Turbulent particle transport in magnetized fusion plasma

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Abstract

Understanding the mechanisms responsible for particle transport is of the utmost importance for magnetized fusion plasmas. A peaked density profile is attractive to improve the fusion rate, which is proportional to the square of the density, and to self-generate a large fraction of non-inductive current required for continuous operation.

Experiments in various tokamak devices (ASDEX Upgrade, DIII-D, JET, TCV, TEXT, TFTR) indicate the existence of a turbulent particle pinch. Recently, such a turbulent pinch has been unambiguously identified in Tore Supra very long discharges, in the absence of both collisional particle pinch and central particle source, for more than 4 min (Hoang et al 2003 Phys. Rev. Lett. 90 155002). This turbulent pinch is predicted by a quasilinear theory of particle transport (Weiland J et al 1989 Nucl. Fusion 29 1810), and confirmed by nonlinear turbulence simulations (Garbet et al 2003 Phys. Rev. Lett. 91 035001) and general considerations based on the conservation of motion invariants (Baker et al 2004 Phys. Plasmas 11 992). Experimentally, the particle pinch is found to be sensitive to the magnetic field gradient in many cases (Hoang et al 2004 Phys. Rev. Lett. 93 135003, Zabolotsky et al 2003 Plasma Phys. Control. Fusion 45 735, Weisen et al 2004 Plasma Phys. Control. Fusion 46 751, Baker et al 2000 Nucl. Fusion 40 1003), to the temperature profile (Hoang et al 2004 Phys. Rev. Lett. 93 135003, Angioni et al 2004 Nucl. Fusion 44 827) and also to the collisionality that changes the nature of the microturbulence (Angioni et al 2003 Phys. Rev. Lett. 90 205003, Garzotti et al 2003 Nucl. Fusion 43 1829, Weisen et al 2004 31st EPS Conf. on Plasma Phys. (London) vol 28G (ECA) P-1.146, Lopes Cardozo N J 1995 Plasma Phys. Control. Fusion 37 799). The consistency of some of the observed dependences with the theoretical predictions gives us a clearer understanding of the particle pinch in tokamaks, allowing us to predict more accurately the density profile in ITER.

(Some figures in this article are in colour only in the electronic version)
1. Introduction

In tokamaks, the heat and particle transport occur perpendicularly to the nested toroidal magnetic flux surfaces. We refer to this direction as the radial direction and this is the unique dimension used in what follows. Cross-field heat and particle transport in fusion plasmas are only partly caused by the collisional mechanisms described by neoclassical transport theory. Indeed, the measured heat diffusivity is higher than expected from neoclassical prediction—this is what is called ‘anomalous’ transport (see, e.g. [12]) due to turbulence.

The particle transport case is somewhat different from the heat transport. Indeed, the heat source is often located in the core of the plasma making the distinction between a pinch and a diffusivity term very difficult. In contrast, the particle source is often only located in the outer edge region. In such cases peaked density profiles are nevertheless observed; they are attributed to a particle pinch. A simple way to express both the diffusive and the pinch contributions to the particle flux, \( \Gamma \), is as follows:

\[
\Gamma = -D \nabla n + V n,
\]

where \( n \) is the density profile, \( D \) the diffusion coefficient and \( V \) the pinch velocity. This formulation assumes that \( D \) and \( V \) are weak functions of \( n \) and \( \nabla n \), respectively. This paper reviews particle transport in steady-state conditions, for a review of perturbative particle transport see [13].

A well-known particle pinch mechanism is the neoclassical pinch called the Ware pinch [14]. It has been extensively compared to experimental observations. For example, in [15], the relaxation of the density profile following the injection of a deuterium pellet could be modelled with a pinch velocity of the order of the neoclassical pinch. Whereas, in [16], an additional anomalous particle pinch was needed to reproduce the perturbed experiments.

On the anomalous particle transport side, various theoretical mechanisms responsible have been identified, some due to the magnetic field gradient [17–19] others linked to the temperature gradient [2, 20]. But, experimentally, it is only lately that anomalous particle transport has been unambiguously proven in steady-state conditions. To prove its existence no neoclassical pinch and no central particle source are necessary. These two conditions have been reached in Tore Supra and TCV plasmas [1, 6] where the density profiles remained peaked. It reinforces other experimental observations in favour of the existence of an anomalous inward particle pinch [7, 8, 10, 16, 21].

The objective of this paper is to review the actual understanding of the particle pinch in tokamaks. First, we will review the theoretical predictions of this pinch and present the latest unified approach. Then, we will review the different experimental observations proving the existence of an anomalous particle pinch. The parametric dependences of this pinch found experimentally will be confronted with the theoretical predictions. Finally, the ITER density profile impact on the fusion power prediction will be discussed.

2. Theoretical studies of anomalous particle transport

The particle flux \( \Gamma \) can be divided into two parts—a part generated by the neoclassical transport \( \Gamma_{\text{neo}} \) and a part generated by the microturbulence \( \Gamma_{\text{turb}} \):

\[
\Gamma = \Gamma_{\text{neo}} + \Gamma_{\text{turb}}.
\]

The neoclassical part of the particle flux is well understood and easily modelled. Indeed, a pinch term, called the Ware pinch, has been identified in [14], where it is shown that, due to the conservation of canonical angular momentum, all trapped particles drift towards the magnetic
axis with a radial velocity $V_{\text{neo}}$, such that
\begin{equation}
\Gamma_{\text{neo}} = V_{\text{neo}} n \tag{3}
\end{equation}
with
\begin{equation}
V_{\text{neo}} = -\frac{E_{\varphi}}{B_\theta}, \tag{4}
\end{equation}
where $E_{\varphi}$ is the toroidal electric field inducing the plasma current and $B_\theta$ is the poloidal magnetic field. Other neoclassical effects, such as neoclassical thermodiffusion, also contribute to the pinch and are modelled by neoclassical codes such as NCLASS [22].

The anomalous part of the flux generated by the microturbulence can be divided into diffusive and convective terms as in equation (1):
\begin{equation}
\Gamma_{\text{turb}} = -D_{\text{turb}} \nabla n + V_{\text{turb}} n, \tag{5}
\end{equation}
where $D_{\text{turb}}$ is the diffusion coefficient and $V_{\text{turb}}$ is the velocity pinch due to microturbulence.

Various theoretical predictions for $\Gamma_{\text{turb}}$ have been made, and two main mechanisms have been identified: turbulence equi-partition (TEP) and thermodiffusion.

The first mechanism is based on the existence of Lagrangian invariants in the presence of turbulence and has a broad applicability. In the particular case of a tokamak plasma, it implies that an inhomogeneous magnetic field induces an inhomogeneous density profile [23]. Since the gradient of the magnetic field traps some of the particles in the magnetic well, the trapped particles are responsible for instabilities especially sensitive to the magnetic field inhomogeneity. This mechanism is described well in [19], where it is found that the density of trapped electrons is proportional to $1/q$. $q$ is the safety factor and is equal to the ratio between the poloidal magnetic flux and the toroidal magnetic flux. $q$ increases from the core to the edge, therefore a density proportional to $1/q$ leads to a peaked profile. When the effects of collisions and passing ions are included, the magnetic field gradient term is proportional to magnetic shear, as shown in [16, 17]. In our particular case, we will refer to the TEP mechanism as the curvature pinch, since the pinch velocity is proportional to the curvature drift [24]. The collisions can detrapsome of the trapped electrons and, therefore, weaken the role of the curvature pinch.

The second mechanism is the thermodiffusion. It predicts a velocity pinch for sufficiently high $\eta = (\nabla T/T)/(\nabla n/n)$ [20, 25–28] due to cold particles diffusing faster than hot particles. This mechanism is also observed in electron drift wave turbulence simulation in a slab geometry [29].

Other analyses based on a more global approach of the drift wave equations for electrons and ions have found that both mechanisms—curvature pinch and thermodiffusion—contribute to the particle pinch. For example, an analytical approach to the linear drift kinetic equation is presented in [4], two-dimensional simulations in [2], computed particle trajectories in [30], quasilinear approach in [9] and three-dimensional non-linear fluid model with a complementary analytical analysis in [3]. These more global approaches show that the curvature particle pinch is always directed inwards, but that the thermodiffusion term can be directed inwards when ion temperature gradient (ITG) modes are dominant and outwards when the TEM are dominant. In particular, the curvature pinch mechanism presented in [17] can be recovered in the limit of small electron pressure fluctuations, i.e. in the case where trapped electrons behave as ‘test particles’ in turbulence mainly due to ion modes [3]. These various effects are illustrated in figure 1 from [3] where the density profiles are the result of microturbulence simulations where the flux is fixed rather than the gradient. The particle flux is thus maintained at zero so that any density peaking is the signature of a turbulent pinch. If only the curvature terms are taken into account, a peaked profile is observed. Then, when the thermodiffusion mechanism is implemented, the impact on the profile varies strongly with the ratio of electron heating versus
ion heating, $S_{pe}/S_{pi}$. When $S_{pe}/S_{pi} = 1$, the density profile is similar to the curvature only case. When the ion heating dominates the electron heating, the inward pinch is reinforced and the density profile is more peaked; the ITG modes are dominant. In contrast, when the electron heating dominates, the density profile becomes hollow, the TEM dominates and the thermodiffusion term due to the TEM is stronger than the inward pinch due to the curvature.

A simple way to represent this unified theoretical view of the pinch due to drift waves in the case of no central particle source and no neoclassical pinch is to write:

$$\nabla n = \frac{V_{turb}}{D_{turb}} = -C_q \frac{\nabla q}{q} + C_T \frac{\nabla T_e}{T_e},$$

(6)

$C_q$ and $C_T$ are the respective factors weighting the two contributions to the anomalous particle pinch. In the non-linear approach illustrated in figure 1, the thermodiffusion can become the dominant pinch mechanism; this is not observed in a quasilinear approach [9], where $C_q$ is found to be always of the order of unity, whereas $C_T$ is lower except for specific density and temperature gradients. In contrast to the quasilinear approach, the non-linear simulations presented in figure 1 do not take the impact of collisions into account, impact that could weaken the outward thermodiffusion term. Another open issue concerns the role of passing electrons. In the approaches presented above the passing electrons are either not taken into account [2, 3], or are expected to have no impact on the particle pinch [4]. But recent simulation results [31] show that passing electrons lead to an inward pinch in cases where TEMs drive an outward pinch.

3. Experimental evidence for anomalous particle pinch

To conclude unambiguously on the existence of an anomalous particle pinch, two important conditions have to be reached experimentally: the particle source inside a given radius must be negligible and the neoclassical pinch due to the loop voltage inducing the current must also be negligible. As far as the particle source issue is concerned, the continuity equation is as follows:

$$\frac{\partial n}{\partial t} = -\nabla \cdot \vec{\Gamma} + S,$$

(7)

where $S$ is the particle source. Under steady-state conditions, $\partial n/\partial t = 0$, and $\int_0^r S \, dr' = \Gamma$, so if the particle source inside the radius $r$ is null, then $\Gamma = 0$. Under such conditions, $\nabla n/n = (V_{neo} + V_{turb})/D$. 

\[\text{Figure 1. Normalised density profiles versus normalised radius, } \rho, \text{ when varying the ratio of electron to ion heating } S_{pe}/S_{pi} = 0.5, 1 \text{ and 2.} \]
In tokamaks, the sources of particles have different origins. In order to identify the anomalous particle pinch, the first obvious thing to do is to avoid any central fuelling due to the use of neutral beam injection, and any gas puffing techniques, such as pellet injection, supersonic injection, etc. We also have to control the source of electrons coming from heavy ionization. When such conditions have been avoided, the only remaining particle source is due to wall out-gassing. This is due to various phenomena: recycled deuterium coming from the carbon tiles, direct ionization of molecular hydrogen, molecular dissociation and creation of Franck–Condon pairs and charge-exchange of atomic deuterium, which include molecules, atoms, charge-exchange, etc. Each of these mechanisms has a penetration length that varies from less than 1 cm to up to 10 cm. Therefore, the larger the tokamak is, the easier it is to localize well the particle source at the edge. The particle source evaluation needs Monte Carlo codes, such as [32, 33].

The neoclassical pinch issue. To prove unambiguously the existence of an anomalous particle pinch, an experiment has to be performed without neoclassical pinch, so that: $\nabla n/n = V_{\text{turb}}/D$. To reach this situation experimentally, the plasma current circulating in the torus has to be fully non-inductively driven for a time longer than the resistive diffusion time. It is only under such conditions that a zero loop voltage can be reached across the whole plasma.

These experimental conditions have been reached in two tokamaks. In TCV, the full current drive situation was obtained using fast electrons from electron cyclotron waves (ECCD) [6] for over 4 s versus a current diffusion time of 0.4 s. In Tore Supra, the full current drive provided by the fast electrons produced by lower hybrid waves (LHCD) was maintained for over 4 min, more than 80 times the current diffusion time [1]. The density profile obtained under such conditions is shown in figure 2.

Once these experimental conditions have been reached, one can deduce from the density profile peaking the ratio between $V_{\text{turb}}$ and $D$, as shown in equation (6).

In cases where the neoclassical pinch is not zero, the analysis of experimental results may lead to contradictory conclusions. In some cases, the observed pinch is found to be in agreement with $V = V_{\text{neo}}$, and $D$ proportional to the heat diffusivity ($\chi_{\text{eff}}$), as found for JET high density H-modes [11, 30, 34] and in ASDEX Upgrade for high density plasmas [35, 36]. Earlier results from a perturbative particle transport study in TFTR [12] had also shown that a neoclassical pinch could explain the observed density profile relaxation. But, in most other cases, the density profile cannot be explained by a neoclassical pinch alone. This is the case for JET L-modes [7, 11], DIII-D plasmas [8], TEXTOR RI modes [37] and of course for the TCV and Tore Supra results [1, 6]. An anomalous particle pinch was also invoked much earlier to explain hydrogen density modulation experiments in TEXT [16], where a velocity an order of magnitude above the neoclassical pinch velocity, scaling as $1/(nq)$, was needed to reproduce the observed peaked density profiles.

4. Parametric dependences of the anomalous particle pinch

After having demonstrated the existence of an anomalous particle pinch in section 2, the parametric dependences of this pinch are now compared to the theoretical predictions presented in section 1. In particular, the predicted dependence versus $\nabla q/q$, i.e. versus the current profile shape, is tested, as well as the dependence versus the temperature profiles. Finally, since different behaviours are expected depending on the role of trapped particles, the dependence versus collisionality is also tested.
The particle flux is expressed as follows:
\[
\Gamma = -D \left( \frac{\nabla n}{n} + C_q \frac{\nabla q}{q} - C_T \frac{\nabla T_e}{T_e} \right) + V_{\text{neo}}.
\] (8)

Density profile scaling as \(1/q^\eta\), \(\eta\) being between 0.5 and 1, has been widely tested. In TFTR, the density profiles are well fitted by the curvature pinch formula from [17]. In JET L-mode plasmas with LHCD [7], the peaking of the temperature profiles and of the current profiles are uncorrelated; therefore, their respective effects on the density profile can be identified. It is found that, \(n_{\text{eff}}/\langle n_e \rangle\) scales as the plasma normalized internal inductance \(l_i\), figure 3. This scaling is independent of \(\nabla T_e/T_e\) and the effective collisionality. The density profiles can be fitted using the curvature pinch formula from [17].

In TCV [6], density profiles can be described with a curvature pinch only, with \(0.4 \leq C_q \leq 1\). Most of the density profiles in Ohmic and ECH discharges, can also be reproduced with suitable combinations of the Ware pinch and the anomalous curvature pinch, assuming for instance \(C_q = 0.45\) and \(D/\chi_{\text{eff}} = 0.4\). In DIII-D also [8], most of the density

Figure 2. (a) Tore Supra density profile of shot #30428 between 6 and 20 s measured by reflectometry (blue circles). The red line is the density profile reconstructed using a very simple one-dimensional penetration model with the \(D\) and \(V\) shown in (b). The particle source has been computed by the eirene code [32]. (b) \(D\) (blue diamonds), has been taken to be equal to a value inferred for the Ohmic phase and in general agreement with the particle confinement time of 100 ms found in Tore Supra. \(V\) (red squares), matches \(D\) to reproduce the measured density profile. \(V\) is two orders of magnitude more than the neoclassical pinch velocity (green triangles).
profiles can be reproduced using an anomalous pinch term based on the adiabatic invariants approach.

There is less evidence for a correlation between the density profile and the temperature gradient, in particular because it is often difficult to decorrelate the temperature profile from the current profile. Nevertheless, a density flattening, called ‘pump-out’, is commonly observed in response to central electron heating by ECH or LH. In some cases, the flattening can be explained by neoclassical pinches. The neoclassical thermodiffusion may be reversed if the axisymmetrical configuration is lost, leading to hollow density profiles [38], as seen in TCV with strong central ECH, when a saturated (1, 1) island is present [40]. In other cases, the ‘pump-out’ seems consistent with the existence of an anomalous thermodiffusion term. This is the case for the hollow density profiles observed with strong central LH in ASDEX [41]. In ECH ASDEX Upgrade plasmas (including L- and H-modes), the nature of the unstable modes calculated with a gyrokinetic code [39] is found to be consistent with a thermodiffusion mechanism [9]. If the ITG modes are dominant, the density profile is not affected significantly by central ECH, whereas when density flattening is observed, $T_e$ increases at fixed $\nabla T_e/T_e$, and reinforces the TEM outward pinch.

In a non-linear microstability analysis of Alcator-C mod [42] density internal transport barriers, the degradation of the confinement is associated with an outward thermodiffusion particle pinch due to the onset of TEM.

In Tore Supra, plasmas with no neoclassical Ware pinch can be studied in conditions in which:

$$\frac{\nabla n}{n} = \frac{V_{\text{turb}}}{D_{\text{turb}}} = \frac{-C_q}{T_e} \frac{\nabla T}{T_e} + C_T \frac{\nabla T_e}{T_e},$$

[5]. In the core, $r/a \leq 0.30$, with $a$ being the minor radius; $\nabla n/e/n_e$ is strongly correlated to $\nabla T_e/T_e$ with $C_T > C_q$. This inward pinch is consistent with the thermodiffusion mechanism since, using a gyrokinetic code Kinezero [43], the ITG modes are found to be dominant in this region. At mid-radius, $0.35 \leq r/a \leq 0.6$; in contrast, the density peaking is found to be
mostly correlated with the magnetic shear $C_q > C_T$, as shown in figure 4(b) and found also at fixed $\nabla T_e / T_e$ with $\nabla q / q$ from 1.5 to 4.5 [5]. At mid-radius, the linear gyrokinetic simulations show that the TEM are dominant.

The collisionality is a parameter expected to have a strong impact on the density profiles, since it affects the trapped particles that are crucial for the curvature pinch [10]. The impact of collisions on TEM is characterized by $\nu_{\text{eff}}$, the ratio of the collisionality acting on trapped electrons to the vertical drift frequency of trapped electrons. In H-modes on ASDEX Upgrade [10] and on JET [34, 40], the density peaking decreases with increasing $\nu_{\text{eff}}$; figure 5 [44].

Neoclassical effects cannot explain the density peaking decrease at increasing collisionality. Indeed, a higher collisionality implies a higher resistivity and, therefore,
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A detailed analysis using quasilinear fluid drift wave modelling [10], finds the $v_{\text{eff}}$ dependence consistent with the curvature driven ITG and TEM. Nevertheless, some puzzling paradoxes are observed. Indeed, in curvature pinch cases, a $q$ profile ($l_i$ or $q_{95}$) sensitivity is expected. But in JET [12] and ASDEX Upgrade H-modes, no such dependence is observed, except for $v_{\text{eff}} < 0.3$ in JET, where the sensitivity to the $q$ profile is recovered. In contrast, in JET L-modes [7], the density peaking strongly depends on $l_i$ (figure 3) but is insensitive to $v_{\text{eff}}$. This could be explained by a $q$ profile sensitivity mainly due to trapped ions rather than trapped electrons, explaining the unseen $v_{\text{eff}}$ dependence. The opposite dependences versus the $q$ profile and $v_{\text{eff}}$ in H- and L-modes are puzzling and could be related to modified transport properties even away from the pedestal region.

5. Discussion

Recently, particle transport has made progress on both experimental and theoretical sides in a wide range of tokamaks. This has led to significant progress in our understanding of the origin of the commonly observed particle pinch.

Indeed, an anomalous contribution to this pinch has been unambiguously determined in two tokamaks [1, 6]. The two anomalous particle pinch contributions predicted by the theory, one scaling with the magnetic field gradient, the other scaling with the temperature gradient, are now understood as specific solutions of the general drift equations [2–4]. Still, the role of passing electrons and collisions have to be further studied. Nevertheless, the theoretical picture proposed in [3, 4] is consistent with several aspects of the experimental observations.

In most of the tokamak plasmas, the density profile in the gradient zone is well fitted by $1/q^n$, $n$ ranging between 0.5 and 1 [5, 7, 8, 17]. Near the magnetic axis, both the neoclassical pinch and the anomalous thermodiffusion pinch play a role [5, 9, 37]. The strength of the thermodiffusion mechanism is not clearly understood; a non-linear approach without collisions [3] predicts it to be dominant in some cases, whereas a collisional quasilinear approach [9] predicts it to be weaker than the curvature pinch as observed experimentally. When a dependence on collisionality is observed, the density peaking always decreases at higher collisionality, which is inconsistent with the Ware pinch mechanism [10, 12]. But, for similar values of $v_{\text{eff}}$, H-modes exhibit clear $v_{\text{eff}}$ dependence and L-modes exhibit no dependence at all [12]. These different behaviours with $v_{\text{eff}}$ are not yet understood.

In the light of the experimental results presented here, the currently used flat density profile for ITER predictions is rather unlikely. To illustrate the impact of a peaked density profile, scaling as $1/q^{0.5}$, ITER fusion power has been simulated using the 0-D scaling law of the code CRONOS [45, 46]. These simulations have been performed in a consistent manner using mixed empirical scaling and the global energy confinement scaling law ITERH-98P(y,2) [47]. An increase of the effective charge (from 1.55 to 1.7) depending on radiated power and on electron density [48] is included. For the ITER reference scenario with 40 MW of neutral beam heating, a fusion power of 530 MW ($Q \sim 13$) is obtained with the effect of inward curvature pinch, to be compared with 400 MW ($Q = 10$) when using the flat density profile currently expected for ITER.

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