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ON THE IONIZATION AND BURNOUT PROCESSES OF A MAGNETICALLY CONFINED PLASMA

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ABSTRACT

The particle and heat balance during plasma start-up are investigated, to specify the conditions for reaching various ion density ranges and high plasma temperatures in cases of a limited heating power. Particular attention is paid to the permeable-impermeable transition regime of plasmas being subject to Ohmic heating and confined in closed or open bottles with a main poloidal field. The ionization and burnout conditions are found to depend critically on the confinement and the filling density. They become optimal in closed bottles under symmetric and stable conditions, where the transition into a fully ionized state should be reached even at moderately large ionization rates, burnout powers, and currents. Start-up methods based on constant as well as on variable filling densities are discussed as means of ion density control.

1. Introduction

The balance of magnetically confined plasmas including start-up, refuelling, purification, achievable temperatures by imposed heating mechanisms, and stability, often depends in a crucial way on the range of ion density being chosen. Thus, there is a need for experiments in which the ion density can be varied within wide limits, also including the transition regime between plasmas being permeable and impermeable to neutral gas. Control of the ion density also becomes important in attempts to achieve optimal conditions for plasma confinement, purity, heating, and stability.

Among the possible ways of creating a fully ionized plasma of desired ion density, we shall in this paper treat some methods by which a given amount of neutral gas atoms is converted into a plasma within a fixed container volume. The present treatment is mainly concentrated to the transition regime between permeable and impermeable plasmas and to confinement in poloidal magnetic fields, but could also be modified such as to apply to other density regimes and field geometries. There are especially two basic phenomena which will be considered. The first concerns the particle balance during the ionization process, with its associated conditions for creating a fully ionized plasma. The second concerns the energy and heat balance during the burnout process, as well as the power input which is necessary for making the transition from a lowly to a highly ionized state. In many cases the power required to sustain an ion density n in presence at an immersed neutral gas of density n_n becomes proportional to the product nn_n. The latter has low values at low and high ionization degrees, with a maximum half-way between which corresponds to the required burnout power threshold. An example of this behaviour will be given in the present paper for an ohmically heated plasma.

2. Basic Conditions

A magnetically confined plasma is considered under the following simplified conditions and assumptions:

(i) The plasma and the neutral gas are contained inside a closed vessel of volume V_0 having ideally reflecting walls. The volume of the magnetic confinement region is $V < V_0$. Thus, the total number of atomic nuclei is constant within V_0.

(ii) The plasma is confined by a strong poloidal magnetic field $\underline{B}$.

(iii) The whole system is axially symmetric. Effects due to deviations from axial symmetry are discussed later in this context.

(iv) The velocity distributions of ions and electrons are isotropic. In closed magnetic bottles this assumption holds as long as the effective collision times of charged particles are somewhat shorter than the particle containment time due to diffusion across $\underline{B}$. In open bottles the same assumption leads to stronger restrictions and holds only when the effective collision times of the particles become smaller than the times of flight between the bottle ends.

(v) The ion and electron temperatures are nearly equal.

(vi) Volume recombination and impurity effects are neglected.

(vii) The immersed neutral gas is assumed to be nearly at rest.

(viii) The macroscopic plasma velocity across $\underline{B}$ is small.

(ix) The only energy source is by Ohmic heating. The latter is provided either by induced low-frequency currents along the poloidal magnetic field, or by toroidal currents induced at frequencies being much higher than the relaxation frequencies of the ionization and burnout processes to be considered here. In both cases the low-frequency part of the toroidal electric field can be neglected in the first approximation. In the particular situation of high-frequency heating we also neglect the corresponding fluctuations in the particle densities and transport coefficients when considering the much slower transient ionization and burnout processes.

(x) The gyro frequencies of ions and electrons are much higher than the corresponding collision frequencies and characteristic macroscopic frequencies.

(xi) According to earlier discussions on laboratory plasmas with low impurity content, bremsstrahlung losses are neglected as compared to those from heat conduction across the magnetic field [1].

(xii) Rotating plasmas are excluded from the main context. In particular, for the partially ionized regions of such plasmas as well as in some other systems with fluid motions, there is a general limitation by Alfvèn's critical velocity $v_c=(2e\phi_i/m_i)^{1/2}$ where ϕ_i is the ionization potential. We shall only consider this limitation in connection with a short discussion on rotating plasmas in Section 5.

2.1. Basic Equations

With conditions and assumptions (i) - (xii) the plasma balance will now be approximated by the following set of earlier defined equations given in SI units [1-3], i.e.,

$$\frac{\partial n}{\partial t} + \mathrm{div}(n\underline{v}) = nn_n\xi \qquad (1)$$

$$nm \frac{d\underline{v}}{dt} = \underline{j} \times \underline{B} - \underline{\nabla} p - nm\alpha\underline{v} \tag{2}$$

$$\eta\underline{j} = \underline{E} + \underline{v} \times \underline{B} - (1/en)\underline{j} \times \underline{B} + (1/2en)\underline{\nabla} p +$$

$$+ (3nk\eta_{ei}/2B^2)\underline{B} \times \underline{\nabla} T \tag{3}$$

$$\frac{3}{2}\frac{\partial p}{\partial t} + \frac{3}{2}\,\mathrm{div}(p\underline{v}) + p\,\mathrm{div}\underline{v} + \mathrm{div}\underline{q} =$$

$$= \eta\underline{j}^2 - nn_n e\phi_i\xi - \frac{3}{2}\,nn_n kT(f_{in}\rho_{in} + f_{en}\rho_{en}) \tag{4}$$

Here n, T, $p = 2nkT$, $\underline{v}$ are the ion density, the plasma temperature, pressure, and fluid velocity, n_n is the neutral gas density, $m = m_i + m_e$, $\underline{j}$ is the current density, $\underline{E}$ the electric field, ξ the ionization rate, $\alpha = n_n\rho_{in}$, $\eta = m_e(\nu_{ei} + \nu_{en})/e^2n \equiv \eta_{ei} + \eta_{en}$, $\eta_{ei} = k_\eta/T^{3/2}$, $k_\eta = 129(\ell n\Lambda)$, $\nu_{en} = n_n\rho_{en}$, $\rho_{in} = \langle\sigma_{in}w_{in}\rangle$ is the effective rate of ion-neutral impacts, $\rho_{en} = \langle\sigma_{en}w_{en}\rangle$ is the effective rate of electron-neutral impacts excluding ionization, ϕ_i is the ionization potential, f_{in} is the effective fraction of ion energy lost in an ion-neutral encounter [3] and f_{en} the corresponding fraction lost by an electron when excitation radiation is included but ionizing impacts neglected, $\underline{q}$ is the heat flow vector leading to a heat flow $\underline{q}_\perp = -\lambda^*\underline{\nabla}_\perp T$ across $\underline{B}$ where $\lambda^* = 5nk^2T\nu_{ii}/4m_i\omega_i^2$ in the strong-field limit with ν_{ii} standing for the ion-ion collision frequency and ω_i for the ion gyro frequency, and the rest of the symbols have their conventional meaning.

Under condition (ix) specified at the beginning of this section, the possibly existing high-frequency parts of eqs. (1) - (4) can be separated from those describing the transient low-frequency processes of ionization and burnout. This corresponds to letting the symbols of eqs. (1) - (4) represent only the low-frequency part of the problem henceforth, whereby the first term of the right hand member of eq. (4) has to be replaced by $\eta(\underline{j} + \underline{j}_0)^2 \simeq \eta j_0^2$. Here j_0 is the effective value of a strong externally imposed poloidal low-frequency or toroidal high-frequency current.

2.2. Classical Diffusion Across the Magnetic Surfaces

We now use subscripts $(_p)$ and $(_t)$ to denote the poloidal and and toroidal low-frequency parts of the present field quantities. The toroidal part of eq. (3) yields

$$- \underline{B}_p \times (\underline{v}_p \times \underline{B}_p) = \eta \underline{j}_t \times \underline{B}_p - (3nk\eta_{ei}/2B_p^2)\underline{B}_p \times (\underline{\nabla}T \times \underline{B}_p) \qquad (5)$$

after vector multiplication by $\underline{B}_p$. From combination with the poloidal part of eq. (2) and with the scalar product between $\underline{B}_p$ and eq. (2), the diffusion velocity $\underline{v}_{p\perp}$ across the magnetic surfaces becomes determined by

$$\left[1 + (nm\eta/B_p^2)(\alpha + \frac{\partial}{\partial t})\right] \underline{v}_{p\perp} \simeq$$

$$\simeq - (\eta_{ei}k/2n^3B_p^2)\underline{\nabla}_\perp(n^4T) - (\eta_{en}/B_p^2)\underline{\nabla}_\perp p \qquad (6)$$

Here it is easily seen that the contributions from α and $\partial/\partial t$ can be neglected on account of condition (x) .

2.3. Neutral Gas Penetration into the Confinement Region

For a plasma body of the characteristic dimension L_b a critical density

$$n_{cf} = 1/\sigma_{cf}L_b; \quad 1/\sigma_{cf} = \left[2kT/m\xi(2\xi + \rho_{in})\right]^{1/2} \tag{7}$$

can be defined where the plasma becomes permeable to fast neutral gas particles when its average ion density $\bar{n} \lesssim n_{cf}$, and becomes impermeable to neutral gas when $\bar{n} \gg n_{cf}$ [3]. A bar on top of a symbol is used henceforth to denote mean value over the plasma volume V . In the case of hydrogen and helium the quantity $1/\sigma_{cf}$ has a nearly constant value of about $3 \times 10^{18}\ m^{-2}$ at temperatures $T \gtrsim 10^5$ K, whereas the same quantity increases steeply as T decreases within the range below 10^5 K . The permeability conditions can , of course, become changed during a burnout process during which the ion density is increasing.

2.3.1. Permeable Plasmas

When $\bar{n} \lesssim n_{cf}$, fast neutral particles will fill the confinement volume almost uniformly [3]. Provided that $\bar{n}$ is not chosen too far below n_{cf}, there will also exist slow neutral particles within a narrow wall-near layer. The contribution of these latter particles to the total number of neutral atoms can then be shown to be negligible according to the theory of Ref. [3]. Under these conditions, and with assumptions (i) of the beginning of this section in mind, the filling density n_{no} of neutral atoms becomes related to the average densities $\bar{n}$ and $\bar{n}_n$ of ions and neutrals by

$$\bar{n}_n = n_{no} - (V/V_o)\bar{n} \tag{8}$$

At laboratory dimensions with $L_b \approx 0.1$ m and during the essential stages of the ionization and burnout processes, eq. (8) should hold with sufficiently good approximation at filling densities $n_{no} \lesssim 10^{20}$ m^{-3} when $T \gtrsim 10^5$ K, and even for higher values of n_{no} when $T < 10^5$ k. Since the possibilities of burnout in the transition region between permeable and impermeable plasmas as well as at lower ion densities are of special interest here, we shall assume the conditions leading to eq. (8) to be valid as a first approximation in the later Sections 3 and 4.

2.3.2. Impermeable Plasmas

When $\bar{n} \gg n_{cf}$ the neutral gas is only able to penetrate into a thin partially ionized wall-near layer of the plasma body. If this condition prevails during the essential stages of the ionization and burnout processes, eq. (8) does no longer hold, and the relations between n_{no}, n, and n_n become more complicated.

2.3.3. The Burnout Process

During the burnout process which converts a neutral gas cloud into a fully ionized plasma in a confinement region being smaller than the volume of the discharge chamber, the situation is sometimes less straight-forward than those described in the previous Sections 2.3.1 and 2.3.2. Thus, ion-neutral collisions with a hot and/or rapidly rotating plasma can produce an outflux of fast neutrals which leave the confinement volume, at the same time as an influx of neutrals enters the same volume from wall-near regions. This produces a depression in neutral gas density within the volume occupied by the plasma, in its turn affecting the burnout conditions [4]. Such density depressions will not be considered in the analysis of the coming Sections 3 and 4.

3. The Particle Balance

In a crude estimate of the particle balance we now integrate eq. (1) over the plasma volume and introduce meanvalues from which

$$\overline{\frac{\partial n}{\partial t}} \approx \overline{n\xi} \left\{ n_{no} - \left[\bar{n}(V/V_o) + (v_{p_\perp} S_\perp / V\bar{\xi}) + (v_{p_\parallel} S_\parallel / V\bar{\xi}) \right] \right\} \tag{9}$$

Here $v_{p_\perp}$ stands for the absolute value of the transverse diffusion velocity in eq. (6),

$$v_{p\parallel} = k_w (kT/m_i)^{1/2} \; ; \qquad k_w = 1 + 1/\sqrt{2\pi} \tag{10}$$

is the effective velocity of plasma escaping along $\underline{B}_p$ to the end walls in the case of an open magnetic bottle [1], and $S_\perp$ and $S_\parallel$ are the corresponding surface areas of the plasma body in the directions across and along $\underline{B}_p$. The value of k_w in expression (10) applies to the case of non-conducting walls and ambipolar plasma streaming, and should be replaced by $k_w \approx 1/\sqrt{2\pi}$ in the case of electrically conducting end walls. With the notation $|-\underline{\nabla}_\perp| \approx 1/L_\perp = S_\perp/V$ and $1/L_\parallel = S_\parallel/V$ as well as

$$\theta_{ei} = kk_\eta / 2\bar{\xi}\sqrt{\bar{T}}(B_p L_\perp)^2 \simeq m_e a_i^2 \nu_{ei} / 4 m_i L_\perp^2 \overline{n\xi} \tag{11}$$

$$\theta_{en} = 2 m_e \rho_{en} k\bar{T} / e^2 \bar{\xi} (B_p L_\perp)^2 \simeq m_e a_i^2 \nu_{en} / m_i L_\perp^2 \bar{n}_n \bar{\xi} \tag{12}$$

$$n_p = k_W(k\bar{T}/m_i)^{1/2}/\bar{\xi}L_{\#} \tag{13}$$

where a_i is the average ion Larmor radius, we obtain

$$\frac{\overline{\partial n}}{\partial t} \simeq \bar{n}\bar{\xi}\left[n_{no}(1-\theta_{en}) - \bar{n}\frac{V}{V_o}(1-\theta_{en} + \frac{V_o}{V}\theta_{ei}) - n_p\right] \tag{14}$$

From expressions (11) and (12) is seen that θ_{ei} and θ_{en} usually become much smaller than unity. In cases where the square bracket of eq. (14) becomes positive the ion density should thus increase up to a limiting value being a function of the corresponding plasma temperature.

Two special cases are now considered:

(i) When $S_{\#} \neq 0$ and there are plasma losses to the ends of an open bottle in presence of a very strong confining field, eq. (14) reduces to

$$\frac{\overline{\partial n}}{\partial t} \simeq \bar{n}\bar{\xi}\left[n_{no} - \bar{n}(V/V_o) - n_p\right] \tag{15}$$

In a discharge which starts at small ion densities $\bar{n}$, it is seen that $\bar{n}$ can only grow when the filling density n_{no} is chosen above the "Poletaev limit" n_p [5,2]. When this is the case, $\bar{n}$ increases up to the steady-state value

$$\bar{n}_\infty = (V_o/V)(n_{no} - n_p) \tag{16}$$

As a numerical illustration we choose hydrogen with $\bar{T} = 4 \times 10^4$ K and $L_{\parallel} = 0.3$ m from which $n_p \simeq 3 \times 10^{20}$ m^{-3}. Thus, the Poletaev limit occurs at rather high densities in open-ended systems with large end losses. Further, due to the saturation of the ionization rate ξ in the range $T \gtrsim 10^5$ K, this limit has a nearly fixed minimum value at high plasma temperatures. In the present example $n_p \simeq 2 \times 10^{19}$ m^{-3} when $10^5 < \bar{T} < 10^6$ K.

(ii) When $S_{\parallel} = 0$ and there is a closed bottle with good confinement corresponding to small values of θ_{ei} and θ_{en}, the density $\bar{n}$ will always increase at small $\bar{n}$, as seen from eq. (14) with $n_p = 0$. It approaches the steady-state value

$$\bar{n}_{\infty} = (V_o/V)n_{no}(1-\theta_{en}) / \left[1 - \theta_{en} + (V_o/V)\theta_{ei}\right] \simeq (V_o/V)n_{no} \qquad (17)$$

The corresponding neutral gas density is related to $\bar{n}_{\infty}$ through eq. (8), in agreement with an earlier analysis of the case $n_{en} << n_{ei}$ [3]. As a numerical illustration hydrogen is chosen with $B_p = 0.5$ tesla and $L_{\perp} = 0.1$ m yielding $\theta_{ei} \simeq 3 \times 10^{-18}/\xi\sqrt{\bar{T}}$ and $\theta_{en} \simeq 2 \times 10^{-28}\, \bar{T}^{3/2}/\xi$. With $\bar{T} = 4 \times 10^4$ K we then have $\theta_{ei} \simeq 5 \times 10^{-5}$ and $\theta_{en} \simeq 5 \times 10^{-6}$ which implies that the right hand member of eq. (17) becomes a valid approximation and the neutral gas density is very small in the final state. Consequently, when the plasma confinement is good, an efficient ionization can be achieved even at temperatures where the ionization rate ξ is far from its saturation (maximum) value.

4. The Heat Balance

We now turn to the heat balance equation (4) under the condition (ix) at the beginning of Section 2 which is further specified by the last paragraph of Section 2.1. With the results of Sections 2.2 and 3 in mind, the treatment can be simplified by the following approximations of the left hand member in eq. (4). First, the convective and conductive transverse heat flows $(3/2)\mathrm{div}(p\underline{v}_\perp)$ and $\mathrm{div}\underline{q}_\perp$ from Coulomb interaction have a ratio of the order of $1/\sqrt{A}$ where A is the mass number. Second, in the case of an open bottle, the end surfaces of area $S_\parallel$ are assumed to be non-conducting, from which the heat transport to these surfaces along $\underline{B}_p$ will be governed by free ambipolar streaming at the velocity v_{pw} of eq. (10). The heat transport from the plasma to the end walls along the magnetic field is then given by the flux of escaping ion-electron pairs, regardless of the fact that there is a high heat conductivity along $\underline{B}_p$ inside the plasma body, being due to the electrons. Most of the latter are namely turned back into the plasma by the ambipolar electric field near the surface of an insulated end wall [1]. Third, the effective value of the imposed heating current j_o is chosen much larger than the part of the total current density being associated with the slow transient process of the plasma balance.

With these approximations a crude meanvalue formation of eq. (4) across the plasma body leads to

$$\frac{3}{2}\frac{\overline{\partial p}}{\partial t} \simeq \overline{nj_o^2} - \bar{n}\bar{n}_n e\phi_i F - Cn^2 k\bar{T}\xi\theta_{ei} - 3k_w(k\bar{T})^{3/2}\bar{n}/\sqrt{m_i}L_\parallel \tag{18}$$

where C is a dimensionless constant of order unity,

$$F = [1 + 2C(k\bar{T}/e\phi_i)\theta_{en}]\bar{\xi} + (3k\bar{T}/2e\phi_i)(\overline{f_{in}\rho_{in}} + \overline{f_{en}\rho_{en}}) \tag{19}$$

and θ_{ei}, θ_{en}, k_w are defined by eqs. (11), (12), (10).

It should be observed that the last term of eq. (19) includes the losses due to excitation radiation produced by electron-neutral collisions. In situations of insufficient power input, the plasma temperature and the ionization rate may further remain at a sufficiently low level for the plasma to become penetrated by substantial amounts of neutral gas. This, in its turn, can lead to very large excitation radiation losses, such as those being observed in the form of a "Lyman-alpha-catastrophe" in certain experiments [6].

From eq. (18) is further seen that, for every value of $\bar{n}, \bar{n}_n, T$, the pressure p can be made to rise when sufficiently large heating currents j_0 are imposed. The latter will, in their turn, depend on the response to the imposed electric field $\underline{E}$, being governed by the boundary conditions and the plasma properties in a way not to be discussed in detail here. It should only be mentioned that rather high electric fields may be needed in certain cases of high-frequency heating by toroidal currents, as pointed out by Tennfors [7].

4.1. Large Longitudinal Losses

In the case where the last term of eq. (18) represents the most important loss, and where Ohmic heating by Coulomb collisions dominates, the average current density of a steady state becomes determined by

$$j_{0\infty}^2 \simeq 3k_W k^{3/2} \bar{T}^3 \bar{n} / k_\eta \sqrt{m_i} L_\parallel \qquad (20)$$

As a numerical illustration we choose hydrogen with $\bar{T} = 4 \times 10^4$ K, $L_\parallel = 0.3$ m, $\bar{n} = 10^{20}$ m^{-3} and obtain $j_{0\infty} \simeq 10^7$A/m^2. Thus, quite large currents are needed to sustain a plasma with the present type of end losses, even at rather low temperatures.

4.2. Classical Confinement in a Closed Bottle

We now turn to a closed bottle with classical magnetic confinement where θ_{ei} and θ_{en} are much smaller than unity.

4.2.1. The Burnout Process

During the burnout process, i.e. before the neutral density $\bar{n}_n$ has reached a low value compared to $\bar{n}$, the term including F will represent the main loss in eq. (18). Since $\eta_{ei} >> \eta_{en}$ at values of $\bar{n}$, $\bar{n}_n$, $\bar{T}$ being of interest during this process, eq. (18) then reduces to

$$\frac{3}{2}\frac{\overline{\partial p}}{\partial t} \simeq \bar{\eta}_{ei}(\overline{j_o^2} - \overline{n n_n} e\phi_i F\bar{T}^{3/2}/k_n) \tag{21}$$

At a given temperature $\bar{T}$, the transition to a highly ionized state can thus be achieved only when j_o exceeds a burnout value j_{ob} corresponding to the maximum of $\overline{n n_n}$. From eqs. (21) and (8) this value becomes

$$j_{ob} \simeq n_{no}(e\phi_i F\bar{T}^{3/2}V_o/4Vk_n)^{1/2} \tag{22}$$

The obtained result is, of course, not valid for temperatures being low enough for volume recombination and other effects to invalidate the main conditions and assumptions of Section 2.

The obtained burnout condition (22) for the current is related to the maximum of $n n_n$ and has some similarities to an earlier derived result for the momentum and power balance of rotating plasmas [2,4]. In the present case of Ohmic heating, however, it is more convenient to express the final burnout condition in terms of an applied electric field.

4.2.2. The Final State

As soon as the burnout current can be exceeded for all values of T which lead to a highly ionized state, it is seen from equations (18) and (14) that such a state also should be reached. Since there will still be traces of neutral gas left within the plasma, and since there remain losses due to the finite values of θ_{ei} and θ_{en}, a certain heating current is required to sustain the plasma also in its final state. From eqs. (18) and (14) this current becomes

$$j_{o\infty} \simeq n_{no}(e\phi_i F\bar{T}^{3/2}\theta_{ei}/k_n)^{1/2} = j_{ob}(2V\theta_{ei}/V_o)^{1/2} \qquad (23)$$

When $\theta_{ei} << 1$ we thus have $j_{o\infty} << j_{ob}$.

4.2.3. Numerical Examples

The previous deductions are illustrated by an example with hydrogen where $V_o = 1.5$ V, $f_{in} \simeq 0.5$, $B = 0.5$ tesla, and $L_\perp = 0.1$ m. We further choose an average temperature $\bar{T} = 3 \times 10^4$ K at which the build-up time is of the order of $1/n_o\xi < 4 \times 10^{-4}$ s for $n_{no} \geq 3 \times 10^{19}$ m^{-3}. This corresponds to an acceptable ion density growth during the main ionization phase of an experiment. From eqs. (11) and (12) we then have $\theta_{ei} \simeq 2 \times 10^{-3}$ and $\theta_{en} \simeq 2 \times 10^{-6}$. These data lead to the following results:

(i) According to eq. (22) the burnout current density becomes $j_{ob} \simeq 1.7 \times 10^{-15} n_{no}$ A/m^2. With $n_{no} = 3 \times 10^{19}$ m^{-3} this leads to the total currents $j_{ob}S_p \simeq 5000$ A and $j_{ob}S_t \simeq 500$ A in the cases of heating by poloidal currents flowing through the cross section $S_p = 0.1$ m^2 and by toroidal currents through the cross section $S_t = 0.01$ m^2, respectively.

(ii) According to eq. (23) the steady-state current density is determined by $j_{0\infty}/j_{ob} \simeq 5 \times 10^{-2}$ with the data of (i) above. Within the limits of the present theory on a stable permeable plasma in a closed bottle, we then arrive at steady heating currents being as small as $j_{0\infty}S_p \simeq 250$ A and $j_{0\infty}S_t = 25$ A. These values will, of course, be raised considerably in the case of asymmetries and instabilities, and they also become modified when the simple balance of Section 3 for a permeable plasma does no longer hold true.

5. Some Comments on Rotating Plasma Operation

The ionization and burnout processes in rotating plasmas are in some respects more complicated than those described in this paper. Here the following brief comments will be given on this matter:

5.1. Particle Balance Problems

(i) Axially symmetric, isotropic plasmas which are confined in open magnetic bottles should in principle behave like the system described in Section 3. (i), but should have a lower Poletaev limit when the centrifugal confinement reduces the plasma density at the end insulators.

(ii) At sufficiently low densities and collision frequencies anisotropic velocity distributions can be established in open bottles of magnetic mirror type. Under axially symmetric conditions this could lower the plasma losses and the Poletaev limit considerably.

(iii) In closed bottles of "Tornado" type, the Poletaev limit should in principle be removed, in analogy with the case of Section 3(ii), provided that the effects of asymmetries can be neglected [8].

(iv) In presence of asymmetries, such as the spoke-shaped structures observed in rotating plasma experiments during the ionization phase [2], the plasma losses become strongly enhanced by rapid radial drifts due to asymmetric electric fields, at the same time as the ionization rate ξ becomes non-uniformly distributed across the spoke structure and the confinement region [2]. This is expected to increase the Poletaev limit, probably as far as to levels being comparable to those given by eq. (13). Possibly the inhomogeneous spoke structure and its associated critical velocity phenomenon may even become a necessary mechanism for reaching sufficiently high ionization rates within limited parts of the plasma volume [2]. In this connection the questions arise whether the spoke structure can be avoided at low plasma densities by introducing auxiliary discharges in the anode region [9], and/or by introducing a hot cathode to avoid cold-cathode spots.

5.2. Heat and Momentum Balance Problems

(i) The crossed electric and magnetic fields provide an efficient shear heating mechanism, both during the ionization process and at later stages.

(ii) The heat losses of open bottles are reduced by the centrifugal force, but are enhanced by spoke-shaped asymmetries.

(iii) During the burnout process a limit is set by the drag force between the rotating plasma wind and the neutral gas. Under certain conditions this also leads to a decrease in neutral density within the central parts of the plasma, in a way which modifies the burnout conditions as compared to those treated in this paper [2].

6. Conclusions and Discussions

Both with the purpose of investigating the ion density ranges of permeable and impermeable plasmas and of reaching higher plasma temperatures, experimental conditions have to be found for controlling the plasma density. In principle there are two ways, i.e. by methods of creating fully ionized plasmas at a constant filling density, and by methods of variable filling density in the form of matter injection or extraction.

6.1. Constant Filling Density

In the present paper some investigations have been presented on the case of constant filling density, in a range where the plasma is mainly permeable to neutral gas. The results, which in the first place apply to a poloidal confining magnetic field, are summarized and discussed in the following way:

(i) In closed magnetic bottles it should become possible to make the transition to a fully ionized state even by means of moderately large burnout currents, powers, and ionization rates being much smaller than the saturation value $\xi \simeq 10^{-14}$ m^3/s, provided that a good axial symmetry of the plasma can be preserved. In the final state of a holding mode, the steady heating current and its corresponding power input should then become much smaller than that required for burnout.

(ii) In open magnetic bottles and under symmetric plasma conditions there exists a "Poletaev limit" of the filling density which has to be exceeded for the ionization degree to be able to increase from low to high levels. The required ionization rates and burnout currents are larger in open than in closed bottles, at least when the latter are based on a main poloidal field.

(iii) Plasma asymmetries and instabilities produce enhanced losses and should lead to substantial increases in the required ionization rates, burnout currents, and power levels. Transverse drift motions due to asymmetries may introduce a high Poletaev limit, also in closed bottles.

(iv) Poloidal low-frequency current systems provide a simple and efficient heating mechanism at moderately high temperatures, being based on a poloidal electric field and the density-independent part η_{ei} of the resistivity. A disadvantage is due to possibly arising azimuthal inhomogeneities from a thermal instability by which the heating current becomes channelled to certain sectors of the perimeter [10-12]. In a fully ionized plasma this instability can be suppressed within certain parameter ranges, by such effects as transverse heat conduction and additional plasma heating of cool regions [12], as well as by convection due to charged particle drifts in an inhomogeneous magnetic field or in a weak superimposed transverse electric field, and possibly also by anomalous transport effects. Further, in the range of ionization degrees below the burnout level, there are stabilizing effects from nn_n as being roughly demonstrated by the bracket in the right hand member of eq. (21). In the case of an asymmetric local increase in T within a certain plasma sector, the ionization degree and nn_n should namely tend to increase there. Since also $F/T^{3/2}$ increases steeply with T in the range $T < 10^5 K$, this would in its turn require an increased local electric field $E_0 \simeq \eta j_0$, as compared to that prevailing in adjacent cooler plasma sectors. On the other hand E_0 is a given axially symmetric poloidal field, being imposed by external means. Consequently, the current j_0 will instead decrease within the hot plasma sector, and this tends to suppress the temperature asymmetry. Above the burnout level the variation of nn_n with increasing n is on the other hand destabilizing, but in this range combination with additional heating methods may become possible, as expressed later under point (vii).

(v) Toroidal high-frequency current systems have the advantage of good axial symmetry. The main problem with these systems is the poor current response to an imposed toroidal high-frequency electric field, this response being impeded by the polarization effects due to the equivalent electric capacity of the plasma [5].

(vi) Rotating plasma systems have the advantage of a large available power input. On the other hand, there are disadvantages due to an existing Poletaev limit, longitudinal plasma losses in open bottles, difficulties in support shielding against the rotating plasma wind in internal ring systems, and asymmetries spoke formation during the ionization and burnout processes.

(vii) To achieve controllable, symmetric, and ideal conditions during the ionization and burnout processes, it may become necessary to combine in a proper way the three systems under (iv)-(vi) above, and possibly also to extend the ranges of density variation by means of the methods described in the next section.

6.2. Variable Filling Density

In this connection the following methods of variable filling density should be mentioned:

(i) There now exists a well-known technique of high-energy beam injection, as well as of plasma gun injection.

(ii) Pulsed neutral gas injection can be used to produce a gradual and slow increase in the ion density of a confined fully ionized plasma [13]. By this method the large burnout powers at high filling densities are avoided.

(iii) A gradual decrease in the ion density by pumping away neutral gas from the wall-near layers of a confined plasma provides means for studying the entire set of plasma states, from the impermeable to the permeable range. This also removes the possible difficulties of plasma start-up at low densities, at the same time as such densities may have to be reached in attempts to rise the temperature by means of a limited heating power.

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ON THE IONIZATION AND BURNOUT PROCESSES OF A MAGNETICALLY CONFINED PLASMA

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The particle and heat balance during plasma start-up are investigated, to specify the conditions for reaching various ion density ranges and high plasma temperatures in cases of a limited heating power. Particular attention is paid to the permeable-impermeable transition regime of plasmas being subject to Ohmic heating and confined in closed or open bottles with a main poloidal field. The ionization and burnout conditions are found to depend critically on the confinement and the filling density. They become optimal in closed bottles under symmetric and stable conditions, where the transition into a fully ionized state should be reached even at moderately large ionization rates, burnout powers and currents. Start-up methods based on constant as well as om variable filling densities are discussed as means of ion density control.

Key words: Magnetic confinement, ionization, plasma start-up, burnout.