Classical Drifts in the Tokamak SOL and Divertor: Models and Experiment

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ABSTRACT

Classical drifts offer potential explanations for a variety of physical phenomena such as poloidal asymmetries, possible extra pinch or outward flow of plasma depending on the $B_r$ direction, non-ambipolarity of radial plasma flow and current flow towards the target. Their incorporation into 2D numerical codes is promising to greatly improve their predictive capacity. The paper contains basic analysis of drift and fluid flows, modified boundary conditions, ad-hoc models and summary of experimental results which are widely regarded as having their origin in drift motions. Unresolved issues of plasma transport are highlighted.

1. INTRODUCTION

Guiding centre drifts can significantly influence poloidal asymmetries of the distribution of plasma parameters over the magnetic surface; radial, poloidal and toroidal plasma flows, perpendicular electric conductivity etc. Simple estimates (Section 2.1) show that drifts' relative contribution to both radial and poloidal transport of particles scales as the ratio of ion poloidal Larmor radius, calculated for ion sound speed, $\rho_{s\theta} = c_s / \omega_{i\theta}$, to the radial decay length of plasma parameters. This ratio is often of the order of unity in the tokamak SOL.

The importance of drifts in the edge transport has demanded their inclusion in 2-dimensional codes (see e.g. [1,2,3]). Recent developments in this area were reviewed in [4]. The present paper focuses on basic analysis of drift and fluid flows, boundary conditions in the presence of drifts, ad-hoc models and experimental results which are widely regarded as being (at least partly) caused by the drifts. Qualitative understanding of the phenomena, obtained from simple models, is essential for both general understanding of the SOL behaviour and code development. Throughout the paper, a single null X-point divertor configuration is assumed. For simplicity, formulae are written in the form which assumes the right angle between the poloidal projection of the magnetic separatrixes and the surface of the divertor target.
2. DRIFT FLOWS IN THE TOKAMAK EDGE

2.1 Main drift motions; fluid and guiding centre approaches

The main guiding centre drifts are depicted on Fig.1 for the case of "normal" toroidal field \( B_t \) direction, when the ion \( \mathbf{B} \times \nabla \mathbf{B} \) drift is directed towards the X-point. The direction of all the drifts is reversed for the opposite \( B_t \) direction. The direction of the plasma current, which is assumed to be parallel to the toroidal field on Fig.1 (the helicity of field lines is indicated by the segment of the field line near the outer side), does not affect any equations. Since the fluid approximation, based on momentum conservation equations, is more universal and is more often used in models and numerical codes, it seems appropriate here to demonstrate the relationship between the fluid and guiding centre flows.

The momentum conservation equation reads [5]:

\[
\frac{\partial}{\partial t} (nm \mathbf{V}) + \nabla \tilde{P} = en(\mathbf{E} + \mathbf{V} \times \mathbf{B}) + \mathbf{R}
\]

(1)

In this and the following equations of this section the electric charge \( e \) is assumed to be positive for ions and negative for electrons. \( \mathbf{R} \) is the friction force exerted on particles of a specified type by another, and \( \tilde{P} \) is the stress tensor. Separating contributions by parallel directional momentum flux and parallel chaotic energy (pressure), the stress tensor given in [6], which neglects viscosity and cross products \( nmV_{i\parallel} \), can be written in the form:

\[
\tilde{P} = (p_1 + nmV_{i\parallel}^2)\mathbf{BB} / B^2 + p_{\perp}(\tilde{\delta} \cdot \mathbf{BB} / B^2)
\]

(2)

Here \( \tilde{\delta} \) is a unit tensor. Perpendicular velocity \( V_{\perp} = \mathbf{V} - \mathbf{B}(\mathbf{V} \cdot \mathbf{B}) / B^2 \) can be extracted from the term \( \mathbf{V} \times \mathbf{B} \) of Eq.(1) by vector multiplying it by \( \mathbf{B} \) and using formulae of tensor analysis to open \( \nabla \tilde{P} \). By adding parallel flux density \( nV_{i\parallel} \mathbf{B} / B \) to the perpendicular one, the total flux density for steady state conditions (\( \partial / \partial t = 0 \)) can thus be obtained:

\[
n\mathbf{V} = nV_{i\parallel} \frac{\mathbf{B}}{B} + \frac{n}{B^2} \mathbf{E} \times \mathbf{B} + \frac{1}{eB^2} \mathbf{R} \times \mathbf{B} + \frac{1}{eB^2} \mathbf{B} \times \nabla p_{\perp} + \frac{(p_1 - p_{\perp} + nmV_{i\parallel}^2)}{eB^3} \mathbf{B} \times \mathbf{B} \cdot \nabla \left( \frac{\mathbf{B}}{B} \right)
\]

(3)

The last term in this equation is often ignored in the models. It cannot, however, be neglected in the SOL, where even for isotropic pressure, \( p_1 = p_{\perp} \), the term with \( nmV_{i\parallel}^2 \) can make a large contribution to the ion radial flux, comparable to the radial component of the \( \mathbf{B} \times \nabla p_{\perp} \) flux.

The notion of drift guiding centre flux only makes sense in a quasi homogeneous plasma, where magnetic field components and plasma parameters do not vary substantially on the scale of the ion Larmor radius. The charged particle distribution function can then be expanded into
fast rotations of Larmor circles and averaged guiding centre drift (see e.g. [5] and Refs. therein). The flux due to the superposition of Larmor circle rotations is given by the so called magnetisation flux $\text{curl}\mathbf{K}$, where $\mathbf{K} = -\left(\frac{p_\perp}{eB^2}\right)\mathbf{B}$. Total (fluid) particle flux density can be represented as:

$$n\mathbf{V} = n\langle \mathbf{v}_{dr} \rangle + \text{curl}\mathbf{K} + n\mathbf{V}_{\text{coll}}$$

(4)

Here $\mathbf{v}_{dr}$ is the drift velocity of a Larmor guiding centre and $\langle \ldots \rangle$ denotes averaging over the distribution function; $n\mathbf{V}_{\text{coll}}$ gives particle flux due to electron-ion collisions equal to $\frac{1}{eB^2}\mathbf{R} \times \mathbf{B}$ - the third term on the r.h.s. of Eq.(3). This form of the total particle flux clearly shows that, apart from the guiding centre contribution, $n\langle \mathbf{v}_{dr} \rangle$, and the Coulomb collisions, the remainder of the particle flux is automatically divergence-free as $\text{curl}$ of a vector. In many cases, where only accumulation or sink of particles or energy in space is of interest, it may be more convenient to treat the plasma as an ensemble of Larmor guiding centres. It may also be of benefit in numerical calculations on finite mesh, since the part of the flux which has to be made divergence-free, $\text{curl}\mathbf{K}$, is much greater than $n\langle \mathbf{v}_{dr} \rangle$, thereby causing problems of accumulating errors in the numerical scheme. Elimination of divergence-free terms from the numerical scheme was used as a basis for implementation of drifts into the EDGE2D code at JET [3].

An expression for the drift velocity of an individual particle can be found in [5]. Its integration over the distribution function gives:

$$n\langle \mathbf{v}_{dr} \rangle = n\langle \mathbf{v}_{dr} \rangle \frac{\mathbf{B}}{B^2} + \frac{n}{eB^2} \mathbf{E} \times \mathbf{B} + \frac{p_{1}}{eB^3} \mathbf{B} \times \nabla B + \left(\frac{p_{1} + n m V_1^2}{eB^3}\right) \mathbf{B} \times \mathbf{B} \cdot \nabla \left(\frac{B}{B}\right) + \frac{p_{1}}{eB^3} \left(\mathbf{B} \cdot \text{curl}\left(\frac{\mathbf{B}}{B}\right)\right)$$

(5)

The second term on the right hand side of this equation describes the $\mathbf{E} \times \mathbf{B}$ drift, poloidal and radial components of which are shown on Fig.1, the third term - $\mathbf{B} \times \nabla B$ (or simply $\nabla B$) almost vertical drift with its direction for ions shown on Fig.1 (for the electrons it has an opposite direction), the fourth term is the curvature drift, which is also almost vertical and is approximately in the same direction as the $\nabla B$ drift. The last term makes a small correction to the parallel guiding centre flux which is offset by parallel component of $\text{curl}\mathbf{K}$.

Depending on whether the particle flux is considered according to Eq.(3), or according to Eq.(5) with Eq.(4) implied, the plasma is treated in fluid or guiding centre approximations. As one can see, the main difference between the two approaches (apart from the term with ion-electron friction, missing in Eq.(5), which could, however, be easily added by virtue of the fact
that Coulomb collisions are predominantly distant and only slightly disturb the Larmor orbit) is that the fluid form (3) explicitly contains a large pressure gradient term: the diamagnetic flux \( \mathbf{B} \times \nabla p_{\perp} / eB^2 \). All other terms with pressure in both Eqs. (3) and (5) are much smaller, only of the order of \( \lambda_p / R \) of the diamagnetic flux, where \( \lambda_p \) is the pressure decay length and \( R \) - major radius. In the literature the diamagnetic flux is often called the diamagnetic drift. This flux, however, is almost entirely dominated by the magnetisation flux, \( \text{curl}\mathbf{K} \), and should not be confused with any guiding centre drift. The difference between these two fluxes, \( (\mathbf{B} \times \nabla p_{\perp} / eB^2 - \text{curl}\mathbf{K}) \), is of the order of \( \lambda_p / R \) of each of them individually. The non-divergence-free part of the diamagnetic flux cannot, of course, be totally ignored. It is of the same order as the \( \nabla B \) and curvature drifts and is responsible for generation of Pfrish-Schlüter currents (Section 5.1). Their origin and spatial distribution in the SOL and divertor has recently been studied by Schaffer et al. [7].

Equivalence between the two approximations can be verified by substituting \( n\langle \mathbf{v}_{dr} \rangle \) from Eq.(5), and an expression for the magnetisation flux:

\[
\text{curl}\mathbf{K} = \frac{\mathbf{B} \times \nabla p_{\perp}}{eB^2} - \frac{2p_{\perp}}{eB^3} \mathbf{B} \times \nabla B - \frac{p_{\perp}}{eB^2} \text{curl}\mathbf{B},
\]

into Eq.(4). Then, with the help of the vector formula: \( \mathbf{B} \times \nabla \mathbf{B} + B\text{curl}\mathbf{B} - \mathbf{B}(\mathbf{B} \cdot \text{curl}(\mathbf{B} / B)) = \mathbf{B} \times \nabla (\mathbf{B} / B) \), Eq.(4) takes exactly the form (3).

Similar equivalence exists between fluid and drift energy fluxes ([5], p.262). If collisions are neglected (this includes parallel collisional heat conduction), then the total energy flux can be written both as a sum of fluid convective and heat fluxes:

\[
\mathbf{q} = \frac{5}{2} nT\mathbf{v} + \frac{5}{2} \frac{nT}{eB^2} \mathbf{B} \times \nabla T,
\]

and as a sum of guiding centre convective flux and a \( \text{curl} \) of a vector:

\[
\mathbf{q} = \frac{5}{2} nT\langle \mathbf{v}_{dr} \rangle - \text{curl} \left( \frac{5}{2} \frac{nT^2}{eB^2} \mathbf{B} \right)
\]

### 2.2 The contribution of drifts to edge transport

If the difference between ion and electron temperatures is ignored, poloidal fluxes for both diamagnetic and \( \mathbf{E} \times \mathbf{B} \) drifts can be estimated as: \( \Gamma_\theta^{dr} = p / eB \lambda_{SOL} \). Parallel flow with the ion sound speed causes poloidal flux: \( \Gamma_\theta^i = n c_s B_\theta / B \). By replacing pressure with \( n m c_s^2 \), the ratio of the two fluxes can be expressed as:

\[
\frac{\Gamma_\theta^{dr}}{\Gamma_\theta^i} \approx p_{\perp} / \lambda_{SOL}.
\]
where \( \rho_{s\theta} = c_s / \omega_{i\theta} \) and \( \omega_{i\theta} = eB_\theta / m_i \) - ion poloidal gyro frequency.

The same scaling can be deduced for the ratio of radial components of diamagnetic and \( \mathbf{E} \times \mathbf{B} \) drift fluxes to the anomalous radial transport in the scrape-off layer [8]. Radial flux due to the diamagnetic and \( \mathbf{E} \times \mathbf{B} \) drifts can roughly be estimated as \( \Gamma_r^{dr} \approx p / eBa \), where minor radius \( a \) serves as a characteristic length of poloidal variation of plasma pressure in the SOL. Radial anomalous particle flux is \( \Gamma_r^{an} \approx nD_\perp \lambda_{SOL} \). By employing standard expression for the SOL width: \( \lambda_{SOL} = \sqrt{D_\perp qR / c_s} \) (see e.g. [9]), their ratio can also be expressed as:

\[
\frac{\Gamma_r^{dr}}{\Gamma_r^{an}} = \frac{\rho_{s\theta}}{\lambda_{SOL}}
\]

(10)

Since this ratio is typically of the order of unity in the scrape-off layer, drift fluxes can strongly affect the overall flux pattern in this region.

According to the above scaling, one should expect a smaller drift contribution to the plasma transport in low power high density discharges, which have low temperature throughout the SOL and divertor. As a function of plasma temperature and radial diffusion coefficient, the \( \rho_{s\theta} \lambda_{SOL} \) ratio scales as \( T^{3/4} D_\perp^{-1/2} \) (a relation \( \lambda_{SOL} = \sqrt{D_\perp qR / c_s} \) is implied here), and a very steep dependence of \( D_\perp \) on \( T \) is needed to offset the effect of the \( T^{3/4} \) dependence. Experimentally found dependence of \( D_\perp \) on \( T \) is consistent with the scaling \( D_\perp \sim T^\alpha \), where \( \alpha = 0 \sim 1 \) [10,11]. Ionisation of neutrals in the scrape-off layer at high density (low temperature) provides an additional widening of the SOL, further reducing the \( \rho_{s\theta} \lambda_{SOL} \) ratio. In the experiment, the scrape-off layer usually gets wider in high density discharges.

On the other hand, even during the plasma detachment from the target, drift contributions should make a significant impact on plasma behaviour in the main SOL provided it is sufficiently "hot", and large poloidal variation of electron temperature exists which gives rise to e.g. radial \( \mathbf{E} \times \mathbf{B} \) drift. So far conditions with "hot" SOL plasma and detachment from the target have not been realised in the experiment [12].

3. BOUNDARY CONDITIONS AT THE TARGET IN THE PRESENCE OF DRIFTS

In the absence of drifts, according to Chodura [13] and Riemann [14] the ion parallel flow reaches or exceeds the ion speed at the entrance to the magnetic pre-sheath layer (m.p.s.). This criterion is deduced from the requirement of smooth non-oscillating potential drop towards the surface of the target, and represents an extension of the Bohm criterion [15] for the case of magnetic field oblique to the surface.

In the presence of poloidal \( \mathbf{E} \times \mathbf{B} \) drift the boundary condition on the minimal parallel velocity at the m.p.s. entrance is modified. Stangeby, Chankin and Hutchinson [16,17] have shown that it is the combined effect of the \( \mathbf{E} \times \mathbf{B} \) drift velocity and parallel velocity with which ions approach the surface that has to achieve a certain minimal value to ensure a smooth
transition of electric potential through the m.p.s and the Debye sheath. For the simplest case of flux surfaces perpendicular to the material surface in the poloidal cross section, the approximate form of the boundary condition on the minimal (in absolute value) Mach number of parallel plasma flow (the complete form of the criterion derived in [16] is too cumbersome to present here), which only includes the effect of the poloidal $E \times B$ drift is:

\[
\begin{align*}
M_{in} &= 1 - M_{E,in} \\
M_{out} &= -1 - M_{E,out}
\end{align*}
\]  

(11)

for inner and outer targets. Here $M - V_i / c_s$, $M_E = \pm E_r / B_\theta c_s$ and it is assumed that positive velocities (both parallel and perpendicular) are directed from the outer to the inner side (see $V_i$ in Fig.1). Therefore, the sign of $M_E$ is negative for normal $B_t$ direction (poloidal $E \times B$ drift is directed towards the outer target) and positive for reversed $B_t$ direction. The parallel ion velocity has to be adjusted to compensate for the velocity of the poloidal $E \times B$ drift towards/away from the target, so that the total velocity with which ions approach the surface would be equal to $c_s \cdot B_\theta / B$, unless there are specific reasons why this velocity should be even higher in absolute value (such as the onset of a free-standing sonic transition some distance away from the target). The above boundary condition has also been used by Cohen and Ryutov [18].

According to Eq.(11), the poloidal diamagnetic flux does not affect the boundary condition on the minimal parallel ion velocity. This is because the diamagnetic flux is almost divergence-free, as was demonstrated in Section 2.1. Moreover, this conclusion can even be extended to the m.p.s. layer, where drift approximation is not valid, as was shown by Chankin and Stangeby [19]. What happens then to the poloidal diamagnetic flows when they come closer to the target? After reaching the entrance to the m.p.s., the diamagnetic flows cross through the magnetic pre-sheath in the direction normal to the magnetic surfaces and then continue "poloidally" on private magnetic surfaces [19]. The neglect of the diamagnetic flux in the boundary condition has, however, recently been disputed by Claassen and Gerhauser [20,21].

In numerical codes the criterion on the minimal ion velocity is often imposed on the total poloidal ion velocity before the m.p.s. entrance. The contribution from the ion poloidal diamagnetic flux then must be included, keeping in mind, of course, that it does not reach the target but, after entering the m.p.s., is diverted along the target surface. The boundary condition then takes the form:

\[
V_{i\theta} = \frac{B_\theta}{B} M c_s + \frac{B_\theta}{enB^2} \nabla_r p_t
\]  

(12)

Cohen and Ryutov [22] considered the boundary condition for the current density towards the target, and their conclusion is essentially that the relation between the parallel
current density and potential drop across the magnetic pre-sheath and the Debye sheath is unaffected by the drifts:

\[ j_{\parallel} = en_{s}\left[1 - \exp\left(e(\Phi - \Phi_{f})/T_{e}\right)\right] \]

(13)

Here \( \Phi \) is the potential before the m.p.s. entrance, and \( \Phi_{f} \) is the floating potential (see e.g. [9]). Again, for the poloidal component of the current density at the m.p.s. entrance, one should correct for the diamagnetic current density:

\[ j_{0} = \frac{B_{0}}{R} en_{s}\left[1 - \exp\left(-\frac{e(\Phi - \Phi_{f})}{T_{e}}\right)\right] + \frac{B_{0}}{B^{2}} \nabla_{r} p \]

(14)

Boundary conditions for the parallel energy flux at the m.p.s. look very similar to the ones in the absence of drifts (see e.g. [23]), and are formulated in [3]:

\[ \frac{5}{2} p_{\alpha} M c_{s} + q_{1} = \beta p_{\alpha} M c_{s} \]

(15)

Here \( q_{1} = k_{B} \nabla_{r} T_{a} \) is the parallel heat flux density of species \( \alpha \). Coefficient \( \beta \) is \( 5/2 \) for ions and \( (2 + e\Phi / T_{e}) \) for electrons.

4. EXPERIMENTS WITH TOROIDAL FIELD REVERSAL: EXPLANATION OF IN-OUT ASYMMETRIES THROUGH THE INFLUENCE OF DRIFTS

The direction of the toroidal field determines the direction of drift flows. It is therefore widely accepted that poloidal asymmetries in the SOL, which are sensitive to the field direction, are caused by the drifts.

4.1 Summary of experimental observations

The average shift in the asymmetries, independent of the field direction, is for denser and colder plasma at the inner target. The temperature asymmetry in favour of the outer target is explained by higher total power flow through the outboard part of the magnetic surface [24]. This is generally attributed to the geometrical toroidal effect of the surplus of its area over the innerboard one. The Shafranov shift and poloidal asymmetry in anomalous heat transport in favour of the outer side [25] may also contribute. As pressure \( p_{e} = n_{e} T_{e} \) tends to reach equilibrium along the field lines: \( p_{e}^{in} = p_{e}^{out} = p_{e} \), higher density \( (n_{e}^{in} = p_{e} / T_{e}^{in}) \) is achieved at the inner target, while higher power \( (P_{target}^{out} - n_{e}^{out}(T_{e}^{out})^{3/2} = p_{e}(T_{e}^{out})^{1/2}) \) flows towards the outer target.
The effect of the toroidal field direction on the power asymmetries, as seen in experiment, is the following. Excessive power lead to the outer strike zone is usually observed in the "normal" toroidal field direction (ion $\nabla B$ drift directed towards the target in single null X-point discharges), whereas for the "reversed" toroidal field (ion $\nabla B$ drift is directed away from the target) power sharing between the targets is much more symmetric [26-35,8]. This tendency is illustrated by Fig.2, representing results on power to the target and radiated power in-out asymmetries in JT-60U obtained by Asakura et al. [30]. At medium and high densities the total power flowing to the outer divertor branch ($P_{\text{target}} + P_{\text{rad,div}}$), after the local strike radiation is taken into account, is slightly larger than that to the inner divertor branch. Such a nearly equal balancing of the total power sharing between the two divertor branches is not a universal case, however, and can vary depending on the machine/regime parameters. More symmetric heat load to the target in reversed $B_T$ discharges at medium and high densities is explained by the asymmetry in the radiated power, which is inboard dominated in the normal and outboard dominated - in the reversed $B_T$ plasmas. Possible explanations for inboard dominated $P_{\text{target}}$ at low $\beta_p$ in reversed $B_T$ case will be discussed in Section 4.2. Power to the target and radiated power asymmetries are substantially suppressed at very high densities.

Asymmetries in plasma parameters at the target in the field reversal experiments were studied in [35-41,8]. Ref. [35] by Hutchinson et al. provides a good illustration of tomographic reconstruction of radiative power in Alcator C-Mod. Fig.3, replicated from this Ref., shows radiative loss pattern in Ohmic discharges for the two field configurations at $\bar{\beta}_p = 1 \times 10^{20} \text{ m}^{-3}$, which is a medium density for this machine. The normal field plasma is dominated by radiation from the inner divertor, while for the reversed field plasma the radiation zone is shifted towards the outer divertor. Langmuir probe measurements and the $D_\alpha$ emission revealed much denser and cooler plasma at the inner side with higher recycling of neutrals there. The sign of $n_e$ and $T_e$ asymmetries changed following the field reversal, but the averaged shift in the asymmetries, independent of the field direction, was for denser and colder plasma at the inner side, as expected. For higher density, the effects of the field reversal were found to be much less dramatic.

A dedicated series of L-mode discharges with both $B_T$ directions and a wide variation of the toroidal field and plasma current has been performed in JET [42,38,8]. The results are broadly consistent with those obtained from other machines. In/out ratios of radiated power, $H_\alpha$ emission and peak ion saturation current density as a function of $q_95$ are shown in Fig.4 [8]. The field reversal has largely eliminated strong asymmetries in these parameters, and the distribution of $H_\alpha$ and $P_{\text{rad}}$ has become slightly shifted to the outer side. Fig.5 [8] demonstrates profiles of $j_{sat}$, $T_e$ and $n_e$ obtained by target Langmuir probes during the diagnostic radial sweep of the X-point for the pair of discharges with $q_95=3.6$. More symmetry is achieved in the plasma parameters distribution between the targets in the reversed field plasmas due to shifts in $n_e$ and $T_e$ distributions in the opposite direction: $n_e$ from the inner to the
outer side, and $T_e$ from the outer to the inner side. These shifts in plasma parameters have been previously identified as the most robust features of field reversal [36].

At very high densities, closer to the plasma detachment from the targets, the formation of a MARFE state, or density limit disruption, the $B_t$ reversal has much weaker affect on all the divertor asymmetries (see e.g. [8,35,41]), however, the detachment density operation window is narrower in reversed field plasmas [43,44].

4.2. Theoretical considerations

There have been a number of attempts to relate observed shifts in the in-out asymmetries caused by the field reversal, with the effect of classical drifts [25,45-48,18, 34-36,38,8]. Staebler [49] recently demonstrated that large asymmetries can be spontaneously (following a small initial perturbation) generated due to radiation and/or passage of parallel currents in the scrape-off layer.

Poloidal $\mathbf{E} \times \mathbf{B}$ drift, caused by the radial electric field, drives the plasma towards the outer side in the normal, and inner side - in the reversed $B_t$ configuration. Convective power flux associated with this drift is, therefore, in the right direction to explain changes in target power asymmetries [35]. This drift, however, should cause profound changes in poloidal pressure/density distribution which are NOT supported by the experiment. Influence of the poloidal $\mathbf{E} \times \mathbf{B}$ drift on the SOL structure has been first analysed by Tendler and Rozhansky [45] and later in [18,48]. This drift introduces an extra flux of parallel momentum in the perpendicular direction within the magnetic surface. The stress tensor (2) should be corrected for this flux by adding the correspondent $nmV_kV_\perp$ components (this should not, however, affect the validity of Eq.(3) since it will introduce only much smaller terms in it compared to those already included). The sum of parallel components of Eq.(1) for ions and electrons in the simplest case of cylindrical geometry can then be written as:

$$\frac{\partial}{\partial s\parallel} \left( p + nmV_\parallel^2 \right) + \frac{\partial}{\partial s\perp} \left( nmV_\parallel V_\perp \right) = 0$$  \hspace{1cm} (16)

Here $p = p_i + p_e$ and $m$ is ion mass. Derivatives over both parallel and perpendicular coordinates can be replaced by the derivative over the poloidal angle using: $\partial/\partial s\parallel = B_\parallel/ B \cdot \partial/ \partial \theta, \partial/\partial s\perp = \partial/ \partial r \partial \theta$. This leads to the following conservation equation:

$$p + nmV_\parallel^2 + nmV_\parallel V_E = \text{const}(\theta),$$  \hspace{1cm} (17)

where $V_E = \pm E_r / B_\parallel$ and positive velocities (both parallel and perpendicular) are assumed to be directed from the outer to the inner side (see $V_\parallel$ in Fig.1). Therefore, the sign of $V_E$ is negative for normal $B_t$ direction (poloidal $\mathbf{E} \times \mathbf{B}$ drift is directed towards the outer target) and positive for
reversed $B_i$ direction. Diamagnetic flux’s contribution to Eq.(17) is neglected as it is almost completely divergence-free, as discussed in Section 2.1.

Eliminating $V_i$ and $V_E = M_E \epsilon_s$ in Eq.(17) by using boundary conditions (11) and replacing $nmc_s^2$ with pressure $p$, the in-out pressure asymmetry can be obtained:

$$\frac{p_{in}}{p_{out}} = \frac{2 + M_{E,out}}{2 - M_{E,in}}$$

(18)

The pressure asymmetry between the strike zones causes plasma flow along the field lines from high to low pressure side giving rise to a net toroidal momentum of the order of $E_r / B \eta_s$ [45,18]. Thus, instead of the poloidal rotation, the poloidal $E \times B$ drift (which, in fact, is in the direction perpendicular to $B$) generates toroidal rotation and pressure asymmetry between the strike zones. The direction of the toroidal velocity is always along the main plasma current direction, irrespective of the toroidal field direction. In [50] it is argued that toroidal rotation in the same direction can also be driven near the separatrix by anomalous radial transport in the presence of shear of the poloidal rotation.

For normal $B_i$ direction ($M_E$ negative), according to Eq.(18), $p_{in}/p_{out} < 1$, and for reversed $B_i$ direction ($M_E$ positive) $p_{in}/p_{out} > 1$. Since the power flux, associated with the poloidal $E \times B$ drift, is convective, it must be the density asymmetry (not temperature!) that is primarily affected by this drift. The sign of this density asymmetry (as well as associated asymmetries in $P_{rad}$ and $H_\alpha$) is, however, opposite to the one observed in the experiment. Moreover, the inevitable losses of power due to local hydrogen recycling and impurity radiation (both increase with increase of density) should reduce electron temperature at the side to which the plasma is driven by the poloidal $E \times B$ drift, and the expected effect on the $T_e$ asymmetry is again opposite to the experimental trend. Regarding experimental pressure asymmetry, no significant changes were detected in Alcator C-Mod [35], while electron pressure asymmetry in favour of the ion drift side (i.e. against the expected effect of the poloidal $E \times B$ drift) was observed in JET [51]. Thus, except for the target power asymmetries, the poloidal $E \times B$ drift alone would predict changes in all other important asymmetries which are against the experimental trends.

The drift which is in the right direction to explain the $n_e$, $T_e$, $P_{rad}$ and $H_\alpha$ asymmetries, is the radial $E \times B$ drift (see Fig.1), as was first pointed out by Hinton and Staebler [25]. According to the analysis performed in [48], radial $E \times B$ drift should dominate over the poloidal $E \times B$ drift in determining the overall flux pattern in high recycling plasmas. It brings more plasma particles to the inner strike zone in normal $B_i$ discharges, and to the outer strike zone - in reversed $B_i$ discharges. Due to the viscous drag experienced by the return parallel flow along the SOL which is induced by the radial drift fluxes, radial $E \times B$ drift also causes pressure imbalance between the targets, but of the opposite sign compared to the poloidal drift. The
increased density at the divertor to which radial $E \times B$ drift supplies particles, increases hydrogenic recycling energy losses, as well as the impurity radiation. This causes the temperature drop and further increase in radiation losses associated with low temperatures (for low $Z$ impurities). The temperature drop causes further increase density owing to the tendency for pressure equilibration (see equations (73,14) of [48]), thereby providing a positive feedback.

The dependence of electron temperature asymmetry on the $B_t$ direction seems to provide the strongest evidence in favour of the radial $E \times B$ drift. The direction of the convective power flow carried by this drift, however, makes it more difficult to explain the experiment. Consistent explanation of the asymmetries through the radial $E \times B$ drift can only be achieved if local radiation power losses near the strike to which the plasma is driven by the drift, is larger than the convective power flow to it. Whether this is the case can only be established by the detailed 2D code calculations which include drifts.

An explanation for the observed power asymmetries through the poloidal diamagnetic energy flux $\mathbf{q} \sim \mathbf{B} \times \nabla_e T_e$ was offered in [47,34]. When the radial ion temperature gradient $\nabla_e T_i$ is steeper than $\nabla_e T_e$, as was observed in ASDEX Upgrade and often observed on other machines, the poloidal energy flux is directed towards the outer target in normal, and towards the inner target - in reversed $B_t$ plasmas, in agreement with the shifts in the measured asymmetries of power conducted to the target. This explanation, however, is founded on the erroneous assumption that the poloidal diamagnetic flows reach the surface of the target. This is not the case, as was pointed out in Section 3. The diamagnetic flows, after reaching the entrance to the m.p.s., cross through the magnetic presheath in the direction normal to the magnetic surfaces and then continue "poloidally" on private magnetic surfaces. They do not reach the target and therefore cannot affect in-out power and particle flux asymmetries.

Radial and poloidal $E \times B$ drifts seem to be unable to explain the rather symmetric power distribution between the targets in normal $B_t$ configuration at low densities in JT-60U [30] (see Fig.2) and hot ion mode regimes in JET (see e.g. [52]), or even the inner target receiving larger heat flux than the outer one [53,28]. A number of other examples of similar behaviour of power and density asymmetries is presented in [36]. Low density and/or high power discharges present significant challenge for the explanation of divertor asymmetries. Indeed, the SOL plasma in these conditions is relatively "hot", and poloidal asymmetries in $T_e$ should be rather small, thus reducing the effect of the radial and increasing the effect of the poloidal $E \times B$ drift. The latter increases plasma pressure (hence, also power to the target) at the outer side in normal $B_t$ configuration. This is not what is observed in the experiment. To explain the discrepancy, an extra force is needed which would compress the plasma at the inner side. A candidate for such a force - the influence of edge toroidal momentum in the direction of the main plasma current on divertor asymmetries - was suggested in [54]. In the experiments reviewed in [54], the toroidal momentum was largest in low density plasmas. To qualitatively explain the experimental data,
the toroidal momentum has to be introduced into the model as an \textit{external} boundary condition on the parallel velocity at the separatrix. The toroidal momentum which originates in the SOL due to the pressure asymmetry caused by the poloidal $\mathbf{E} \times \mathbf{B}$ drift, discussed above, cannot perform this role since it intrinsically requires higher plasma pressure at the \textit{outer} side in normal $B_t$ discharges, opposite to experimental observations.

Higher power to the inner target in low density reversed $B_t$ discharges in JT-60U shown on Fig.2, can naturally be explained by the poloidal $\mathbf{E} \times \mathbf{B}$ drift which drives the plasma to the inner side. The toroidal momentum should be weaker in the reversed $B_t$, since in the field reversal experiments in JT-60U (also in JET) the toroidal field and plasma current were reversed simultaneously, so that counter-injection was applied in the reversed, whereas co-injection - in normal $B_t$ discharges. It has to be noted, however, that the assumption that the drift (in this case, the poloidal $\mathbf{E} \times \mathbf{B}$ drift) which drives the plasma towards one particular target, increases the power flow to this target, is only valid in low density "isothermal" (no poloidal variation of temperature) plasmas. In high density plasmas, as discussed above, the increased local radiation has to be taken into account, and no simple conclusions on whether the drift towards the target can increase power flow towards this target, can be made.

As was pointed out in Section 4.1, there exists convincing evidence for a much weaker effect of the field reversal on divertor asymmetries at very high densities, when the plasma in the scrape-off layer is rather "cold". Such a trend can be explained from the scaling $\rho_s \theta / \lambda_{SOL}$ which reflects the role of drifts in the SOL transport (Section 2.2). There was, however, no explanation so far for lower density limit in conditions close to plasma detachment in reversed field configuration compared to the normal one (there is evidence from JET that reversed $B_t$ plasmas do not necessarily have larger impurity content and radiation power loss, so that earlier disruption can not be easily explained by "power starvation" of the edge plasma, and some other explanation has to be found).

5. ELECTRIC CURRENTS IN THE SOL

Experimental and theoretical aspects of divertor bias experiments have recently been reviewed by Staehler [55]. The focus here will be on the contribution of \textit{classical drifts} to the electric current. There can be two main classical contributions to the radial current: one due to the distribution of pressure and parallel convectional energy flux over the magnetic surface (these terms will be referred to as pressure-related terms), and another - due to ion-neutral collisions.

5.1 Currents due to pressure-related terms

Radial current density can be obtained by multiplying Eq.(3) by an electric charge $e$ and summing up ion and electron components. The ion and electron components of the $\mathbf{E} \times \mathbf{B}$ drift
and fluxes due to their mutual friction force $R_{ie} \equiv -R_{ei}$ will cancel each other out, and the result will be:

$$j = j_i \frac{B}{B} + \frac{1}{B^2} B \times \nabla p_\perp + \frac{(p_i - p_\perp + n m v_i^2)}{B^3} B \times B \cdot \nabla \left( \frac{B}{B} \right)$$  \hspace{1cm} (19)

As opposed to equations in Section 3, $p_i$ and $p_\perp$ are now assumed to be total (ion plus electron) pressures.

Provided the pressure is isotropic and constant along the field lines, and parallel velocity is small (which is usually the case in the plasma core), the only perpendicular current following from Eq.(19) is the current within the magnetic surface perpendicular to the magnetic field, accounted for by the second term on the r.h.s. In the torus this current, which itself is not constant along the magnetic surface due to $\nabla_r p / B \sim R = R_0 + r \sin \theta$, does not satisfy the continuity equation and must be supplemented by parallel Pfirsch-Schlüter current of magnitude (see e.g. [56], p.73):

$$j_r = 2 q \sin \theta \times \nabla_r p / B,$$  \hspace{1cm} (20)

where $q$ is the safety factor. This current consists of opposite parallel flows of ions and electrons, proportional to their radial pressure gradients. Its direction, projected onto the poloidal cross-section, is from the bottom to the top of the torus along the magnetic surface for normal, and from the top to the bottom - for reversed $B_t$ configuration. The contribution of the ion part of this current (ion parallel flow) to the overall flux pattern in the scrape-off layer of DITE (in limiter configuration) has been experimentally identified by Hugill [57] from analysis of Langmuir probe data.

On open field lines significant pressure gradients along the magnetic surface can result in non-zero surface averaged radial current. This current was first analysed by Rozhansky and Tendler [58,59] in cylindrical geometry, with neglect of the last term on the r.h.s. of Eq.(19). The main result of this theory was obtained by integrating radial current density:

$$j_r = \frac{1}{B} \frac{\partial p}{\partial \theta}$$  \hspace{1cm} (21)

over the magnetic surface from the entrance of the inner m.p.s. to the entrance of the outer m.p.s. For small values of $M_E$, the pressure asymmetry between the two sides from Eq.(18) would give for the pressure difference: $\Delta p \approx M_E p$. By writing $p = n m c_s^2$, dividing the integrated current density by the poloidal circumference $2 \pi r$ and remembering the definition of $M_E$, the averaged radial current density would be:
\[ \langle j_r \rangle \approx E_r \frac{nmce}{2\pi rBB_0} \]  

(22)

Inclusion of anomalous viscosity into the analysis leads to higher radial current required to establish the same pressure asymmetry, introducing coefficient \( K > 1 \) into the relation (22) [58,59].

Eq.(22) gives the averaged current density through the main SOL plasma, excluding contributions from magnetic pre-sheaths at the targets. They, however, can be significant [19]. Moreover, since Eq.(21) is in fact an expression for the radial component of the diamagnetic flux, it should give zero net radial current through the magnetic surface in cylindrical geometry. Recently, Chankin and Stangeby considered net radial current in toroidal geometry [60] with inclusion of both pressure dependent terms in Eq.(19) into the analysis. Their numerical results demonstrated very weak dependence of the net radial current on the \( E_r \). The direction of the current is inwards for the normal, and outwards - for the reversed \( B_t \) configuration. The origin of such a result can be understood from Fig.6. Due to the plasma sink towards the target plate, strong up-down pressure difference is formed, of the order of the pressure itself, \( \Delta p \approx p \). This generates the total current per unit toroidal length through both inboard and outboard sides of the magnetic surface, \( \approx \frac{p}{B} \). The current is directed inwards on the outboard side, and outwards - on the inboards side of the torus, as shown on Fig.6 for the case of normal \( B_t \) direction (the directions are opposite for reversed toroidal field). Due to toroidal variations of both the toroidal field \( B \) and surface area, the poloidal average of the above current should be \( \approx \epsilon \frac{p}{B} \), where \( \epsilon = r / R \) is the toroidicity. For the surface averaged radial current density one can, therefore, obtain:

\[ \langle j_r \rangle = en_c \epsilon \frac{\rho_i}{2\pi r} \]  

(23)

This estimate is quite close to numerical results of [60]. Weak dependence of the current on \( E_r \), obtained in the calculations, confirms the above interpretation that this current is caused by up-down pressure asymmetry, which is weakly affected by the poloidal \( \mathbf{E} \times \mathbf{B} \) drift. The latter influences mainly in-out asymmetries. For cylindrical geometry (\( \epsilon = 0 \)), \( \langle j_r \rangle = 0 \) was found in the calculations [60], as expected.

Estimate (23) is only valid in the main SOL. Closer to the separatrix, poloidal pressure distribution becomes more uniform and large parallel velocities are eliminated. Also, viscous forces must become important due to the large shear of the poloidal rotation near the separatrix. The flux surface average radial current must be zero at the separatrix flux surface. Radial divergence of the radial current, \( d\langle j_r \rangle / dr \), should, therefore, create two distinct regions of current flow to the target plates. Near the separatrix, current to the target should be positive for normal, and negative - for reversed \( B_t \) plasmas. Further away from the separatrix, on the
assumption of pressure exponential decay, negative parallel current should flow to the target in normal, and positive - in reversed B_t configuration. Some details of the profiles of current flow towards the target in divertor magnetic configuration are discussed in Section 6.

Pressure up-down asymmetry creates vertical electric field and, therefore, radial \( \mathbf{E} \times \mathbf{B} \) drift. This ambipolar flow cancels electron component of the pressure-related current [60], so that the current (23) is actually carried only by ions. The dependence of the direction of this current on the toroidal field direction then implies that better particle confinement in the SOL should be expected in normal compared to the reversed \( B_t \) case. This effect was proposed in [60] as an explanation for lower power threshold power for the I-H transition in normal \( B_t \) configuration (see e.g. [61,62]). By employing up-down asymmetries plus the toroidal effects of higher area and lower magnetic field on the outboard side, to get non-zero radial flows, the above explanation shares the same basis with earlier work by Hinton and Tang [63,64] on the dependence of the neoclassical ion heat transport through the separatrix on the \( B_t \) direction.

5.2 The influence of interaction with neutrals

A large number of papers have been dedicated to the influence of ion-neutral collisions on the radial current. Tsui [65] calculated the width of the poloidal velocity shear layer around the separatrix, and Lingertat et al. [66] modelled excessive heat flux conducted to the target by non-ambipolar electron flux just outside the separatrix which is caused by the divergence of the radial current. Both models used the same expression for the radial current and offer explanations for certain features of experiments on TEXT [65] and JET [66]. Radial current was obtained effectively by adding the term \( \mathbf{R}_{i-n} \times \mathbf{B} / B^2 \), which describes the damping of the poloidal rotation due to ion-neutral interactions (mainly charge-exchange), into the r.h.s. of Eq.(19). Other, pressure dependent terms, however, were ignored in both models. With the poloidal friction force given by: \( \mathbf{R}_{i-n} = -nmV_{\perp i-n}v_{i-n} \), where \( V_{\perp i} = \pm(1 / enB \cdot dp / dr \cdot E_r / B) \) ("+" for normal and "-" for reversed \( B_t \), with positive poloidal velocity assumed to be from the outer to inner divertor target, as defined in Section 3), the radial current density was obtained:

\[
J_r = nmV_{\perp i-n}v_{i-n} / eB
\]  

(24)

This current is carried only by ions and is positive, i.e. outward, for both \( B_t \) directions (the field reversal changes the sign of \( \mathbf{R}_{i-n} \) so that \( \mathbf{R}_{i-n} \times \mathbf{B} \) doesn't change). The magnitude of this current, as it will be shown in the next section, is substantially less than the one originating due to pressure-related terms. For comparison with results of biasing experiments we introduce the radial conductivity as: \( \sigma = dj_r / dE_r \). It is equal to:

\[
\sigma = nmV_{i-n} / eB^2
\]  

(25)
Considerable enhancement of radial conductivity, compared to the above expression, can be achieved when the influence of ion-neutral interactions on pressure-related terms (i.e. poloidal pressure distribution and parallel convectional energy) is taken into account. Ion-neutral interactions thus perform the role of a trigger, with pressure-related terms being responsible for the bulk of the current in the \( j_r(E_r) \) dependence. As was first demonstrated by Boozer [67] (see also later paper [68]) for the core region, when ion-neutral friction force is added to the r.h.s. of Eq.(1), it can be transformed to give the following relation between toroidal friction force and surface averaged radial current density:

\[
\frac{\partial}{\partial r}(nmV_\phi) = \langle j_r \rangle B_\theta - nm v_{i-n} V_\phi
\]  

(26)

In the regime dominated by neoclassical viscosity rather than ion-neutral friction force, toroidal velocity \( V_\phi \) should react on the applied radial electric field so as to keep the combination \( (V_\phi - E_r / B_\theta) \) constant, in the steady state conditions. Radial conductivity, therefore, is:

\[
\sigma = nm v_{i-n} / eB_\theta^2
\]  

(27)

which is larger than the one given by Eq.(25) by the factor of \( (B / B_\theta)^2 \). In the regime dominated by ion-neutral collisions, Yoshikawa [69] found that the cylindrical result (25) for the radial conductivity must be multiplied by a Pfirsch-Schlüter factor \( (1+2q^2) \) with \( q \) being the safety factor.

Certain similarities in the plasma flux pattern between the core and the SOL, such as the build-up of the toroidal rotation rather than the poloidal one as a response to the externally applied radial electrical field, pose questions as to whether above factors \( (B / B_\theta)^2 \) or \( (1+2q^2) \) of the enhancement of the radial conductivity may apply to the result (25) in the scrape-off layer. For example, local radial conductivity (away from the m.p.s. layers) which follows from the model of Weynants [70] definitely exhibits the enhancement factor \( (B / B_\theta)^2 \). The contribution of the m.p.s. to the radial current, however, has not been included in this work and would reduce it substantially. Generally, straightforward translation of the results obtained in the core, to the SOL region, is not valid due to the interaction between the SOL plasma and the targets. Pressure asymmetry between the strike zones, for example, should lead to the exchange in toroidal momentum between the plasma and the target due to opposite toroidal directions with which the plasma approaches the two strike zones.

### 5.3 Comparison with experiment

Classical contributions to radial conductivity were compared with experimental results by Lachambre et al. [71]. Experiments were performed on Tokamak de Varennes (TdeV), which is
a small machine with minor radius of 0.27m and typical toroidal field of B=1.5T. A schematic cross-section of TdeV with plasma biasing electric connections is shown on Fig 7. Radial plasma current was driven by applying voltage to neutralisation plates with respect to the grounded guard limiters. A phenomenological model based on the assumption that the ion mobility is responsible for the radial current was applied to I-V characteristics obtained in density, current and toroidal field scans.

It was found that the best fit to the experimental data from TdeV was a scaling law of the form:

$$\sigma \sim n_i^{1.84} I_p^{1.59},$$

(28)

where $n_i$ is the density at the separatrix. None of the theoretical scalings agrees with such a dependence on plasma parameters. In absolute terms, the dependence (22) falls short of the experimental conductivities by about a factor of 5 at low densities and differs even more at high densities. Since the correct expression for the net radial current, which includes m.p.s. contributions, Eq.(23), has no dependence on the radial electric field at all, pressure related terms in Eq.(19) cannot provide an explanation for the experiment.

The parametrical dependence of $\sigma$ in Eq.(28) is closest to the one given by Eq.(27). Both show strong dependence on the plasma current ($I_p \sim rB_0$), and density of neutrals, which enters the expression for the collision frequency $\nu_{i-n} = (\sigma v_{i-n}) n_n$, was roughly proportional to plasma density in TdeV. However, in absolute value the dependence (27) was found to be short by a factor of about 10 of what is required to explain the experiment, after the neutrals' toroidal acceleration due to plasma toroidal rotation was taken into account. One should also remember that Eq.(27) is likely to overestimate the neutrals' contribution to radial conductivity, as the enhancement factor $(B / B_0)^2$ has no justification in the SOL.

Overall, experiments on TdeV provide convincing evidence that at least on this machine radial conductivity has an anomalous nature. More direct experiments (biasing) are needed to extend this conclusion to larger size machines like JET, since the nature of plasma turbulence may strongly depend on the machine/plasma parameters. Alternatively, indirect experimental data such as current profiles to the target plates have to be compared with predictions of models (yet to be developed!) which correctly account for both pressure-related and neutral friction terms and their mutual influence in the SOL, or with the results of 2D numerical codes which include drifts. There is some indirect evidence (see e.g. Refs. [65,66] where comparison between model calculations and experimental data near the separatrix position has been done; also Section 6 on target current profiles) that classical contributions to radial current may be very important.
6. PROFILES OF ELECTRIC CURRENT AT THE DIVERTOR TARGET

Electrically conducting divertor target allows for local non-ambipolarity of plasma flow onto its surface. Due to the variety of physical mechanisms that may cause the current flow, fine structure of target profiles of electric current can be rather complicated and difficult to resolve in experiment. Fig.8 illustrates schematically target current density profiles for the three main mechanisms which were discussed in literature, for both normal and reversed toroidal field direction.

The thermoelectric current (dashed lines on Fig.8) has been first predicted by Harbour [24] and observed in JET [72]. This current is driven by parallel electric field caused by difference in Debye sheath drops $\approx 3T_\text{e}/e$ at the two sides of an open magnetic field line. Another source of parallel current is the electron pressure asymmetry, which can be important when temperatures at the strike zones are nearly equal (full set of equations is given in [73] for the current flow and in [74] for the heat flux to the plates). The target with lower $T_\text{e}$ and/or higher $p_\text{e}$ receives positive current from the plasma. As was found in experiments in JET [51] and JT-60U [75], the thermoelectric current is almost always parallel to the main plasma current, regardless of the $B_\varphi$ direction. This is due to higher $T_\text{e}$ at the outer side in normal $B_\varphi$ plasmas, and more equal temperature distribution plus higher pressure at the inner side, or higher $T_\text{e}$ at the inner side - in reversed $B_\varphi$ plasmas. Therefore, considering in-out asymmetries, the thermoelectric current usually flows from the outer to the inner side for normal, and from the inner to the outer side - for reversed $B_\varphi$ plasmas. Minor corrections [76] for Eq.(3) in [51] lead to the following equation for the parallel current density towards the target:

$$j_\parallel = - \frac{\sigma_\parallel T_A}{eL_\parallel} \left[ k \left( \frac{T_B}{T_A} - 1 \right) - \frac{1}{T_A} \int_A^{B} \frac{B}{B} \frac{V_\parallel}{n_\text{e}} \frac{P_\text{e}}{n_\text{e}} + \ln \left( \frac{1 + j_\parallel/j_{3A}}{1 - j_\parallel/j_{3A}} \right) \frac{T_B}{T_A} \right]$$  \hspace{1cm} (29)

In this equation side A is assumed to have lower electron temperature, $T_A < T_B$, $j_3$ is parallel ion saturation current density to the target plate, $k \equiv (k + 0.85 - \alpha) \approx 4$, $k \equiv \frac{1}{2} \frac{2m_i}{\pi m_\text{e}}$, $L_\parallel$ - connection length between the two strike zones, and inverse averaged parallel resistivity is defined as:

$$\overline{\sigma}_\parallel = \frac{e^2\lambda_\parallel L_\parallel}{m_\text{e}} \left[ \int_A^{B} \frac{dL_\parallel}{n_\text{e}} \frac{\tau_\parallel}{e} \right]^{-1}$$  \hspace{1cm} (30)

Other symbols are conventional and may be found in [73]. Around the separatrix position the connection length $L_\parallel$ approaches infinity logarithmically, and $j_\parallel \to 0$. This is schematically
reflected on Fig.8 as $|j|\$ dipping to zero at the separatrix position. Several other mechanisms to drive parallel currents in the SOL were examined in [74].

There are also two important contributions to the parallel current which arise from the divergence of the surface averaged radial current, $d\langle j_r \rangle / dr$. Radial current density due to pressure-related terms far away from the separatrix is given by Eq.(23), while the radial current density due to damping of the poloidal rotation by ion-neutral interaction with the neutrals is given by Eq.(24). It can be demonstrated that in typical conditions the former should be much larger than the latter. Employing the scaling for the poloidal/perpendicular velocity: $V_{\perp} = c_s \rho_s / \lambda_{s_{\text{sol}}}$, the current due to ion-neutral interactions can be expressed as:

$$ j_r \big|_{i_{-n}} = \varepsilon n p_s^2 V_{i-n} / \lambda_{s_{\text{sol}}} \tag{31} $$

The scaling for the ratio of the two currents is:

$$ \frac{j_r \big|_{i_{-n}}}{j_r \big|_p} = \frac{2 \pi \rho_s V_{i-n} r}{\varepsilon \lambda_{s_{\text{sol}}} c_s} \tag{32} $$

Making an upper estimate for the ion-neutral collision frequency as $V_{i-n} = c_s / \pi q R$ for conditions of strong recycling of neutrals in the scrape-off layer (SOL average for $V_{i-n}$ and $c_s$ is implied), the maximum estimate for the above ratio is $\frac{2}{q} \frac{\rho_s}{\lambda_{s_{\text{sol}}}}$. Therefore, the current due to the ion-neutral collisions can in typical conditions be neglected compared with the current originating due to pressure-related terms. The radial current caused by ion-neutral collisions is carried by ions and directed away from the separatrix both in the main SOL and in the private region (the poloidal $\mathbf{E} \times \mathbf{B}$ drift has different directions in the main SOL and the private region, as shown on Fig.1). Its direction does not depend on the direction of the toroidal field. The parallel current to the target, arising from $d\langle j_r \rangle / dr$, is shown by dotted lines on Fig.8. Negative current flows towards the target around the separatrix position, and positive current - further away from the separatrix [66].

The pressure-related current, as discussed in the previous section, results from the ion $\nabla R$ drift and the up-down asymmetry of plasma pressure. Its direction is inwards for the normal and outwards - for reversed $B_t$ direction. Near the separatrix, current to the target should be positive for normal, and negative - for reversed $B_t$ direction. Further away from the separatrix, negative parallel current should flow to the target in normal, and positive - in reversed $B_t$ direction. In divertor magnetic configuration with an X-point considerable fraction of the pressure-related current can be continued through the separatrix into the private region and be deposited onto the target just inside the separatrix. Vertical arrows on Fig.8 show the direction of the ion $\nabla B$ drift through the boundary between the divertor SOL and private regions (electron $\nabla B$ drift has an opposite direction to the ion one). It is clear that positive charge is
flowing into the private region for normal and negative - for reversed $B_I$ configuration. Due to short connection length in the private region this should create narrow peaks of the current density just inside the separatrix. The peaks should have positive sign for the normal and negative - for the reversed $B_I$ direction. The current density profile caused by pressure-related terms is schematically shown by dash-dotted lines on Fig.8. The total amount of current flowing into the private region can easily be estimated from guiding centre approximation:

$$ I \approx en \frac{2(T_e + T_i)}{eBR_o} \times (R_{out} - R_{in}) \times 2 \pi R_o \times 4 \pi p (R_{out} - R_{in}) / B $$

(33)

where $R_{in}$ and $R_{out}$ are major radii of inner and outer separatrix positions at the target.

Narrow current density peaks inside the private region with opposite signs for normal and reversed $B_I$ have been observed in JET by Schaffer et al. [7]. This paper also gives theoretical treatment of the parallel currents in fluid approximation (all parallel currents originating due to divergence of pressure-related currents are referred to as Pfirsch-Schlüter currents in [7]). Fig.9, replicated from [7], shows current at zero volts measured by Langmuir probes imbedded into the target surface obtained during the horizontal sweep of the X-point for the inner strike point (ISP) and outer strike point (OSP) regions for discharges in both normal and reversed toroidal field. Also profiles of ion saturation current density, electron pressure and electron temperature are shown. The profiles are mapped to the outer midplane and plotted as a function of $r_{mid}$, the midplane radial distance form the nominal separatrix position. Negative $r_{mid}$ correspond to private region. The position of the separatrix could not be determined with sufficiently high precision and can be slightly different from the one shown on Fig.9. Narrow current density peaks of opposite signs for the two different field configurations, superimposed on the thermoelectric current can clearly be distinguished on Fig.9.

CONCLUSIONS

Classical drifts are expected to make a significant impact on particle, energy and electric current flow pattern in the SOL and divertor. Their incorporation into 2D numerical codes is promising to greatly improve their predictive capacity. Drifts offer potential explanations for a variety of physical phenomena such as poloidal asymmetries, possible extra pinch or outward flow of plasma depending on the $B_I$ direction, non-ambipolarity of radial plasma flow and current flow towards the target.

At present, common understanding among theoreticians is almost reached on the critical issue of modified boundary conditions, and degree of sophistication of the codes is being increased by inclusion of more drift terms. Experimental results on the effects of the toroidal field reversal on divertor asymmetries should provide a proving ground for testing the codes. Present qualitative understanding is insufficient to reliably explain all experimental aspects of
these experiments. For example, the $n_e$, $T_e$, $P_{rad}$ and $H_\alpha$ asymmetries between the strike zones can be explained by the effect of the radial $E \times B$ drift, whereas power to the target asymmetries have their more straightforward explanation through the effect of the poloidal $E \times B$ drift. The codes will have to correctly treat drift particle and energy flows, as well as radiative impurity losses and the momentum loss by plasma-neutral interactions, in order to describe the experiment.

The main unresolved issues of the plasma transport related to drifts' implementation into the codes are the absence of knowledge on the contribution of anomalous radial current and its dependence on the machine/ regime parameters (will it be important in ITER, for instance?) and edge toroidal momentum which comes as a (so far, free) boundary parameter for the parallel ion velocity at the separatrix. Apart from the drifts, uncertainty in the poloidal distribution of anomalous transport coefficients, of course, still remains an impediment for increasing predictability of codes. More experimental effort and comparison between measurements and code results is needed to resolve these issues.

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Fig. 1. Schematic illustration of the main guiding centre drifts for the case of normal toroidal field ($B_t$) direction (ion $B \times \nabla B$ drift is directed towards the X-point). The direction of all the drifts is reversed for the reversed $B_t$ direction.
Fig. 2. Fractions of (a) radiation power loss, (b) heat load and (c) total loss power in the inboard divertor for normal and reversed $B_t$ directions in JT-60U as a function of $\bar{n}_e$ at safety factor $q_{eff} = 3.5$. 
Fig. 3. Tomographic reconstruction of radiative power loss for moderate density plasma ($\bar{n}_e = 1 \times 10^{20} \text{ m}^{-3}$) in Alcator C-Mod for normal (a) and reversed (b) $B_t$ direction.
Fig. 4. q-dependence of in-out asymmetries in local radiated power, $H_\alpha$ intensity and peak ion saturation current density in JET.

Fig. 5. Profiles of ion saturation current density, electron temperature and density across the target, measured by Langmuir probes in JET.
Fig. 6. Directions of radial current density caused by the pressure drop toward the target, on the inboard and outboard parts of the magnetic surface, for the case of normal $B_{L}$. For the reversed $B_{L}$ direction these directions reverse.

Fig. 7. Schematic cross-section of Tokamak de Varennes with plasma biasing electric connections.
Fig. 8. Schematic illustration of target current density profiles due to the thermoelectric current (dashed lines), radial divergence of currents caused by ion-neutral interactions (dotted lines) and pressure-related terms (dash-dotted lines) for normal (a) and reversed (b) $B_t$ directions.
Fig. 9. $j_q$ profiles for inner strike point (ISP) and outer strike point (OSP) regions for normal and reversed $B_t$ discharges in IFT. Also shown are probe-derived profiles of ion saturation parallel current density $j_{sat}^+$, electron pressure $p_e$ and electron temperature $T_e$. The profiles are mapped to the outer midplane and plotted as a function of $r_{mid}$, the midplane radial distance from the separatrix.