PFC/JA-94-19

Effect of Diamagnetic and E x B Drifts on Divertor Plasma Flows

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July 1994

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This work was supported by the US Department of Energy under contract DE-FG02-91ER-54109.

Submitted for publication in: Physics of Plasmas

Effect of Diamagnetic and E×B Drifts on Divertor Plasma Flows

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Abstract

Due to the influence of the diamagnetic and $\vec{E} \times \vec{B}$ drifts affecting the inertia term in the plasma momentum equation a strong variation of the plasma pressure along the magnetic field lines can appear similar to experimental observations of the "detached divertor" regimes. Thus, drift effects can be considered as an alternative physical mechanism of divertor plasma detachment. Drifts can play an especially important role when the mechanism of plasma detachment employing plasma-neutral interaction becomes ineffective e. g. low density plasmas transparent to the neutrals or for large (ITER scale) tokamaks where the efficiency of neutral-wall interaction is considerably reduced.

I. Introduction

Recent experiments on most diverted tokamaks (JET, DIII-D, C-MOD, ASDEX/U, JT-60U) have found "detached divertor" operating regimes[1-5]. These regimes are characterized by a plasma pressure drop in the divertor volume along the magnetic field lines; very low heat and particle fluxes to the divertor plates; high radiation losses from the X-point region, and low plasma temperature and density in front of the divertor plate. Because of very low heat loads to the divertor plates the detached divertor regimes look quite promising from the ITER divertor design point of view. One of the most attractive features of these regimes is the plasma pressure drop in the divertor region.

In Ref. 6 it was shown that detached divertor regimes can be related to the reduction of the heat flux into the hydrogen recycling region below some critical value due to impurity radiation losses. This automatically leads to a decrease of the plasma flux onto the target and a plasma pressure drop which (depending on plasma parameters and divertor geometry) is driven by neutral gas viscosity and neutral pressure[7] or the friction between the plasma and the neutral gas[8].

Although this model of the divertor plasma detachment qualitatively fits the majority of the experimental observations in current tokamaks this mechanism of plasma detachment probably can not be realized easily in ITER size tokamaks. It also raises problems with the interpretation of the plasma pressure drop in detached divertor regimes at low plasma densities where the divertor plasma becomes transparent for the neutrals.

In this paper we show that due to the influence of self-consistent diamagnetic and $\vec{E} \times \vec{B}$ drifts a strong variation of the plasma pressure along the magnetic field lines can appear in a manner similar to that observed in the experiments. Thus drift effects can be considered as an alternative physical mechanism of the divertor plasma detachment.

In this context it is necessary to mention here that the effect of prescribed constant radial electric field on plasma flow near the target was recently discussed in Ref. 9.

In Chapter II a simplified set of plasma fluid equations is specified and the physical mechanism of plasma pressure drop along the magnetic field lines due to the influence of the drifts is discussed. The solutions of these equations are given in Chapter III. The results obtained in the paper as well as their possible relation to detached divertor regimes are discussed in Chapter IV.

II. Equations

We will use slab geometry and assume that x, y and z are the "radial", "poloidal" and "toroidal" coordinates, so that the magnetic field vector is in (y,z) plane. To simplify the problem we only will analyze plasma flows in the divertor volume assuming that the width of the scrape-off layer as well as the profile of the heat flux into the divertor region is determined upstream by heat and particle sources and anomalous heat and particle transport across magnetic field lines (see Fig. 1). Since the divertor region is relatively small (in comparison with the whole scrape-off layer) we neglect at first the influence of anomalous transport processes, and particle and momentum sources in this region so that the steady state plasma and electron momentum equations along magnetic field lines and the continuity equation can be written as

$$\nabla \cdot \left(Mn\vec{V}_{\perp} V_{\parallel} \right) + \nabla_{\parallel} \left(n(T_e + T_i) + MnV_{\parallel}^2 \right) = 0 \,, \qquad (1)$$

$$\nabla_{\parallel}(nT_e) + \alpha n \nabla_{\parallel} T_e - en\nabla_{\parallel}\Phi = 0 \,, \qquad (2)$$

$$\nabla \cdot (n\vec{V}) = 0 \,, \qquad (3)$$

where n and $\vec{V}$ are the plasma density and velocity; $\vec{V}_{\perp}$ and $V_{\parallel}$ are the plasma velocities across and along the magnetic field lines; M is the ion mass; T_i and T_e are the ion and electron temperatures respectively; α is the thermal force coefficient (depending on the plasma effective charge α can vary in the range $0.71 \div 1.5$); $\nabla_{\parallel}$ is the gradient along magnetic field line, $\nabla_{\parallel}(\ldots) = b\nabla_y(\ldots)$, $0 < b = B_P/B_T << 1$; B_P and B_T are the poloidal and toroidal magnetic field strengths. We assumed in Eq. (2) that there is no electric current. Note that Eq. (1) is valid even in the presence of neutral recycling provided that neutral gas does not bring additional momentum into the plasma.

Taking into account $\vec{E} \times \vec{B}$ and diamagnetic drifts and the smallness of b one gets for the velocities V_x and V_y

$$V_x = -\frac{c}{B_T}\left\{\frac{\partial\Phi}{\partial y}+\frac{1}{n}\frac{\partial}{\partial y}\left(\frac{nT_i}{e}\right)\right\},$$

$$V_y = \frac{c}{B_T}\left\{\frac{\partial\Phi}{\partial x}+\frac{1}{n}\frac{\partial}{\partial x}\left(\frac{nT_i}{e}\right)\right\}+bV_{\parallel}\ ; \tag{4}$$

here we assumed that the toroidal magnetic field is in the z direction.

One also has to impose boundary conditions for the potential drop in the sheath. Since we assumed no electric current one has

$$e(\Phi_t-\Phi)\Big|_{sheath} = -\gamma T\Big|_{sheath}\ , \tag{5}$$

where $\gamma \approx 3.5$ and Φ_t = const. is the target potential.

Equations (1)-(5) have to be closed by the energy balance equations. In their general form these equations can be written as

$$\nabla\cdot\vec{q}_e = -Q_e\ ,\quad \nabla\cdot\vec{q}_i = -Q_i\ , \tag{6}$$

where $\vec{q}_e$ and $\vec{q}_i$ are electron and ion heat fluxes, and Q_e and Q_i are electron and ion energy sinks.

Before proceeding with the solution of Eqs. (1)-(6) let us try to understand qualitatively under what conditions one can have a strong plasma pressure variation along the magnetic field lines which is not balanced by the conventional inertia term $MnV_{\parallel}^2$.

For our qualitative analysis it is useful to introduce an integral plasma momentum Π. Integrating Eq. (1) over the "radial" x coordinate and assuming that there is no outflow of parallel plasma momentum onto the side walls one has

$$\Pi(y) \equiv \int dx\left(P + MnV_yV_{\parallel}/b\right) = \text{const.} \tag{7}$$

where, $P = n(T_i + T_e)$, is the total pressure. While the first term in the integral expression in Eq. (7) is always positive, the second one (which contains cross field convection due to the drifts) can be positive or negative. Let us assume, for example, that in the divertor region plasma density and temperatures, and electrostatic potential are symmetric with respect to x=0 (assumed to be the separatrix location). We represent the plasma flux along the y coordinate as a sum of a symmetric, $\Gamma_g(x,y)$, and an asymmetric term, $\Gamma_u(x,y)$

$$nV_y = \Gamma_g(x,y) + \Gamma_u(x,y) . \tag{8}$$

Then from Eq. (4) one has

$$bnV_{\|} = \Gamma_g(x,y) + \Gamma_u(x,y) - \Gamma_{u,drift}(x,y) , \tag{9}$$

where the last term on the right hand side is related to $\vec{E} \times \vec{B}$ and diamagnetic drifts and describes the part of the asymmetric component of the plasma flow along the magnetic field lines. Introducing the symmetric function $\zeta_g(x,y)$ through $\Gamma_{u,drift}(x,y) = (1 + \zeta_g(x,y))\Gamma_u(x,y)$, one can re-write expression (7) for the integral plasma momentum Π in the form

$$\Pi = \Pi_p + \Pi_g + \Pi_u = \text{const.} , \tag{10}$$

where

$$\Pi_p = \int dx P , \quad \Pi_g = \int dx \frac{\Gamma_g^2}{nb^2} , \quad \Pi_u = -\int dx \frac{\zeta_g \Gamma_u^2}{nb^2} . \tag{11}$$

In a 1D approach the sum $\Pi_p + \Pi_g$ can be interpreted as the conventional plasma momentum $P + MnV_{\|}^2$. The term Π_u is the result of the plasma flow asymmetry and can only appear in a 2D treatment of the problem. The most important things are that: 1) while the conventional terms Π_p and Π_g are always positive, the term Π_u can be positive or negative depending on the sign of ζ_g and 2) since the asymmetric part of the plasma flow does not contribute to plasma flux onto the target, the magnitude of Π_u is not limited by plasma recycling although the magnitude of Π_g can be.

Let us consider how the plasma pressure drop in detached divertor regimes can be interpreted by the expression (10). Since the plasma pressure drop in these regimes can not be explained by supersonic plasma flow near the target and, moreover, the particle flux decreases when we are going from the target to the mid plane we can neglect the influence of the Π_g term in Eq. (10). Then a plasma pressure drop can only be balanced by Π_u which is possible for $\zeta_g > 0$ or an increasing asymmetric part of nV_y or $V_{||}$, or both. These asymmetries are related to $\vec{E} \times \vec{B}$ and diamagnetic drifts and, hence, their magnitudes will increase in detached divertor regimes when we are going from the target to the mid plane. Therefore, it is possible in principle to balance an increasing Π_p term by the decrease of Π_u provided regimes with $\zeta_g > 0$ are allowed (which will be shown below). A crude estimation of the variation of the electrostatic potential across the scrape-off layer, $\delta\Phi$ which allows for the balance of the Π_p term by Π_u can be obtained from Eq. (7). Assuming that $V_{||}$ is of the order of sound speed one gets

$$e\,\delta\Phi / T \sim \Delta / \rho_{pi} , \qquad (12)$$

where Δ is the scrape-off layer width and ρ_{pi} is the poloidal ion gyroradius.

III. Solution

Since our objective is focussed on showing that Eqs. (1)-(6) allow solutions with $\zeta_g > 0$ and strong concomitant variation of plasma momentum along the magnetic field lines, we will consider the class of the solutions which can be represented in the form

$$T_e(x,y) = T_e(n) , \quad T_i(x,y) = T_i(n) , \tag{13}$$

$$\Phi(x,y) = \varphi(n) + \varphi_d(x) , \tag{14}$$

where plasma density n depends on both x and y coordinates and $T_{e,i}(n)$, $\varphi(n)$ and $\varphi_d(x)$ are unknown functions. We will see below that the shear of the "radial" electric field driven by the function $\varphi_d(x)$ gives rise to the plasma pressure drop along the magnetic field lines.

We will show that the two distinct simplified limits which are of physical interest, $T_i = 0$ and $T_i = T_e$, give rise to very similar results; therefore we will only consider the case $T_i = 0$ ($T_e = T$) below.

Substituting expressions (4), (13), (14) into the continuity equation one has

$$V_x = -\frac{c}{B_T}\frac{\partial\varphi}{\partial n}\frac{\partial n}{\partial y} , \tag{15}$$

$$V_y = \frac{c}{B_T}\frac{\partial\varphi}{\partial n}\frac{\partial n}{\partial x} + \frac{g(x)}{n} , \tag{16}$$

$$bV_{\parallel} = -\frac{c}{B_T}\frac{d\varphi_d}{dx} + \frac{g(x)}{n} , \tag{17}$$

where $g(x)$ is a function arising from the integration of the continuity equation over the y coordinate. Then, making the following normalization of the functions $e\varphi \rightarrow M\varphi'$, $T \rightarrow MT'$, $g \rightarrow bg'$ and $(ebB_T/Mc)x \rightarrow x'$ (to simplify the notation we will henceforth omit to keep the prime) one obtains from plasma momentum equation (1)

$$-\frac{dg}{dx}\frac{d\varphi}{dn} + \frac{d^2\varphi_d}{dx^2}n\frac{d\varphi}{dn} = \frac{g^2}{n^2} - \frac{d(nT)}{dn} , \tag{18}$$

and, keeping in mind *anzatz* (13), (14), the electron momentum equation can be written in the form

$$T + (\alpha + 1) n \frac{dT}{dn} = n \frac{d\varphi}{dn} . \tag{19}$$

Let us consider the solutions of Eqs. (18), (19) for the special case

$$g(x) \equiv g_0 n_0 = \text{const.}, \qquad \frac{d^2 \varphi_d}{dx^2} \equiv \psi_d = \text{const.}, \tag{20}$$

where the constants n_0 and g_0 are the effective plasma density and sound speed. (Below we will use n_0 as a normalization constant to simplify the writing of the solution Eqs. (18), (19)).

Putting expressions (19), (20) into Eq. (18) one has

$$\psi_d \left(\frac{d(nT)}{dn} + \alpha n \frac{dT}{dn} \right) = g_0^2 \left(\frac{n_0}{n} \right)^2 - \frac{d(nT)}{dn} . \tag{21}$$

Note, that to get the analogous equation for the more common case $T_i(n) \neq 0$ one has to replace in Eq. (21) the pressure nT by $n(T_i + T_e)$. For the case $T_i = T_e$ by a simple renormalisation of the constants α and g_0 one can reduce the common equation to Eq. (21).

The general solution of Eqs. (19), (21) can be written in the form

$$T(n) = T_\beta \left\{ (n_0 / n)^2 + \sigma (n_0 / n)^\beta \right\} , \tag{22}$$

$$\varphi(n) = T_\beta \left\{ \frac{1 + 2\alpha}{2} (n_0 / n)^2 - \sigma \frac{1 - \beta(1 + \alpha)}{\beta} (n_0 / n)^\beta \right\} , \tag{23}$$

valid for $\beta \neq 2$, where $\sigma = \pm 1$,

$$\beta \equiv (1 + \psi_d) / (1 + (1 + \alpha)\psi_d) , \tag{24}$$

$$T_\beta = g_0^2 (1 - \beta(1 + \alpha)) / (\alpha(2 - \beta)) . \tag{25}$$

The constant n_0 is chosen either to equilibrate the two terms in the brackets in Eq. (22) for $\sigma = +1$ or to cancel them for $\sigma = -1$.

For $\psi_d = -1/(1+2\alpha)$ (i. e. $\beta = 2$) the solution of Eqs. (19), (21) is

$$T(n) = T_2 (n_0 / n)^2 \ln(n / n_0) , \tag{26}$$

$$\varphi(n) = -T_2 (n_0 / n)^2 \left\{ \frac{1}{4} + \frac{1+2\alpha}{2} \ln(n_0 / n) \right\} , \tag{27}$$

where

$$T_2 = g_0^2 (1+2\alpha) / \alpha . \tag{28}$$

We are interested in "detached" solutions $dT / dn > 0$, which are similar to those obtained in the experiments, and, of-course, $T > 0$.

For $\sigma = -1$ these solutions only exist for $\psi_d > -1/(1+\alpha)$ (i. e. $\beta > \beta_* = 1/(1+\alpha)$). They are restricted by the inequalities

$$n_0 < n < n_* = n_0 (|\beta|/2)^{1/(\beta-2)} , \tag{29}$$

$$0 < T < T_* = g_0^2 \frac{\beta(1+\alpha) - 1}{\alpha\beta} (|\beta|/2)^{2/(2-\beta)} . \tag{30}$$

For $\sigma = 1$ these solutions only exist for $-1 < \psi_d < -1/(1+\alpha)$ (i. e. $\beta < 0$) and are restricted by the inequalities

$$n > n_* , \qquad T > T_* . \tag{31}$$

Thus, the solutions we are interested in ($dT / dn > 0$ and $T > 0$) are only possible for $\psi_d > -1$. For the range $-1 < \psi_d < -1/(1+\alpha)$ these solutions correspond to $\sigma = 1$ and are restricted by Eq. (31). For $\psi_d > -1/(1+\alpha)$ they correspond to $\sigma = -1$ and are restricted by Eqs. (29)-(30).

The expressions (22)-(31) determine the dependencies of plasma temperature and electrostatic potential on plasma density but do not specify their space dependencies. To get these dependencies one has to solve the energy balance equations (6). We will assume here that the electron energy sink term Q_e is determined by impurity radiation losses and has the form $Q_e = Q_{rad}(n, T)$.

There are two limits when the energy balance equation can be drastically simplified. In the first limit of high plasma temperature, the heat transport is determined by electron heat conductivity along the magnetic field lines ($\vec{q} = -\kappa_e \nabla_{||} T$, where $\kappa_e \propto T^{5/2}$ is the heat conduction coefficient). In this case, from the energy balance equation one gets

$$\left(\kappa_e b \frac{dT}{dy}\right)^2 = q_{up}^2 - \int_n^{n_{up}} \left(2\kappa_e(T(n'))Q_{rad}\left(n', T(n')\right)\frac{dT(n')}{dn'}\right) dn' , \qquad (32)$$

where $q_{up} = q_{up}(x)$ and $n_{up} = n_{up}(x)$ are the upstream heat flux along the y coordinate and plasma density respectively. In the second limit of low plasma temperature, heat transport is determined by convection. In this case, from the energy balance equation and taking into account expressions (15), (16) one has

$$|g_0| \frac{5}{2} \frac{n_0}{n} \frac{dT}{dn} \frac{dn}{dy} = Q_{rad}(n, T) . \qquad (33)$$

Using relation (22) and solving Eqs. (32), (33) one finds the spatial dependence of plasma temperature and density.

Let us now consider the boundary condition (5) which, taking into account relations (13), (14), (20), can be re-written in the form

$$\varphi_0 - \psi_d \frac{(x - x_0)^2}{2} = \left(-\gamma T(n) + \varphi(n)\right)\Big|_{sheath} , \qquad (34)$$

where φ_0 and x_0 are the constants determined by the target potential and the $\varphi_d(x)$ function respectively. It is easy to see that the boundary condition (34) actually determines the dependence of the plasma density in front of the target, $n_t(x)$, on the x-coordinate. We will assume that the separatrix coordinate x_{sep} corresponds to $x_{sep} = x_0 = 0$.

For $\psi_d > -1/(1+\alpha)$ (e. i. $\beta > \beta_*$) Eq. (34) can be written in the form

$$-\varphi_0 + \psi_d \frac{x^2}{2} = g_0^2 \frac{\beta(1+\alpha)-1}{\alpha(2-\beta)} \left\{ \omega_- \left(\frac{n_0}{n_t^{(-)}} \right)^\beta - \omega \left(\frac{n_0}{n_t^{(-)}} \right)^2 \right\} , \qquad (35)$$

where $n_t^{(-)}(x)$ is the plasma density in front of the target; $\omega = \gamma - (1+2\alpha)/2$ and $\omega_- = \gamma - \left(\beta(1+\alpha)-1\right)/\beta$. One can see that since $\gamma > 1+\alpha$ and $\beta > \beta_*$, both ω and ω_- are always positive.

For $\psi_d > -1/(1+\alpha)$ the plasma density variation is restricted by (29). That means that the solution of Eq. (35) only exists within a finite range of the x-coordinate, $-\delta_- < x < \delta_-$. One can show that for $\gamma > 1+\alpha$ and $\beta > \beta_*$ the right hand side of Eq. (35) is positive and reaches its maximum for $n_t^{(-)} = n_-$ where $n_* > n_- = n_0(\beta\omega_-/2\omega)^{1/(\beta-2)} > n_0$. Therefore a solution of Eq. (35) is only possible for the width δ_- is restricted by inequality $\delta_-^2 \le \Delta_-^2$ where

$$\Delta_-^2 = \frac{2g_0^2}{|\psi_d|} \frac{(\beta(1+\alpha)-1)}{\alpha\beta} \left\{ 2\omega \left[\frac{\beta}{2}\left(1 - \frac{1}{2\omega}\right) + \frac{1}{2\omega} \right]^{2/(2-\beta)} - 1 \right\}$$

$$< \frac{4g_0^2}{|\psi_d|} \frac{1+\alpha}{\alpha} (\gamma - 1 - \alpha) \ . \qquad (36)$$

Additional restrictions on the solutions (22)-(28) are related to the imposed requirements that in the sheath within the region $-\delta_- < x < \delta_-$ the velocity V_y is directed onto the target and the Bohm criterion (which can be written here as $MV_{||}^2 > T$) is satisfied. In order to get a maximum plasma pressure drop it is necessary to start with the lowest value of $n_t^{(-)}$. Thus we assume that $n_t^{(-)} \ge n_0$ and examine these restrictions for $x = 0$.

From Eqs. (16), (17) one finds that V_y is directed onto the target if $g_0 < 0$ and the Bohm criterion is satisfied since $T(n = n_0) = 0$ while $(V_{\|}(n = n_0))^2 = g_0^2$. Thus solutions (22)-(28) for $\psi_d > -1/(1+\alpha)$ $(\sigma = -1)$ satisfy Eqs. (1)-(3) as well as the boundary conditions.

Let us now estimate the change of conventional total (kinetic plus static) plasma pressure $P_\Sigma = MnV_{\|}^2 + nT$ along the magnetic field lines for $x = 0$. From Eq. (22) one has

$$P_{\Sigma,x=0}(n) = \frac{g_0^2}{\alpha(\beta-2)}\left\{(2\alpha+1)(\beta-1)\frac{n_0}{n} - \left(\beta(\alpha+1)-1\right)\left(\frac{n_0}{n}\right)^{\beta-1}\right\}. \quad (37)$$

It is easy to show that within the range $n_0 < n < n_*$ the right hand side of Eq. (37) has a minimum, $P^{min}_{\Sigma,x=0} = g_0^2$, reached at for $n = n_0$ and a maximum, $P^{max}_{\Sigma,x=0}$, obtained for $\beta < 2$ at $n = n_p$:

$$n_0 < n_p = n_0\left((\beta(\alpha+1)-1)/(2\alpha+1)\right)^{1/(\beta-2)} < n_*, \quad (38)$$

$$\frac{P^{max}_{\Sigma,x=0}}{P^{min}_{\Sigma,x=0}} = \frac{2\alpha+1}{\alpha}\frac{n_0}{n_p}. \quad (39)$$

Note, that for the case $\alpha = 0$ (e. i. $\beta = 1$) from Eqs. (38), (39) one has $P^{max}_{\Sigma,x=0}/P^{min}_{\Sigma,x=0} = 1$. For $\beta > 2$ the maximum occurs at $n = n_*$:

$$\frac{P^{max}_{\Sigma,x=0}}{P^{min}_{\Sigma,x=0}} = \frac{\beta(2\alpha+1)-1}{\alpha\beta}\frac{n_0}{n_*}. \quad (40)$$

One can see from expressions (29), (38)-(40) that the ratio $P^{max}_{\Sigma,x=0}/P^{min}_{\Sigma,x=0}$ increases with increasing β and for $\beta \rightarrow \infty$ one has

$$P^{max}_{\Sigma, x=0} / P^{min}_{\Sigma, x=0} \leq (2\alpha + 1)/\alpha \; . \tag{41}$$

With the thermal force coefficient bounded by $0.71 \leq \alpha \leq 1.5$ this yields a maximum $P_{\Sigma, x=0}$ drop between 3.4 and 2.7 (for the case $T_i = T_e$ one would get maximum $P_{\Sigma, x=0}$ drop between 4.8 and 3.3). The maximum of the ratio $P^{max}_{\Sigma, x=0} / P^{min}_{\Sigma, x=0}$ is achieved for the most negative of the possible ψ_d values. Note, that according to Eq. (17) $V_{||}(x \leq \delta_-, y = 0)$ can be much lower than $V_{||}(x = 0, y = 0)$ due to the effect of the $\vec{E} \times \vec{B}$ drift. Therefore, the ratio $P^{max}_{\Sigma, x\leq\delta_-} / P^{min}_{\Sigma, x\leq\delta_-}$ can in principle be even higher than $P^{max}_{\Sigma, x=0} / P^{min}_{\Sigma, x=0}$.

For $-1 < \psi_d < -1/(1+\alpha)$ (e. i. $\beta < 0$) it is possible to show that at $x = 0$ the Bohm criterion is not satisfied for the solutions considered above. That means that for $-1 < \psi_d < -1/(1+\alpha)$ the possible solutions satisfying the boundary condition and the Bohm criterion have to be within the region $|x| > \delta_+$, where δ_+ is determinded by boundary condition (34) and the Bohm criterion. The investigation of these solutions is beyond the scope of this paper, however it is necessary to mention here that the maximum plasma pressure which can be achieved for these solutions is not restricted, in principle.

Let us discuss now how to extend the solution of Eqs. (1)-(3) for $\psi_d > -1/(1+\alpha)$ outside the regions $-\delta_- < x < \delta_-$ where boundary conditions can not be satisfied for the solutions (23) discussed above. It is easy to see that Eqs. (1)-(3) allow the solution

$$T(x,y) = \text{const.}\,, \; V_{||} = V_{out}(x)\,, \tag{42}$$

and plasma density

$$n = n_{out}(x) \tag{43}$$

which decreases with increasing $|x|$ (we assume here that $V_{||}$ is directed onto the target and the Bohm criterion is satisfied). One can see that it is easy to match the solutions (22) with solutions (42), (43) if

$$n(|x| = \delta_- - 0) = n(|x| = \delta_- + 0) = n_{out}(\delta_-) \,, \tag{44}$$

since in this case $V_x(|x| = \delta_-) = 0$ and there is no link between regions $|x| < \delta_-$ and $|x| > \delta_-$.

IV. Conclusions

It was shown that due to the influence of the shear of the self-consistent diamagnetic and $\vec{E} \times \vec{B}$ drifts affecting the inertia term in the plasma momentum balance a strong drop of the total plasma pressure $P_\Sigma = MnV_{||}^2 + nT$ (by a factor 3-5 or even higher) along the magnetic field lines can be supported in a manner similar to that observed in detached divertor plasma experiments. Therefore, drift effects can be considered as competitive physical mechanisms of the divertor plasma detachment especially important for the cases where the mechanism of plasma detachment employing plasma-neutral interaction becomes ineffective: low density plasmas transparent to the neutrals, or ITER scale tokamaks where the efficiency of neutral-wall interaction is considerably reduced. The typical radial scale length of plasma parameter variation for these regimes is of the order of the poloidal ion gyroradius (see Eqs. (12) and (36)). The solutions obtained here for the simplified plasma fluid equations can be verified by 2D edge plasma modeling including a fuller treatment of electrostatic potential effects.

We further note that in this paper we have only considered the influence of laminar diamagnetic and $\vec{E} \times \vec{B}$ drifts on the inertia term. However, the turbulence at the tokamak plasma edge can be strong enough so that fluctuating quantities can be comparable to the laminar values. In this case one has to consider in Eq. (7) the time averaged $MnV_yV_{||}$ term. The role of fluctuations can be substantial and can lead to further balancing of the plasma pressure drop.

The effect of the $MnV_yV_{||}$ inertia term on the pressure balance may also be important for the treatment of Multifaceted Asymmetric Radiation

From the Edge (MARFE) [10] characterized by strong temperature variation along the magnetic field lines.

Acknowledgments

The authors are gratful to Dr. Peter Catto for his careful checking of the analysis and solution, and stimulating discussions. This work was performed under DOE grant DE-FG02-91-ER-54109 at MIT.

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Figure Caption

Fig.1. Geometry of the problem.

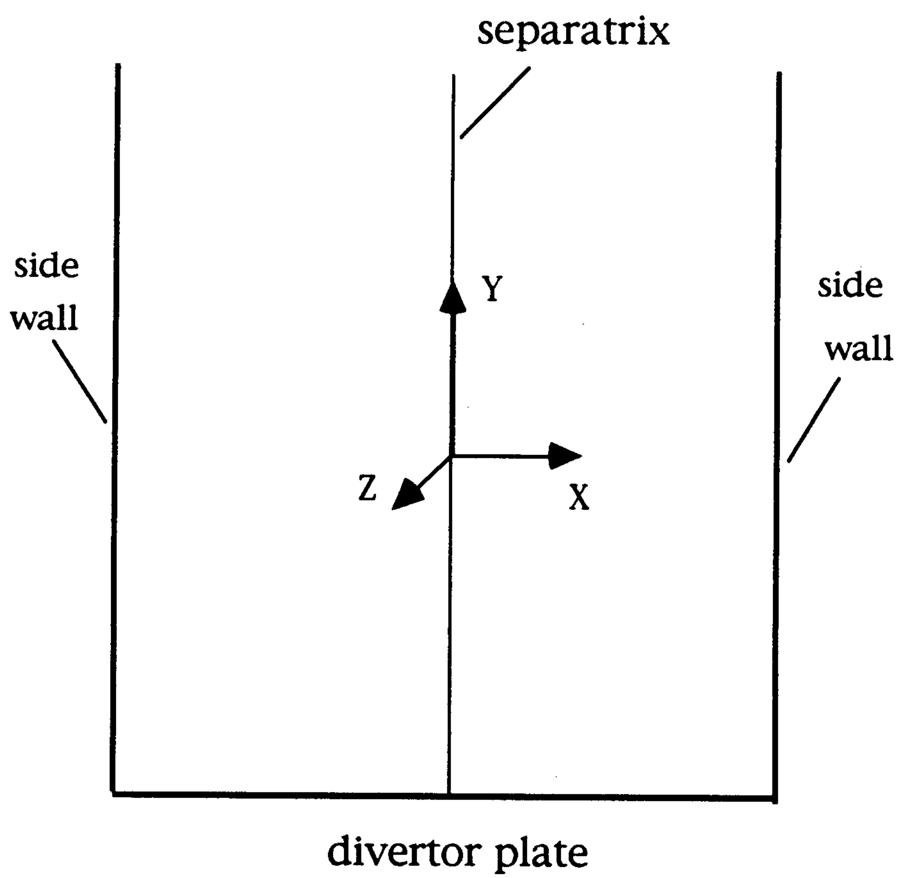

Fig. 1