Effects of island chains on transport through changes in the radial electric field (TJ-II stellarator)

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Abstract. The opening of island chains in magnetically confined fusion plasma is a natural and common process. The effects on transport, however, are difficult to assess because of the complexity of the phenomenon, which often happens dynamically. In this work we investigate one basic way in which transport may be altered due to the presence of magnetic islands: the modification of the ambipolar radial electric field $E_r$ due to the distortion of the flux surfaces and the presence of a separatrix. For the latter, we first adapt the metric coefficients to a new configuration that excludes the island region; then we also perturb the original $E_r$ with a neoclassical scaling taken from Shaing’s formulation [Phys. Plasmas 9 3470 (2002)]. Some examples are shown for TJ-II plasmas with dynamic configuration scans, i.e., with some low-order rational of the rotational transform moving through the plasma minor radius during one discharge. It is suggested that a modification of transport properties other than neoclassical is necessary in order to explain the experimental results.

1. Introduction

One-dimensional transport codes for magnetic fusion plasmas are based on the idea of toroidal nesting of magnetic flux surfaces all over the confined plasma. Such ideal magnetic configurations often provide a good approximation to the transport problem, but it is also known that the topology of toroidal nesting is always broken to some degree by stochastic layers or magnetic island chains [1]. These broken-topology regions can destroy transport and stability properties, as is well known, or else provoke non-destructive transport alterations with their own properties. This might be particularly interesting for stellarator plasmas where the externally imposed rotational transform avoids the most explosive tearing-mode instabilities. For example, the TJ-II stellarator has proven that low-order rational values of the rotational transform, or magnetic resonances, can be present and moved throughout the plasma without destroying confinement. Furthermore, it seems that magnetic resonances offer a good opportunity for transport control, including the formation or destruction of transport barriers associated to sheared ExB flows. See a brief summary about the TJ-II experience in [2].

Once accepted that magnetic resonances can coexist with good transport properties in helical devices, there are many ways in which they can alter confinement: from the increased containment of particles inside a static forced island, where estimated transport coefficients are found notably smaller than outside the island zone [3], to the possible use of the low order rational layers associated to the rotational transform profile as controllers of transport...
barrier formation/destruction [4]. The study of transport properties inside well defined magnetic islands (i.e., cases where the magnetic configuration develops a dominant helicity characterized by poloidal, \( m \), and toroidal, \( n \), mode numbers) in currentless devices has been done considering externally perturbed configurations, as in the LHD [5]. Now we concentrate on the regions outside the magnetic island chains (or eventual stochastic layers), where traditional 1-D transport codes can be used after appropriately modifying the metric coefficients [6], in order to study how the modification of the radial electric field according to the notions described by Shaing in a series of papers (see [7] and references therein) would affect transport.

The paper is henceforth organized as follows. In section 2 we use the formulation provided in [8, 9, 10] in order to define modified electron and ion fluxes due to the distortion of flux surfaces near magnetic islands in low-collisionality plasma regimes. These fluxes are added perturbatively to the neoclassical background fluxes. In section 3 we apply the model to TJ-II plasmas heated by ECH, where the results can be confronted with experiments of dynamical configuration change.

2. Adapting Shaing’s model

Transport studies in the TJ-II device have been mostly done with the ASTRA package [11] using fixed metric coefficients and complemented with “peripheral” modules [12] for other tasks such as the calculation of particle and heat sources or the evaluation of the radial electric field. For the latter, at each radial location we obtain the ambipolar fluxes taking to steady state the simplified equation for non intrinsically ambipolar fluxes \( \Gamma^{na}_j \),

\[
\partial_r E_r = \left| e \right| \epsilon_\perp (\Gamma^{na}_e - Z_i \Gamma^{na}_i),
\]

where \( \epsilon_\perp \) is the perpendicular dielectric constant for the plasma and \( Z_i \) is the ion charge. A contribution to these fluxes can come from different formulations of the neoclassical transport. In what follows, we use Beidler’s expressions [13] for the electron (\( j = e \)) and ion (\( j = i \)) neoclassical fluxes \( \Gamma^{NC}_j \). The modification now consists of adding locally, to each \( j \)-th species, another contribution \( \Gamma^I \) due to the presence of an island chain,

\[
\Gamma^{na}_j(E_r) = \Gamma^{NC}_j(E_r) + \Gamma^I_j.
\]  (1)

This approach is useful in many cases and has already served as a guide to interpret core-heat changes in presence of magnetic resonances in low-density TJ-II plasmas [14]. However, in this paper we want to investigate the effects due to the distortion of the flux surfaces in the vicinity of magnetic islands. A theory has been proposed and developed by Shaing [7]. It considers the modified neoclassical viscosity in a distorted tokamak magnetic configuration and provides appropriate neoclassical fluxes \( \Gamma^{na}_j = \Gamma^{NC}_j(E_r) \). Considering that the island chains may change size and location in the plasma, this formulation (which should be nonetheless developed for the particular TJ-II geometry) would require complicated calculations for each case. Instead, here we follow a perturbative approach in the sense that we take \( \Gamma^I_j \) as the additive contribution in Eq. 1 and all we ask from Shaing’s formulation is the scaling and a magnitude of the effect that is non-negligible in our plasmas. From the beginning, therefore, we warn that our implementation makes sense under the working assumptions:

(i) The distortion of flux surfaces near magnetic islands in the TJ-II stellarator provokes local changes in the neoclassical fluxes that follow the scalings provided by Shaing in [8, 9, 10].

(ii) The fluxes in presence of the distorted magnetic configuration can be treated as a neoclassical background (transport in absence of magnetic islands) perturbed by an island effect as in Eq. 1.
2.1. Island fluxes

The formulae used for the $\Gamma_j^\nu$ starting from Shaing’s expressions have been written in [15, 16]. Here we simply remind them and indicate briefly how they are obtained: we (i) substitute the helical coordinate $\Psi$ of the original formulation by a new radial coordinate $\rho$, which can be done outside the islands separatrix where the topology of nested toroids is preserved; (ii) use the tokamak minor radius $r$, which shows up as a consequence of the original tokamak geometry, in expressions for the inverse shear length $L_s^{-1} = d\ln q/dr$, the ExB drift $\omega_E = \Phi' / Br$ and a transit frequency $\omega_t = \nu_t / r$. In these definitions, $\Phi$ is the plasma electric potential, $\nu_t$ the thermal speed, $q = 1/\nu$ the inverse rotational transform and a prime means derivative with respect to the normalized coordinate $\rho$. The resulting radial fluxes show the typical collisionality, $\nu$, dependences in non-axisymmetric devices, i.e., $\propto \nu$ for the lowest collisionalities, $\propto \nu^{-1}$ for the largest collisionalities (but still in regimes where the populations of bouncing particles are important) and an intermediate regime $\propto \sqrt{\nu}$. Using Bohm coefficients $D_B = T_j/(16eB)$ we have

$$\Gamma_\nu^0 = -1.76 n_j \nu_j \omega_j^2 D_B \left( \frac{\delta_w}{\omega_E} \right)^2 \varepsilon^{-1/2} H_\nu(\rho) \left( \frac{p_j'}{p_j} + \frac{e_j \Phi'}{T_j} - \frac{T_j'}{2T_j^2} \right)$$

(2)

$$\Gamma^0_{\sqrt{\nu}} = -65.3 \frac{n_j}{\nu_j} \omega_j^2 \sqrt{\frac{\delta_w}{\omega_E}} D_B \frac{\delta_w^2 \sqrt{m}}{\langle RB_p \rangle^{1/2}} \left( \frac{r_w}{L_s} \right)^{1/2} |\omega_E|^{-3/2}$$

$$\times H_{\sqrt{\nu}}(\rho) \left[ \lambda_1 \left( \frac{p_j'}{p_j} + \frac{e_j \Phi'}{T_j} \right) - \lambda_2 T_j \right] \left( \Psi' \right)^{-1/2}$$

(3)

$$\Gamma^0_{1/\nu} = -4 \frac{n_j}{\nu_j} \omega_j^2 D_B \langle RB_p \rangle^2 \left( \frac{r_w}{L_s} \right)^{2} \langle \delta_w^2 \rangle^2 \varepsilon^{3/2}$$

$$\times H_{1/\nu}(\rho) \left( \frac{p_j'}{p_j} + \frac{e_j \Phi'}{T_j} - \frac{5T_j'}{2T_j^2} \right) \left( \Psi' \right)^{-2}$$

(4)

where new $\rho$-dependent functions $H_\nu$, $H_{\sqrt{\nu}}$ and $H_{1/\nu}$ substitute the original $G$, $J^{-1/2}$ and $H$ geometry-dependent functions defined by Shaing. Other magnitudes in Eqs. 2-4 are the major radius $R$, the particle-species density $n$, temperature $T$ and pressure $p$, the island width $r_w$, $\delta_w = r_w/R$, the poloidal mode number of the island helicity, $m$, and the aspect ratio $\varepsilon$ (its square root is representative of the fraction of toroidally trapped particles $f_t$, so we have set $\varepsilon = f_t^2$ taking $f_t$ from [17]).

It must be warned that expressions 2-4 neglect the flux surface average (on the perturbed surfaces) of the helical component of the particle fluxes, which have an angular dependency proportional to the perturbed magnetic flux associated to the opening of islands, namely $\delta \chi$, and also proportional to the mode number $m$ associated to their helicity. Therefore our “radial” fluxes require a small enough perturbation $\delta \chi \rightarrow 0$ (or the trivial limit $m \rightarrow 0$ that neglects angular deformations). Provided this, we can take $(\Psi')^{-1} \propto L_s / r_w$, the magnetic shear length over the island width, so the fluxes in the three regimes are independent of the magnetic shear but, nonetheless, always proportional to $r_w^{-2}$. Indeed, the main dependencies with island-related parameters are $\Gamma_\nu^0 \propto r_w^2$, $\Gamma^0_{\sqrt{\nu}} \propto \sqrt{\frac{m \pi}{n} r_w^2}$ and $\Gamma^0_{1/\nu} \propto \left( \frac{m \pi}{n} \right)^2 r_w^2$. Hence, only the poloidal mode number $m$ can unbalance the collisionality regimes (the $\nu$ values change little comparatively).

Recalling our working assumptions, all the unknown coefficients will be used to make a non-negligible contribution of the fluxes relative to the original neoclassical background. Thus, the $\lambda_j$ coefficients, related with kinetic integrals, will be embedded in a factor for the $\sqrt{\nu}$ collisionality regime. The $H$-functions, which also depend on the extension of the perturbation, will be taken as Gaussians centered at the resonant location in such way that they define the plasma region where the modifications of $E_r$ due to the distorted geometry are notable. Their magnitude, as
with the \( \lambda_j \), are considered in global factors for each collisionality regime.

The fluxes are obtained considering \( \nu \)-dependent global diffusivities in front of the thermodynamic gradients,

\[
\begin{align*}
D_\nu & \equiv k_\nu \nu \\
D_{\sqrt{\nu}} & \equiv k_{\sqrt{\nu}} \sqrt{\nu} \\
D_{1/\nu} & \equiv k_{1/\nu} / \nu,
\end{align*}
\]

where the variables \( k \) for each collisionality regime are deduced from the fluxes 2–4. As a first exercise, we have neglected the thermodynamic forces associated to \( T' \), so \( k_{\sqrt{\nu}} = 1 \).

\[
\omega_2 \tau_j \Omega_j D_B j \left( \delta \right) \epsilon - \frac{1}{2} H_{\nu} (\rho) \left( \Psi' \right)^{-1/2},
\]

\[
\begin{align*}
k_{\nu} & = 1.76 \frac{\omega_2}{\Omega_j} D_B j \left( \frac{\delta}{\omega_E} \right)^2 \epsilon^{-1/2} H_{\nu} (\rho) \\
k_{\sqrt{\nu}} & = 65.3 \frac{\omega_2}{\Omega_j} D_B j \delta^2 \sqrt{m} (RB p)^{1/2} \left( \frac{r_w}{L_s} \right)^{1/2} |\omega_E|^{-3/2} \lambda_1 H_{\sqrt{\nu}} (\rho) (\Psi')^{-1/2} \\
k_{1/\nu} & = 4 \frac{\omega_2}{\Omega_j} D_B j (RB p)^2 \left( \frac{r_w}{L_s} \right)^2 (m \delta_w)^2 \epsilon^{3/2} H_{1/\nu} (\rho) (\Psi')^{-2}.
\end{align*}
\]

We have used these diffusivities to construct a joint function for all the long mean-free-path range. Calling \( A \) and \( B \) the crossings between the \( \nu - \sqrt{\nu} \) and \( \sqrt{\nu} / \nu \) regimes, we find \( A = (k_{\sqrt{\nu}} / k_{\nu})^2 \) and \( B = (k_{1/\nu} / k_{\sqrt{\nu}})^{2/3} \), from which we build the joint function

\[
D(\nu) = k_{\nu} \frac{\sqrt{A \nu}}{\sqrt{A + \nu}} \frac{B^{3/2}}{\nu^{3/2}}.
\]

2.2. Calibration factors

**Table 1.** Factors to approach Shaing fluxes at different collisionality regimes for distorted tokamak geometry to the corresponding neoclassical stellarator (TJ-II) background.

| Regime: \( \nu \) | \( \sqrt{\nu} \) | \( 1/\nu \) |
|-------------------|-----------------|-------------|
| Factor:           | 10000 20 4      |             |

As mentioned above, the fluxes Eqs. 2–4 correspond to a model tokamak geometry and their magnitude is supposed to be much smaller than that of a non-optimized helical device like the TJ-II stellarator. In order to still investigate the possible effect of the flux-surface distortions due to the opening of magnetic islands in the TJ-II, we apply factors to the \( k \) coefficients in Eq. 5 (see table 1) to make the fluxes non-negligible with respect to the neoclassical background. However, we must recall that the extent of the different collisionality regimes depends on \( m \).

Since a measure of the distortion of the flux surfaces in the vicinity of magnetic islands, for a given width, comes from the spectral decomposition of the perturbed flux-surface coordinates, the factors in table 1 have been obtained for the smallest non-trivial deformation, \( m = 1 \).

The process to “calibrate” Shaing’s fluxes is better explained from figure 1. The upper panels show model density and temperature profiles for a generic ECH plasma of the TJ-II stellarator. The left bottom panel shows, for electrons (blue) and ions (green) in log-log scale, the diffusivity that corresponds to each regime (Eq. 5, dashed lines) and the one resulting from the joint formula Eq. 7 (continuous lines) using plasma parameters from the location \( \rho I \) of an...
Figure 1. Top: model density (left) and temperature (right) profiles for a generic ECH TJ-II plasma. Bottom left: diffusivities as a function of collisionality for plasma parameters at $\rho = 0.68$; modified Shaing’s values for electrons (blue) and ions (green), and Beidler’s monoenergetic coefficients for both species (red). Bottom right: electron and ion diffusivity profiles from the modified Shaing’s model (lines) and Beidler’s model after convolution (dashes). An $m = 1$ island is located at $\rho I = 0.68$ with $\delta = 1$ cm.

$m = 1, r_w = 1$ cm wide island. Also shown (in red) are the monoenergetic diffusivities obtained from our implementation of Beidler’s formulation, from which we also obtain the electric field that is imposed in Eq. [6]. The coefficients in table 1 provide a good match of the collisionality regimes, as can be appreciated. However, note that the diffusivities from Beidler’s model are obtained after convolution of the monoenergetic coefficients. The right-bottom panel shows the radial profiles that correspond to the electron and ion diffusivities from Beidler’s neoclassical background (dotted lines) and Shaing’s island-due fluxes (continuous lines). The latter are only notable around $\rho I = 0.68$. This checking has been done for different TJ-II model plasmas (e.g. higher density neutral-beam heated plasmas) and using parameters from different minor radii.

The values in table 1 have been found to be appropriate in all cases and so we keep them to proceed with the studies. The apparently enormous value of the factor for the $\nu$ regime is forced by the matching of the different collisionality regimes, which from the neoclassical point of view must justify the big $1/\nu$ fluxes that correspond to the lowest collisionalities when $E_r$ is small.

3. Transport examples with TJ-II plasmas

Part of the inspiration for the present work comes from dynamic configuration scans in the TJ-II heliac [18]. In these scans, the offset of the $\nu$-profile is changed during the discharge and, consequently, low order rationals move through minor radius. It is also known that such low-order rationals are commonly related with rotating structures with the same helicity as the expected magnetic islands [19]. In low density ECH plasmas two main transport effects are found: reduced effective electron-heat diffusivities around the resonant locations [20] confirming previous results with variable magnetic shear [21], and higher radial electric field [22] at the location of the resonance, at least in the density gradient region of these plasmas. In later
experiments it was found that, given an approximately steady-state ECH plasma, the passage of magnetic resonances through minor radius could be tracked as a local increment of the electron temperature with respect to the local average [4]. Now we check what should be expected from the formulation explained above.

### 3.1. Island-size effects

![Figure 2](image-url)

**Figure 2.** Radial electric field profiles for different effective island widths: no island (green), \( r_w = 1 \text{ cm} \) (blue) and \( r_w = 2 \text{ cm} \) (red). Left: example NBI plasma; the inset shows \( \nu^* (\rho) \) indicating the value above which the effect of the islands is to increase \( E_r \). Right: example ECH plasma with neoclassical root change around \( \rho = 0.7 \). Vertical dotted lines indicate the location \( \rho_s \) of the \( \iota = 5/3 \) value.

As mentioned after Eq. [4], the dependency of the radial fluxes due to the island effect is \( \propto r_w^2 \). Many other parameters can alter the relative importance of such fluxes, like the collisionality, so the effect may be negligible or maximized depending on plasma parameters for a given island width. In figure 2 we show TJ-II examples of NBI (generally larger collisionalities) and ECH (smaller collisionalities) plasmas indicating some main parameters. \( E_r \)-profiles for different effective island widths are distinguished by their colours. Note how the difference between the absence of islands (green) and \( r_w = 1 \text{ cm} \) (blue) becomes much larger when \( r_w \) is doubled. The diagonal transport coefficients for particle or energy transport change notably around the island location due to the change in electric field for the given collisionality. Since the effect in our plasmas generally consists of decreasing \( |E_r| \), it also happens that the transport coefficients increase around the islands. On the other hand, note that there are locations where, depending on collisionality, the effect on \( E_r \) is very small or reversed as it happens near the resonant location \( \rho_s \) in the NBI case (see inset). Finally, the ECH case at smaller collisionalities in figure 2 (right) shows that the position of the change of neoclassical root, from positive to negative \( E_r \), can shift towards smaller radii because the effect is, in this collisionality range, lowering \( |E_r| \) around the island location.

### 3.2. Dynamic configuration scans

One of the advantages of having the models for \( E_r \) included in transport codes is that it allows exploring the parameter space in more realistic conditions. For instance, the results in figure
have been obtained with fixed density and temperature profiles, but the profiles should be consistent with the sources and the transport coefficients. For this reason, we have coupled the models above to a transport code designed to mimic steady-state ECH plasmas of the TJ-II stellarator. In the calculations later shown we have further simplified the transport model by fixing the density and ion temperature profiles. Thus, only the electron temperature evolves with a thermal diffusivity that is dominantly neoclassical (see appendix for a brief indication).

Figure 3. Top: time evolution of profiles of the electron temperature over the mean in the shown time interval, \( T_e/\langle T_e \rangle \), for TJ-II discharge No. 21671 with upwards scan (left) and corresponding simulation (right) including the effect of the moving \( \iota = 5/3 \). The dashed lines indicate the separatrix in terms of the square root of the toroidal flux (values between the dashed lines are not calculated). Red dotted lines indicate \( T_e = \langle T_e \rangle \). Bottom: same for discharge No. 21658 with downwards scan.

Fig. 3 shows experimental results (left) and numerical simulations (right) for two TJ-II discharges with respective upwards (\( \iota_I \) proceeds from edge to center) and downwards (\( \iota_I \) proceeds from center to edge) magnetic configuration scans. The contour plots show the time evolution of profiles of the electron temperature over the mean in the shown time interval, \( T_e/\langle T_e \rangle \). The simulations include the effect of the moving \( \iota_I = 5/3 \) according to the modified Shaing’s effect for the assumed magnetic islands associated to this low-order rational of \( \iota \). In all cases, the dashed lines indicate the separatrix in terms of the square root of the toroidal flux (values between the dashed lines are not calculated), and red dotted lines indicate \( T_e = \langle T_e \rangle \). In the experiments, \( \iota = 5/3 \) is the lowest order but not the only rational to cross the plasma. In the simulation only the 5/3 rational is imposed with effects in the radial electric field. It is evident that the effect on the simulation is opposite to the experimental data: the average electron temperature decreases upon the passage of the island chain through minor radius.
Figure 4. Neoclassical fluxes $\Gamma/n$ for electrons (blue) and ions (green) as function of the electric field for plasma parameters of a typical ECH discharge of the TJ-II stellarator. Shown are the results from Shaing’s model (thin lines), the background neoclassical fluxes [13] (thin dashes), and their sum (thick lines) at three radial locations: $\rho \approx 0.4$ (left), $\rho \approx 0.6$ (center) and $\rho \approx 0.8$ (right).

According to the model, the radial electric field decreases around the resonant locations in these ECH plasmas. The reason of this decrement of $E_r$ is that Shaing’s model in our implementation predicts lower radial electric fields than those imposed by the background neoclassical fluxes. This is illustrated in figure 4 where the quantity $\Gamma_e/n_e$ is plotted as a function of the radial electric field at three locations for Shaing’s (thin lines) and Beidler’s (thin dashed lines) fluxes for electrons and ions. In the three radial locations the electric field is in the electron root, but it can be seen that the respective crossings between electron and ion fluxes happen at different $E_r$, giving a lower $\Gamma_e = \Gamma_i$ for their sum (thick lines). A larger island width would simply give more weight to Shaing’s fluxes making $E_r$ still closer to their crossing. A smaller radial electric field in these collisionality regimes gives larger neoclassical thermal transport, which partly explains why the island provokes smaller average temperatures in the simulations in figure 3.

We must here recall that the decreasing electric field contradicts the experimental findings, where $E_r$ was found to increase at, supposedly, the location of the magnetic resonances [22]. In these experiments, an increment $\sim 1$ kV/m was found in all cases independently of whether the background $E_r$ (in absence of low order rationals of $\iota$) was already positive or negative. Additionally, the magnitude of the change in $E_r$ did not show a clear dependence on poloidal mode number either ($m = 3$ and 4 were identified). The model is sensitive to $m$ but this might be compensated by the island width if we naturally assume that larger widths are associated to smaller $m$. In order to compare with the model, however, we must bear in mind that the geometrical effects alone can provoke transient changes in the radial electric field even disregarding the effects proposed by Shaing. For instance, if an island chain of non-negligible volume shows up in the plasma and the electric potential and electron temperature can be considered approximately constant inside the islands, there is a sudden change in $E_r$ that relaxes in a typical transport timescale. We have performed several calculations to quantify this effect. Taking for instance $\rho = 0.6$ in ECH plasmas and using effective widths $\sim 1$ cm, every time the island chain appears near the core or the edge, $E_r$ increases $\sim 1$ kV/m quite independently of the initial $E_r$ value as seen in the experiments. Similar results are obtained if the electron density also evolves in the transport model. Still, this transient increment is found almost simultaneously in most of the plasma, again at odds with the experimental result that different reflectometry channels (radial locations) are delayed among them by a time considerably larger than the confinement timescales. Therefore, we cannot invoke this mechanism alone as an explanation for the experimental findings.
4. Conclusions

We have investigated the modifications of the radial electric field due to the deformation of flux surfaces that appears following the opening of magnetic islands in plasmas like those of the TJ-II heliac. In order to make transport studies in variable configuration scans, where the islands may change size and location, Shaing’s theory [8, 9, 10] has been adapted so the dependencies are preserved and the effect is assumed to be non-negligible in the background neoclassical fields. The modification of the geometry when magnetic islands open is also taken into account [6]. Under these working assumptions, it is found that the expected effects in dynamic configuration scans of the TJ-II device are not compatible with the experimental results. We conjecture that the effects on transport of low order rationals of the rotational transform in TJ-II plasmas, particularly in low density ECH discharges, are dominated by phenomena of a different nature. For instance, a coupling of parallel and perpendicular transport might provide a better explanation because the larger speeds of electrons would explain why the experimental electric fields turn larger than the background fields in presence of magnetic resonances. A proper estimate of this effect should possibly take into account that real magnetic islands do not have a single helicity and their separatrix is more a volume than a surface.

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Appendix

Given reasonable sources and sinks for the electron channel, the electron temperature of TJ-II ECH plasmas can be acceptably reproduced using a combination like

\[ \chi_e(\rho) = \chi_{eNC}(\rho) + \chi_{eKP}(\rho) + C\rho^b, \]  

where \( \chi_{eNC} \) is a neoclassical formulation (here we have followed [13, 23]) and \( \chi_{eKP} \) is based on a dimensional scaling for trapped-electron mode turbulence given in [24], i.e., \( \chi_e \sim \gamma/k_{\perp}^2 \) for a growth rate that depends on the perpendicular wavenumber \( k_{\perp} \),

\[ \gamma = f_{tr}^2 \frac{\omega_{te}}{k_{\perp} v_e} = k_{\perp} \rho_s f_{tr}^2 \frac{\rho_s c_s}{L_n} \sqrt{\frac{m_e}{m_i}}, \]

where \( f_{tr} \) is the trapped electron fraction, \( \omega_{te} \) is the electron diamagnetic frequency and \( v_e \) the electron thermal velocity. Considering a roughly constant \( k_{\perp} \rho_s \sim 0.1 \), we write the last expression in terms of the sound speed \( c_s \) and Larmor radius \( \rho_e \), the density gradient scale length \( L_n \) and the ratio of electron to ion masses. The power function in Eq. [8] is normally used with a big exponent (we have set \( b = 10 \)) in order to lower the temperatures near the edge to experimental levels. In the calculations here presented, \( \chi_{eNC} \sim 7\chi_{eKP} \) in most of the plasma. Near the edge, \( \rho \gtrsim 0.7 \), the power function dominates.

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