature is below tens of electron volts. The final plasma state from a successful start-up turns out not to depend sensitively on the detailed plasma state during initiation. However, relatively small levels of low Z impurities (e.g., oxygen) have a strong bearing on the loop voltage required.
- One-dimensional (1-D) magnetic flux diffusion (i.e., skin current effects) can be modeled by nested plasma annuli with assumed temperature profiles. Based on the 1-D results, skin effects can be properly modeled by the 0-D codes. The inclusion of a simplified 1-D skin current penetration model is useful in comparing the code predictions with experiment and is valuable in checking the internal consistency of the calculation.
- MHD instability effects have not been included in the existing start-up codes and may be necessary in the future.

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- For TNS and other large machine design work, a more significant uncertainty than the validity of the simulation codes is the value of the internal plasma inductance, since this is important in determining the volt-second requirements. Experimentally, the internal inductance is reasonably well known for present ohmically heated discharges or those with moderate neutral beam heating. For machines with very large supplementary heating, current profile modification should be possible and the usual estimates based on present experience may not be completely adequate. In the simulation codes, the internal inductance of the plasma is essentially an input variable. For instance, in the 0-D codes, it is determined by the choice of $l/2$. In the 0-D codes with current penetration, it is determined by the form of the temperature profile, such as the experimental profile. In the fully 1-D codes, it is determined by the choice of a heat conduction model as well as a model for impurity transport and radiation. Each approach has its advantages and can be used to evaluate present experiments. Because of limited experimental data, none of these models is sufficiently developed to forecast future machine requirements other than approximately. In general, though, these approximate estimates are sufficient to determine volt-second requirements for TNS at the level of present needs.
- Regarding F. Marcus' atomic physics addition for hydrogen molecular dissociation, it appears that a small admixture of impurities such as oxygen tends to be more significant in the energy balance and thus in determining the voltage requirements for a normal discharge. In evaluating the evolution of partially ionized discharges used in discharge cleaning, Marcus' contribution may be very important, nonetheless.

2. CURRENT UNDERSTANDING BASED ON MODELING OF START-UP IN PLT AND ORMAK

- The start-up code has been successful in evaluating the plasma circuit interaction in PLT and in providing some insight into the evolution of the discharge. In certain discharges studied last year, the presence of low Z impurities, i.e., oxygen and carbon, appears to be very important in the gross energy balance, tending to depress the electron temperature. As a result of the low electron temperature and somewhat enhanced Z_{eff} due to the impurities, the skin current penetration time was short and tended to agree with a model of current penetration using neoclassical resistivity. In more recent discharges after extensive discharge cleaning, the low Z impurities are less important and do not appear to be the dominant loss mechanism. In these discharges, the current penetration time appears to be somewhat enhanced. In general, the codes have been sufficiently evaluated on PLT that they can be used with considerable confidence in evaluating the plasma-circuit evolution.
- As indicated by ORMAK and PLT results, the loop voltage required for full-bore start-up scales linearly with plasma major radius. Nonetheless, further experiments are necessary in evaluating the impurity influx and MHD activity during start-up and their influence on the development of the discharge.
- Gas puffing in PLT and gas puffing and neutral beam injection in ORMAK have indicated a high rate of rise in the central plasma density. Warc pinch effects are currently advanced as a possible explanation.
- The maximum achievable densities in various tokamak experiments have been found to scale linearly with B_T/T (the Murakami scaling). This indicates a nearly linear dependence of the achievable density on input power density.

3. PLASMA SHUTDOWN

- In most tokamak experiments, discharges are terminated by major disruption or by reversing the plasma loop voltage. Detailed experimental and modeling studies of this phase are only beginning.
- Major plasma disruption should be considered in the design studies of the poloidal field power supplies.

4. PLANNED AND PROJECTED START-UP EXPERIMENTS IN PLT, ISX, AND PDX

- Small-radius start-up off the limiter will be pursued in PLT.
- Impurity evolution, plasma wall interaction, and wall conditions will be diagnosed in ISX starting with its beginning discharges.
- Discharge cleaning at $B_T = 800$ G, assisted by microwave excitation at 2.45 GHz and a few kilowatts, will be conducted in ISX.
- The effects of microwave preionization and mild preheating on the start-up loop voltage will be studied in ISX.
- Impurity control with small-radius start-up off a poloidal field divertor will be conducted in the Poloidal Divertor Experiment (PDX).

5. IMPLICATIONS ON PULSED POWER SUPPLY REQUIREMENTS OF THE TNS POLOIDAL FIELD SYSTEMS

- In the case of full-bore start-up, TNS designs of the voltage peak and pulse length should retain some flexibility. Present experimental results are not sufficient to determine an optimal voltage pulse wave form.
- Although research and development (R&D) of pulsed energy storage and power supplies for TNS are definitely needed now, one should note that current understanding of the start-up physics is limited. Planned experiments in the immediate future are expected to decide with better precision the pulsed loop voltage and power supply required. Innovative start-up schemes such as small-radius start-up, microwave-assisted start-up, and ion-core-assisted start-up may result in a more efficient use of voltage and power supplies. It is advisable to retain flexibility in the R&D program in anticipation of these possibilities and the associated savings in the cost of the TNS tokamak.

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FIGURE CAPTIONS

Fig. 1. The assumed mechanisms of microwave (as extraordinary mode) absorption near the ECR and the UHR in a currentless tokamak plasma.

(a) In the low density regime where Eq. (1) is not satisfied, relatively strong absorption at the ECR region (Δ_{EC}) occurs, producing electrons above a few kiloelectron volts which are weakly coupled to the plasma. The remaining reduced wave is then assumed to undergo linear conversion into electron Bernstein modes of small amplitude. The resulting plasma electron temperature is then relatively low (~ 10 eV). (b) In the high density regime when Eq. (1) is satisfied, only weak interaction occurs at the ECR region. The almost undamped wave then linearly accumulates to large amplitude near the UHR and parametrically decays into electron Bernstein and ion-acoustic modes. The Bernstein mode is assumed to be strongly damped to heat the electrons to a few hundred electron volts in a relatively narrow region (Δ_{UH}).

Fig. 2. The dependence of y ($= I/I_0$) on $\gamma(t - t_0)$ according to Eq. (20) for two different values of the initial plasma aspect ratio (R/a_0).

(a) LOW DENSITY

(b) HIGH DENSITY

Fig. 1.

Fig. 2.

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68. CTR Reading Room, c/o Prof. Allan N. Kaufman, Physics Dept., Univ. of California, Berkeley, CA 94720
69. Dr. David A. Dingee, Fusion Programs, Battelle-Northwest, Battelle Boulevard, Richland, WA 99352
70. CTR Reading Room, c/o Prof. C. S. Liu, Dept. of Physics and Astronomy, Univ. of Maryland, College Park, MD 20742
71. CTR Reading Room, c/o Prof. T. Kammash, 103 Research Admin. Bldg., N. Campus, Univ. of Michigan, Ann Arbor, MI 48105
72. CTR Reading Room, c/o Dr. Ravi N. Sudan, Phillips Hall, Cornell Univ., Ithaca, NY 14850
73. Prof. Marshall N. Rosenbluth, Institute for Advanced Study, Princeton, NJ 08540
74. CTR Reading Room, c/o Prof. R. Gross, Plasma Research Lab., Columbia Univ., New York, NY 10027
75. CTR Reading Room, c/o Prof. Roy Gould, California Inst. of Tech., M.S. 116-81, Pasadena, CA 91125
76. Dr. Nicholas A. Krall, Science Applications, Inc., P.O. Box 2354, 1200 Prospect St., La Jolla, CA 92037
77. CTR Reading Room, c/o Dr. Jay P. Boris, Plasma Physics, Naval Research Laboratory, Washington, D.C. 20390
78. Professor A. Simon, Dept. of Mechanical & Aerospace Sciences, University of Rochester, Rochester, NY 14627
79. CTR Library, c/o Dr. Alan F. Haight, United Technologies Research Labs, East Hartford, CT 06108
80. Dr. H. K. Forsen, Exxon Nuclear Co., Inc., 777-106th Avenue, NE, C-000777, Bellevue, WA 98009
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89. Dr. Gareth E. Guest, General Atomic Co., P.O. Box 81608, San
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90. Dr. Claude Mercier, Service du Theorie des Plasmas, Centre
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