

# Diffusive-convective transition for scrape-off layer transport and the heat-flux width

J R Myra, D A Russell and D A D'Ippolito

*Lodestar Research Corporation, Boulder, Colorado*

Dec. 2011

submitted to *Plasma Physics and Controlled Fusion*

---

DOE/ER/54382-66

LRC-11-145

---

**LODESTAR RESEARCH CORPORATION**

*2400 Central Avenue  
Boulder, Colorado 80301*

# Diffusive-convective transition for scrape-off layer transport and the heat-flux width

J R Myra, D A Russell and D A D'Ippolito

*Lodestar Research Corp., 2400 Central Ave. P-5, Boulder, Colorado 80301*

## Abstract

Transport of plasma from the edge pedestal gradient region into the scrape-off layer (SOL) forms the heat exhaust channel. The properties of this channel are critical for future tokamak devices. The SOL heat-flux width is believed to be set by a competition between classical parallel transport and turbulent cross-field transport. In previous work, [Myra J R, *et al.* 2011 J. Nucl. Mat. **415** S605] focusing on modeling of the heat flux width in the National Spherical Torus Experiment (NSTX), the possibility of a transition from quasi-diffusive to convective transport in the SOL was noticed. This transition, and the scaling of the heat-flux width is explored here through additional SOL turbulence simulations using the SOLT code [Russell D A, *et al.* 2009 Phys. Plasmas **16** 122304]. At the transition, the transport becomes intermittent, and the SOL width is broadened due to blob emission. Critical parameters for the transition are investigated, including the power flux into the SOL, the field line pitch, the connection length and the plasma collisionality. An inverse dependence of the heat-flux width on the poloidal field, also seen routinely in experiments, is noted and explained qualitatively.

PACS: 52.35.Ra, 52.65.-y, 52.55.Fa

## 1. Introduction

In magnetically confined plasmas, the hot dense core region must transition to a colder plasma before interacting with material surfaces. This transition takes place in the edge and scrape-off-layer (SOL) regions. The SOL is the region of open field lines where plasma actually makes contact with material surfaces. This region is the main one of interest for the present study. The edge is a somewhat poorly defined region which in present terminology will refer to the closed field line region just inside the separatrix or last closed flux surface (LCFS). The difficulty in defining the edge region precisely reflects the coupling that exists between the core and the SOL, and the present lack of physics understanding of how this coupling works. From the point of view of SOL studies, the edge region is important as the source of SOL plasma, i.e. particles, energy and momentum transport across the LCFS to feed the SOL. Conversely, the plasma exhaust that leaves the tokamak through the SOL affects the density, temperature and flow profiles in the edge region.

The SOL and edge regions affect overall confinement of tokamak plasmas. Edge transport barriers give rise to enhanced confinement modes such as the high-confinement (H) -mode [1]. The transition from low (L) -mode to H-mode seems to be related to mean and/or oscillating (zonal) flows in the edge [2,3]. Further evidence for the importance of the edge and SOL plasmas to tokamak operation comes from the beneficial effects on confinement resulting from lithium-coated plasma-facing surfaces [4,5]. It is very important to understand the physics of the edge and SOL plasmas, both from the point of view of understanding core confinement, and mitigating potentially damaging interactions of plasma exhaust with material surfaces [6].

In the SOL itself, the simplest conceptual model is to regard the SOL plasma profiles as resulting from a competition between parallel and cross-field transport. This competition sets a characteristic SOL width [7,8]. While parallel transport is usually considered to be classical [9], often with flux limits and other corrections [10], turbulent cross-field transport is still more complex. In the far SOL, cross-field transport can be dominated by blob-filaments (often referred to as “blobs”) which are described in a number of review articles [11,12,13]. The motion of blobs is understood as being primarily convective. The physics of the near SOL (close to the LCFS) is more complicated. Here diffusive and convective processes can compete and the environment

is typically more turbulent [14]. There is a strong need for reliable predictive models and corresponding experimental validation of plasma behavior in the near SOL region.

The interactions of plasma exhaust with the surface of the divertor are a key motivation for understanding the near SOL. This topic has motivated a concerted effort on major tokamaks in the USA to measure the scaling of the SOL heat-flux width  $\lambda_q$  and peak power fluxes on the divertor [15,16,17]. Similar studies have also been carried out in Europe. [18] The present paper extends previous modeling work in this area. We begin, as a base case, with plasma conditions from the National Spherical Torus Experiment (NSTX) [19] for a case that was previously modeled [20,21]. Starting from this base case, a number of numerical scaling experiments are carried out by varying code input parameters systematically one at a time. In addition to studying the dependence of  $\lambda_q$  on net power flux into the SOL, the field line pitch, the connection length and the plasma collisionality, we observe a transition in SOL dynamics from a nearly laminar, diffusive-like regime, to a more turbulent convective transport regime. The properties of this diffusive-convective transition are a major focus of our paper.

The present work is related to a few similar numerical studies in the literature. In particular, in a previous investigation [22] different confinement and transport states were identified in a two-dimensional (2D) resistive-g model depending on the strength of the turbulence drive. In Ref. 23 low and high confinement states of a simple magnetized torus were studied. The 2D model employed in these studies is similar to the model presented here. In another earlier approach, a simple one-dimensional model which allowed for competition between mode coupling and profile relaxation was shown to exhibit crossover from diffusive to ballistic turbulent transport governed by a single control parameter [24].

The formation of large-scale structures (including blobs [13], streamers and zonal flows [25]) from turbulence is a topic of considerable contemporary interest [26]. The edge and SOL parameter space is a very large one, and many different types of behaviors are observed and modeled. For example, in contrast to the present study in an NSTX-like parameter regime, similar modeling efforts [27] of the heat flux width and edge-SOL transport in the Alcator C-Mod tokamak [17] produced an entirely different phenomenon: quasi-coherent oscillations that appear to be a key controlling element of separatrix-crossing transport and blob formation in that case.

In Sec. II the physics content of our numerical model and the procedure for our study are discussed. Results for the heat-flux width ( $\lambda_q$ ) scaling and the diffusive-

convective transition phenomena are given in Sec. III. Conclusions are presented in Sec. IV.

## 2. Numerical model

### 2.1. Physics content

The physics model for this study is a three-field fluid turbulence model evolving density  $n$ , electron temperature  $T_e \equiv T$ , and electrostatic potential  $\Phi$ . A simplified approach to the geometry is employed: turbulent structures are assumed to be constant filaments elongated in the direction of the background magnetic field  $\mathbf{B}$ . Thus, the dynamics is contained in the two dimensions perpendicular to  $\mathbf{B}$ . For tokamak plasmas, this plane is considered to be at the outboard midplane of the torus.

Computations are carried out with the 2D Scrape-Off-Layer Turbulence (SOLT) code [28]. Analytical closure relations are employed to describe the parallel dynamics. For example, the relationship of parallel current  $J_{\parallel}$  to  $\Phi$  is prescribed by a sheath closure on open field lines (generalized to allow for collision-dominated regimes), and a Wakatani-Hasegawa-like model [29] for  $J_{\parallel}$  is employed on closed field lines to capture basic drift-wave physics.

The fundamental equations for the SOLT model are the vorticity, continuity and energy conservation (or alternatively temperature) equations

$$\frac{d}{dt} \nabla \cdot \mathbf{n} \nabla \Phi = A_{\text{dw}}(\Phi, n) + \alpha_{\text{sh}} J_{\parallel} - \beta \frac{\partial(nT)}{\partial y} \quad (1)$$

$$\frac{dn}{dt} = A_{\text{dw}}(\Phi, n) - \alpha_{\text{sh}} \left( nT^{1/2} - J_{\parallel} \right) \quad (2)$$

$$\frac{dT}{dt} = -\alpha_{\text{sh}} \frac{q_{\parallel}}{n} \quad (3)$$

Here we employ dimensionless notation (using the Bohm normalization with time-scale  $\Omega_{\text{ci}} = ZeB/m_i c$  and space scale  $\rho_{\text{sr}} = c_{\text{sr}}/\Omega_{\text{ci}}$  where  $c_{\text{sr}}^2 = T_{\text{er}}/m_i$  and  $T_{\text{er}}$  is a reference temperature for the normalization). The simulation plane is denoted as the  $(x, y)$  plane where  $x$  is the radial direction and  $y$  is binormal (approximately poloidal). The radial variables  $r$  and  $x$  are used interchangeably in this paper; the only difference is that the origin of  $x$  is taken to be the LCFS. The convective derivative is written as  $d/dt = \partial/\partial t + \mathbf{v} \cdot \nabla$  where  $\mathbf{v} = \mathbf{e}_z \times \nabla \Phi$ .  $A_{\text{dw}}$  is the drift wave operator

$$A_{dw}(\Phi, n) = \alpha_{dw} \bar{T}^{3/2} \{\Phi - T \ln n\} \quad (4)$$

where for any quantity  $Q$ ,  $\langle Q \rangle \equiv \bar{Q}$  denotes the zonal or  $y$  average part and  $\tilde{Q} \equiv \{Q\} \equiv Q - \bar{Q}$  denotes the fluctuating part. Thus  $A_{dw}$  enforces a Boltzmann response on fluctuations when the coefficient  $\alpha_{dw}$  is large, in the spirit of the Wakatani-Hasegawa adiabaticity parameter. Note that the flux-surface or  $y$ -average of  $A_{dw}$  vanishes (as it must since it arises from  $\nabla_{\parallel} J_{\parallel}$  on closed surfaces).

To close the system of equations, the parallel current and heat flux  $J_{\parallel}$ ,  $q_{\parallel}$  must be expressed in terms of the dynamic variables  $\Phi$ ,  $n$  and  $T$ . A set of closure relations, valid for a range of collisionality regimes from conduction limited to sheath-connected, is employed here and discussed in detail in Appendix A of Ref 21. In particular, the parallel heat flux in the sheath-connected limit takes the form

$$q_{\parallel SL} = s_E n T^{3/2} e^{(\Phi_B - \Phi)/T} \quad (5)$$

where  $\Phi_B \sim 3 T$ , and  $s_E \sim 6$  is the sheath energy transmission factor. In the collisional (conduction limited) case we have

$$q_{\parallel CL} = \frac{3.2 n T^{3/2}}{\Lambda} \quad (6)$$

where the collisionality parameter  $\Lambda$  is defined below. A smooth interpolation between regimes is employed for intermediate cases.

The most important input parameters of the model for the present application are the sheath conductivity parameter

$$\alpha_{sh} = \frac{\rho_{sr}}{L_{\parallel}} \quad (7)$$

where  $L_{\parallel}$  is the midplane-to-divertor-plate connection length in the SOL, the curvature drive parameter

$$\beta = \frac{2\rho_{sr}}{R} \quad (8)$$

where  $R$  is the effective radius of curvature (approximately the major radius of the torus), and the collisionality parameter

$$\Lambda = \frac{v_{ei} L_{\parallel}}{\Omega_e \rho_s} \quad (9)$$

which enters the closure relations for  $J_{\parallel}$  and  $q_{\parallel}$ .  $\Lambda$  is related to the SOL electron parameter  $v_{*e}$  introduced by some authors as  $v_{*e} = L_{\parallel} / \lambda_{ei} = \Lambda(m_i / m_e)^{1/2}$ , where  $\lambda_{ei}$  is the electron mean free path for collisions with ions.

The SOLT code solves the vorticity equation, by splitting Eq. (1) into zonally-averaged and fluctuating parts. The zonally-averaged part is manipulated into a conservative form for zonal (i.e.  $y$ ) momentum conservation. The Boussinesq approximation is employed on the fluctuating part. Small diffusive terms are usually added to Eqs. (1) – (3) to absorb high-wavenumber fluctuations before they cascade to the scale of the numerical grid. Further discussion of the SOLT model is given in previous publications [32].

## 2.2. Procedure

In the present work, our goal is to carry out a SOL simulation with as few artificial features as possible. The role of the edge region is to provide an effective boundary condition at the LCFS. Artificial source terms are added to Eqs. (2) and (3) to maintain the time- and  $y$ -averaged  $n$  and  $T$  profiles *inside the LCFS* against losses to and in the SOL. These source functions are set to zero for  $x > x_s \approx -1$  cm so that the SOL is fueled entirely by transport across the LCFS, i.e. the SOL profiles themselves are determined self-consistently by the balance between perpendicular turbulent transport and parallel losses.

Following closely the procedure used in earlier work [21,27], we impose a mean sheared flow velocity *in the edge region* to simulate the  $E_r$  well that is characteristic of H-mode plasmas

$$v_y = -E_x = \frac{\tau}{n} \frac{d}{dx} n T_e \quad (10)$$

Setting  $\tau = T_i/T_e$  forces balancing of the  $\mathbf{E} \times \mathbf{B}$  and ion diamagnetic flow, which approximates the radial force balance characteristic of H-mode pedestals where the relatively flat core profiles transition sharply to their much smaller edge values. Here we regard  $\tau$  as an order unity control parameter that regulates turbulence. Varying  $\tau$  changes the power crossing the separatrix  $P_{sep}$ . In this way, we effectively achieve a heat flux-driven boundary condition for our SOL simulations. Results will be presented in this paper regarding  $P_{sep}$  as the independent variable.

It is important to recognize that  $P_{sep}$  (i.e.  $\tau$ ) scans presented in this paper *hold the target edge profiles fixed*. (See Fig. 1.) Thus, qualitatively a  $P_{sep}$  scan can be regarded as

a confinement scan: low  $P_{\text{sep}}$  (heat flux) and steep plasma profiles typical of H-mode imply good confinement, while high  $P_{\text{sep}}$  for the same profiles implies poorer confinement.

Our analysis begins with the plasma profiles and machine parameters of a particular discharge on NSTX that was modeled previously [21]. These parameters put our base case in a low collisionality sheath-connected regime. We then vary parameters about this real NSTX base case (keeping mean edge profiles for  $x < x_s$  fixed) to explore the parameter space of our numerical model. In particular we carry out  $\tau$  scans for various values of  $b_\theta = B_\theta/B$  (which specifies the field line pitch), SOL connection length  $L_{\parallel}$ , and SOL collisionality  $\Lambda$ . For each run we calculate

$$P_{\text{sep}} = 2\pi R b_\theta \int dr q_{\parallel} \quad (11)$$

where the  $q_{\parallel}$  integral is over the entire SOL. The reference plasma parameters for the base case are  $n_{\text{er}} = 1.9 \times 10^{13} \text{ cm}^{-3}$ ,  $T_{\text{er}} = 125 \text{ eV}$ ,  $B_r = 0.27 \text{ T}$  (outboard midplane) yielding  $\rho_{\text{sr}} = 0.60 \text{ cm}$ ,  $\Omega_{\text{ir}} = 1.29 \times 10^7/\text{s}$ ,  $c_{\text{sr}} = 7.74 \times 10^6 \text{ cm/s}$ . The reference values of density and temperature are taken at the inflection point of the tanh-function fits. (See Fig. 1). Other inputs are  $R = 150 \text{ cm}$ , deuterium plasma (charge  $Z = 1$ , mass  $A = 2$ ), and  $b_\theta = 0.3$  (outboard midplane). Dimensionless SOLT input parameters are  $\beta = 8.0 \times 10^{-3}$ ,  $\mu = 0.1$  (dimensionless viscosity),  $D = 0.01$  (dimensionless density diffusion). The sheath conductivity is  $\alpha_{\text{sh}} = \rho_{\text{sr}}/(a_f + b_f \text{Log}[x(\text{cm})])$  where the connection length  $L_{\parallel}$  in the denominator was fit from an equilibrium reconstruction with  $a_f = 750 \text{ cm}$  and  $b_f = -225 \text{ cm}$ . The adiabaticity parameter on the core-side boundary is  $\alpha_{\text{dw}} = 0.172$ , estimated from NSTX parameters, and its tanh-shaped profile goes to zero at the LCFS. Experimentally, the value  $P_{\text{sep}} = 0.8 \text{ MW}$  was observed for this discharge.

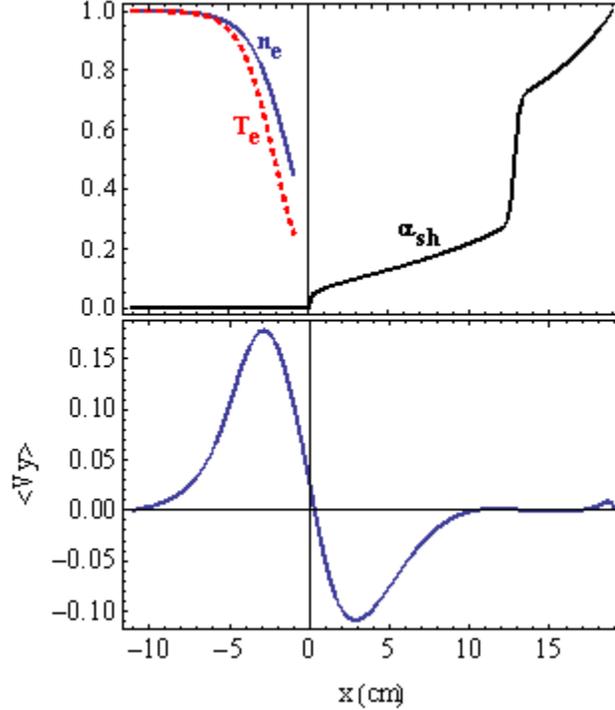


Fig. 1 (upper panel) Target plasma profiles in the edge, imposed for  $x < -1$  cm, and radial variation of the sheath conductivity parameter  $\alpha_{sh}$  in the SOL,  $x > 0$ . Here plasma profiles are normalized to their value at the left boundary and  $\alpha_{sh}$  is normalized to its value at the right boundary. The sharp rise in  $\alpha_{sh}$  at  $x = 12$  is due to an imposed limiter. (lower panel) Typical radial variation of the mean binormal flow velocity  $\langle v_y \rangle$ , imposed for  $x < -1$  cm, but determined dynamically elsewhere. Noting that  $E_r \sim -\langle v_y \rangle$  in our model, a typical  $E_r$  well exists just inside the separatrix.

### 3. Results

Experiments have revealed a consistent inverse dependence of the heat-flux width on the plasma current  $I_p \propto B_\theta$  and relatively less dependence on other parameters for a given tokamak. This observation motivates a theoretical study of the heat flux width scaling, and the choice of parameter scans studied in the next subsection, III A. It will be seen that while qualitative scalings pertain, there is no one power-law scaling that describes all regimes. Rather, we find the interesting result that there is a break in scaling that separates two regimes for the SOL, a somewhat laminar diffusive-like regime, and a turbulent blob-dominated regime. These regimes are investigated further in Sec. III B.

### 3.1. Heat-flux width scaling

In the following, results are presented for the heat-flux width in  $P_{\text{sep}}$  scans using the base case parameters and then for separate scans in which  $b_\theta$ ,  $L_\parallel$  and  $\Lambda$ , are modified by specified amounts. For these heat-flux width scaling studies, we employ an integral definition [30] of  $\lambda_q$ , based on Eq. (9)

$$\lambda_q = \frac{P_{\text{sep}}}{2\pi R b_\theta q_{\parallel\text{peak}}} \quad (12)$$

where  $q_{\parallel\text{peak}}$  is the peak value of the heat flux, obtained at the first grid-point in the SOL.

Fig. 2 shows the results of a study in which  $b_\theta$  is varied, with the base case denoted by  $b_{\theta 0}$ . The parameter  $b_\theta$  enters the simulation inputs in two ways: (i) it appears explicitly in Eq. (11) for  $P_{\text{sep}}$ , and (ii) it affects the connection length  $L_\parallel$  which, for a given flux surface shape, changes with the field line pitch.

As illustrated in Fig. 2, there is a consistent inverse dependence of  $\lambda_q$  on  $b_\theta$  which for low powers can be approximated as  $\lambda_q \propto 1/b_\theta^m$  with  $m \sim 0.5$  at low power but closer to 1 at high power. This inverse dependence is qualitatively similar to the inverse dependence of  $\lambda_q$  on the plasma current  $I_p$  that has been noted in several experiments. [15,16,17] Our previous modeling work [21] on specific experimental shot sequences also showed this trend.

We can understand this dependence from Eq. (11). If  $b_\theta$  is decreased at constant  $P_{\text{sep}}$ , then the integral of  $q_\parallel$  must rise. From Eq. (5) this can only happen if density and temperature at the separatrix rise and/or if the SOL width increases. Both effects occur, and both require stronger levels of turbulence carrying plasma from the edge region into the SOL.

It is interesting that at  $P_{\text{sep}} \sim 2$  MW there is a break from the low power  $\lambda_q \propto 1/b_\theta^{1/2}$  scaling. This break in scaling turns out to be related to the diffusive-convective transition that is the second subject of the present paper.

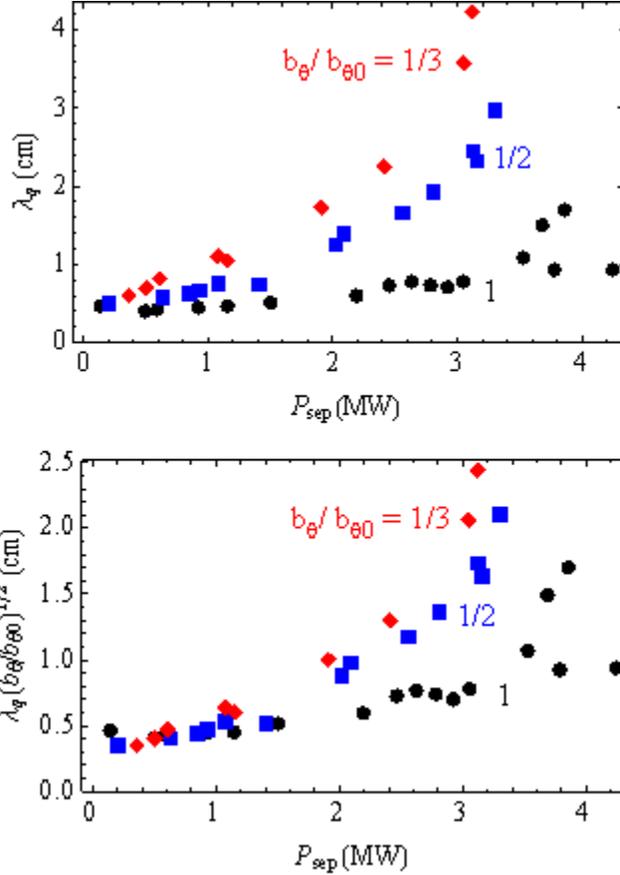


Fig. 2 Upper panel: power variation of the heat flux width for various field line angles: base case with  $b_\theta \equiv b_{\theta 0} = 0.3$  (black disks),  $b_\theta/b_{\theta 0} = 1/2$  (blue squares) and  $b_\theta/b_{\theta 0} = 1/3$  (red diamonds). There is a consistent inverse dependence on  $b_\theta$ . Lower panel: same data with  $(b_{\theta 0}/b_\theta)^{1/2}$  scaling removed. At low powers  $P < 2$  MW, this scaling unifies the results, however at high powers the  $b_\theta$  dependence is stronger.

Although the connection length  $L_{||}$  changes when  $b_\theta$  is changed, further numerical experimentation reveals that  $L_{||}$  itself, is not the main contributor to the  $b_\theta$  scaling of Fig. 2. In Fig. 3 the connection length was artificially varied by changing  $\alpha_{sh}$  while holding  $b_\theta$  fixed. Little explicit dependence resulted. This result is qualitatively consistent with observations on Alcator C-Mod that showed  $\lambda_q$  to be insensitive to  $L_{||}$  in experiments that compared single and double null SOL geometries. This result is not unexpected from Eqs. (5) and (11), which have no explicit  $L_{||}$  dependence in the sheath-connected regime which characterizes the base case. The connection length does enter  $\Lambda$  in the conduction limited regime of Eq. (6). This motivated the next numerical experiment.

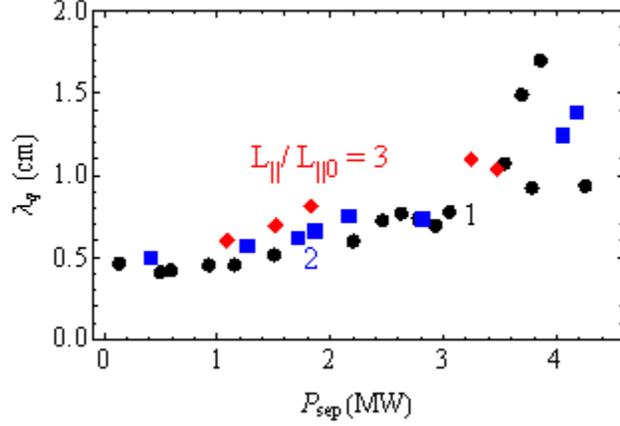


Fig. 3 Power variation of the heat flux width for various connection lengths: base case (black disks),  $L_{||} / L_{||0} = 2$  (blue squares) and  $L_{||} / L_{||0} = 3$  (red diamonds). There is little explicit dependence on the connection length.

In Fig. 4,  $b_{\theta}$  and  $L_{||}$  (i.e.  $\alpha_{sh}$ ) were held fixed, but the collisionality parameter  $\Lambda$  was artificially multiplied by a constant factor in one of the scans. A weak increase of  $\lambda_q$  with  $\Lambda$  is evident. The dependence is stronger at higher  $P_{sep}$ , as was also the case in the  $b_{\theta}$  study of Fig. 2. An increase in the SOL density width with collisionality for blob-driven transport has been previously noted in experiments [31] and simulations [31,32] where it is related to an increase of blob speed [33]. The collisionality dependence of  $\lambda_q$  seen here can be understood qualitatively by an argument similar to that given for the  $b_{\theta}$  dependence. From Eqs. (6) and (11), if  $\Lambda$  is increased at constant  $P_{sep}$  then either  $n$  and  $T$  must go up at the separatrix, and/or the SOL must broaden to compensate. Both effects occur, and both require stronger levels of turbulence carrying plasma from the edge region into the SOL.

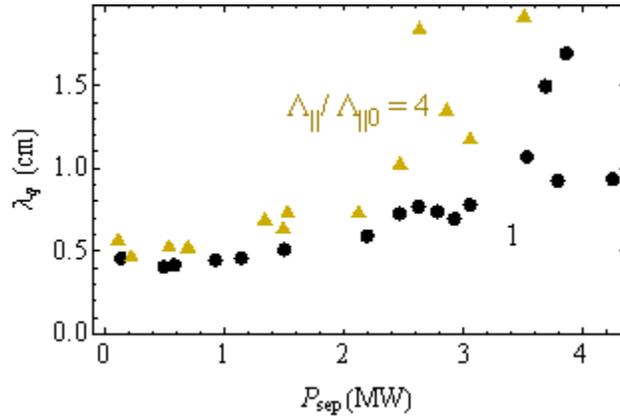


Fig. 4 Power variation of the heat flux width for various collisionalities: base case (black disks),  $\Lambda / \Lambda_0 = 4$  (gold triangles). There is a weak, but consistent increase of  $\lambda_q$  with collisionality.

### 3.2. Diffusive-convective transition

As evidenced from the scans in the preceding section, at a given  $P_{\text{sep}}$  the SOL width broadens with decreasing  $b_\theta$ , and increasing collisionality  $\Lambda$ . The mechanism for the broadening of the SOL width is increased cross-field turbulent transport. In analogy with critical Reynolds number phenomena in neutral fluid turbulence, the turbulence and cross-field transport in the present system appear to undergo a qualitative change from a nearly laminar diffusive-like regime to a turbulent convective regime which tends to be dominated by blob ejection.

Figure 5 shows plasma profile comparisons in the SOL contrasting the low power quasi-diffusive regime and the high power convective regime. The cases correspond to  $P_{\text{sep}} = 0.6$  and 3.9 MW respectively, with  $b_\theta/b_{\theta 0} = 1/2$ ,  $\bar{\mu} = 0.01$  and  $D = 0$ . We choose reduced ad-hoc diffusion and viscosity parameters here and for Fig. 6 so as not to bias the characterization of the diffusive regime. (Smaller time steps are required with these choices, so they are not employed for the broad parameter studies of the previous section.) The low power regime is not strictly diffusive because there is no separation of

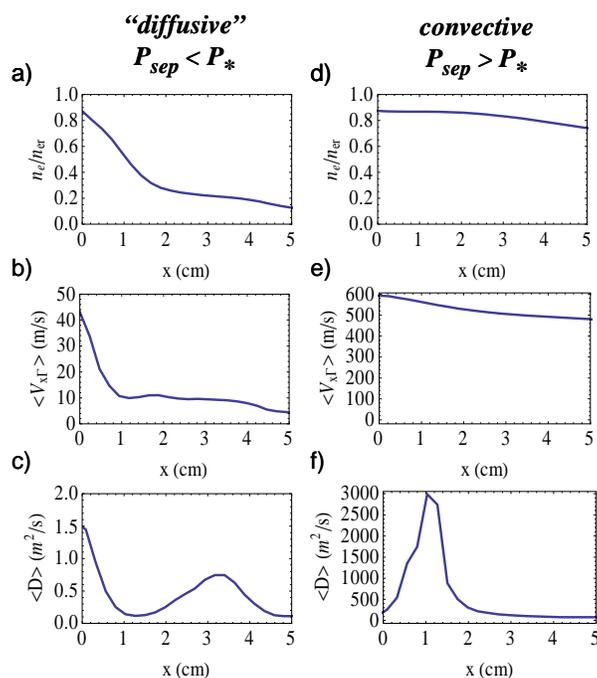


Fig. 5 Profile comparisons in the SOL contrasting the low power quasi-diffusive regime (left column) and the high power convective regime (right column). Shown are density profiles in a) and d), effective radial convection velocity in b) and e), and effective diffusion coefficient in c) and f). Note the flattening of the density and velocity profiles in the convective regime, d) and e).

scales (i.e. there is no step size that is small-scale with respect to the profiles). Nevertheless, an effective diffusion coefficient can be defined as  $D_{\text{eff}} = \Gamma / \nabla \langle n \rangle$  where  $\Gamma = \langle n v_x \rangle$  is the particle flux. The fact that  $D_{\text{eff}}$  varies on a scale similar to that of  $\langle n \rangle$  itself suggests the absence of a truly diffusive process. It is also possible to define an effective radial convection velocity  $V_{\text{eff}} = \Gamma / \langle n \rangle$ . In the high power regime,  $P > P_*$  (where  $P_* \sim 2$  MW from Figs. 2 and 4),  $\langle n \rangle$  and  $V_{\text{eff}}$  are noticeably flattened, and  $V_{\text{eff}}$  appears to be a better descriptor of transport than the wildly varying  $D_{\text{eff}}$ .

Further evidence for a change of character in the transport comes from examining the probability distribution functions (PDFs) of density and flux variations shown in Fig. 6 for the same two extreme cases. Note also the transition from nearly laminar to strongly turbulent dynamics in going from the low power “diffusive” case to the high power convective case.

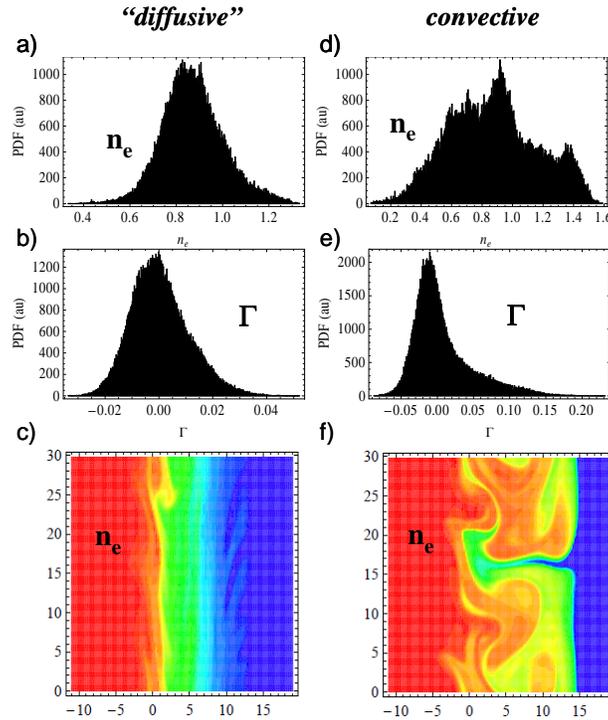


Fig. 6 PDFs contrasting the low power “diffusive” regime and the high power convective regime. Show are the PDFs of density in a) and d), and particle flux in b) and e) The scale of the flux in e) is about 5 times larger than in b) due to the much faster transport. Two-dimensional snapshots in the  $(x, y)$  plane of the density (color palette, log scale) in parts c) and f) illustrate the transition from nearly laminar to strongly turbulent dynamics. The high density core plasma is to the left.

The presence of convective transport and strong PDF tails for the high power case shown in Figs. 5 and 6 is usually indicative of blob formation and ejection, and also manifests itself in the time history of synthetic probe diagnostics in SOLT. Time traces of the density in the SOL are shown in Fig. 7a) comparing a low power ( $P_{\text{sep}} = 1.1$  MW) and high power ( $P_{\text{sep}} = 3.3$  MW) case, from the  $b_{\theta}/b_{\theta 0} = 1/2$  scan of Fig. 2. The bursts seen in the high power case are usually accompanied by blob ejection. A particularly good example of this is shown in Fig. 7b).

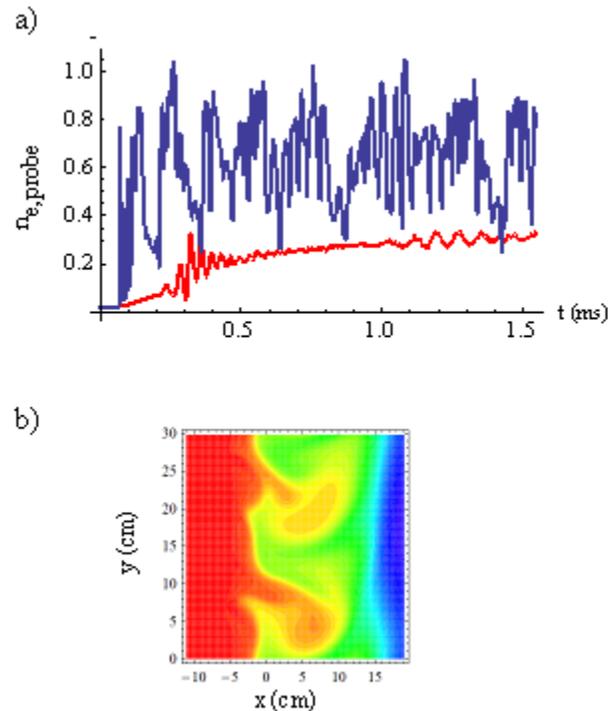


Fig. 7 a) Time histories of the fluctuating density in the SOL for a high power turbulent convective case (upper blue trace), and a low power nearly laminar diffusive case (lower red trace). The bursty character of the signal is evident in the convective case. The bursts are often associated with the emission of one or more blobs. b) Two-dimensional snapshot in the  $(x, y)$  plane of the density (color palette, log scale) at a time corresponding to a strong burst and emission of two blobs in the turbulent convective case. The high density core plasma is to the left.

So far we have seen evidence for a diffusive-convective transition from rather extreme cases, chosen to highlight the differences. In fact, for most of the runs, the transition while qualitative, is rather easily identified by visual inspection of the  $V_{\text{eff}}$  profiles and snap-shots of the dynamics. Before attempting to identify more

quantitatively the critical parameters for the transition, it is useful to examine heuristic scalings for diffusive and convective transport.

For turbulent diffusive transport, we expect a mixing-length scaling to hold, where we can estimate  $D_{\text{eff}} \sim \gamma/k_y^2$  with  $\gamma$  a characteristic growth rate of instabilities sustained by the time-averaged profiles, and  $k_y$  a characteristic wavenumber. A characteristic velocity in the diffusive regime is therefore  $D_{\text{eff}}/\lambda_q$  where we take  $\lambda_q$  as the typical profile scale.

For blob-dominated convective transport, we can estimate the characteristic net convective velocity as  $v_b f_p$  where  $v_b$  is the radial velocity of an individual blob and  $f_p$  is the blob packing fraction (the percentage of space occupied by blobs). Furthermore, using the blob correspondence principle [11,13] we estimate  $v_b \sim \gamma/k_y$ . Thus

$$V_{\text{eff}} \sim \begin{cases} D_{\text{eff}} \nabla \ln q_{\parallel} \sim \frac{\gamma}{k_y^2 \lambda_q} & \text{diffusive} \\ v_b f_p \sim \frac{\gamma}{k_y} f_p & \text{convective} \end{cases} \quad (13)$$

Comparing the estimates for  $V_{\text{eff}}$  with diffusive and convective transport mechanisms we find the ratio

$$\text{diffusion} : \text{convection} = \frac{1}{k_y \lambda_q} : f_p \quad (14)$$

This leads to the following scenario. At sufficiently low  $P_{\text{sep}}$  where  $\lambda_q$  and hence  $k_y \lambda_q$  is small, diffusion dominates convection. As  $P_{\text{sep}}$  is raised, we find numerically that both  $k_y$  and  $\lambda_q$  increase; thus there is a decrease in  $1/(k_y \lambda_q)$ . For sufficiently large  $\lambda_q$ , i.e. with flattened profiles, diffusion cannot supply enough flux to deliver the power through the SOL. At this point the system responds by starting to eject more and more blobs, i.e. increasing the packing fraction  $f_p$ . Therefore, we expect the diffusive-convective transition to occur for  $k_y \lambda_q$  of order unity.

Figure 8 combines the data for four scans. Three of the scans are the same  $b_\theta$  cases considered in Fig. 2. The fourth case is an extreme collisionality case ( $b_\theta/b_{\theta 0} = 1/2$  and  $\Lambda/\Lambda_0 = 16$ ) chosen to provide data in the small  $P_{\text{sep}}$  large  $k_y \lambda_q$  portion of the space. In this figure, open symbols indicate cases in the convective regime. There is good qualitative agreement with the heuristic prediction that the convective regime occupied the region  $k_y \lambda_q > 1$ .

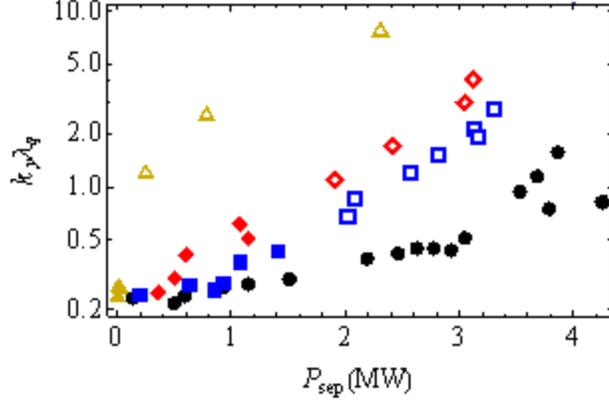


Fig. 8 Power scaling of the heat-flux width normalized to the wavenumber of the fastest growing linear mode. The four cases are the base case with  $b_\theta = b_{\theta 0} = 0.3$  (black disks),  $b_\theta / b_{\theta 0} = 1/2$  (blue squares),  $b_\theta / b_{\theta 0} = 1/3$  (red diamonds), and  $b_\theta / b_{\theta 0} = 1/2$ ,  $\Lambda_{\parallel} / \Lambda_{\parallel 0} = 16$  (gold triangles). Filled (open) symbols indicate the diffusive (convective) regime. The transition occurs at  $k_y \lambda_q \sim 1$

#### 4. Discussion and conclusions

Using SOLT code simulations, we have shown that SOL cross-field transport depends on the net power flux into the SOL,  $P_{\text{sep}}$ , the normalized poloidal magnetic field  $b_\theta$ , the SOL connection length  $L_{\parallel}$  and the plasma collisionality  $\Lambda$ . Below some critical power level,  $P_{\text{sep}} < P_*$ , the near-SOL cross-field transport is quasi-diffusive; whereas for  $P_{\text{sep}} > P_*$  the SOL makes a transition into a convective blob-dominated regime.

The critical power  $P_*$  depends on  $b_\theta$  and  $\Lambda$ . As can be seen from Fig. 8,  $P_*$  increases with  $b_\theta$  and decreases with  $\Lambda$ . One way of characterizing the critical parameters for the transition is that the convective, blob-dominated regime occurs for  $k_y \lambda_q > 1$  where  $k_y$  is a characteristic poloidal wavenumber. The parallel heat flux width  $\lambda_q$  also scales with poloidal magnetic field and collisionality. We find that  $\lambda_q \sim b_\theta^{-m}$  with  $m \sim 0.5$  in the diffusive regime and  $m \sim 1$  in the convective regime. This result is in accord with the inverse scaling of  $\lambda_q$  with plasma current  $I_p$  seen in many experiments. From collisionality scaling studies, we find  $\lambda_q$  increases with  $\Lambda$ , but is nearly independent of the SOL connection length  $L_{\parallel}$  in the sheath-connected regime. This lack of explicit  $L_{\parallel}$  dependence is qualitatively consistent with experiments on Alcator C-Mod.

It is tempting to try to relate the more quiescent, laminar diffusive regime with H-mode confinement, along the lines of what was shown for simple magnetized toroidal plasmas in Ref. 23. Indeed, in our simulations which hold edge (i.e. closed surface)

profiles fixed, and vary  $P_{\text{sep}}$ , low  $P_{\text{sep}}$  corresponds to better confinement than the high  $P_{\text{sep}}$  cases. Also the high confinement low  $P_{\text{sep}}$  cases are those where the strongest  $E_r$  well is imposed on the edge plasma. Furthermore, on NSTX at least, L-mode plasmas produce many more blobs in the SOL than H-mode plasmas [34], which is in accord with our diffusive-convective transition model. At present however, there is insufficient physics in the SOLT code to make more than very speculative remarks about this complex topic.

The present model is rudimentary in many respects, such as its use of 2D geometry and reduced physics. Nevertheless, it is hoped that the observations presented here will further experimental, theoretical and numerical inquiries leading to the development of a fully predictive theory for the SOL heat-flux width and SOL dynamics in tokamak plasmas.

## Acknowledgments

This work was motivated in part by previous modeling [20,21] of data from the NSTX experiment. The authors wish to thank their colleagues J-W Ahn, J Boedo, D P Lundberg, R Maingi, R J Maqueda, D P Stotler, M Umansky and S J Zweben for stimulating discussions of NSTX scrape-off-layer physics and for providing characteristic data for the starting point of these studies. This work was supported by the U.S. Department of Energy (DOE) under grants DE-FG02-97ER54392 and DE-FG02-02ER54678; however, such support does not constitute an endorsement by the DOE of the views expressed herein.

## References

- 1 Wagner F, *et al.* 1982 Phys. Rev. Lett. **49** 1408
- 2 Zweben S J, *et al.* 2010 Phys. Plasmas **17** 102502
- 3 Conway G D, Angioni C, Ryter F, Sauter P, Vicente J, and ASDEX Upgrade Team 2011 Phys. Rev. Lett. **106** 065001
- 4 Canik J M, *et al.* 2011 Phys. Plasmas **18** 056118
- 5 Maingi R, *et al.* 2009 Phys. Rev. Lett. **103** 075001
- 6 Lipschultz B, *et al.* 2007 Nucl. Fusion **47** 1189
- 7 Counsell G F, Connor J W, Erents S K, Field A R, Fielding S J, LaBombard B, Morel K M 1999 J. Nucl. Mater. **266–269** 91
- 8 Itoh S-I and Itoh K 1994 Plasma Phys. Control. Fusion **36** 1845
- 9 Braginskii S I *Transport Processes in a Plasma*, Reviews of Plasma Physics, Vol. 1 (Consultants Bureau, New York, 1965)

- 10 Stangeby P C, Canik J M and Whyte D G 2010 Nucl. Fusion **50** 125003
- 11 Krasheninnikov S I, D'Ippolito D A, and Myra J R 2008 J. Plasma Phys. **74** 679
- 12 Garcia O E 2009 Plasma Fusion Res. **4** 019
- 13 D'Ippolito D A, Myra J R and Zweben S J 2011 Phys. Plasmas **18** 060501
- 14 Zweben S J, *et al.* 2007 Plasma Phys. Controlled Fusion **49** S1
- 15 Lasnier C J, *et al.* 2011 J. Nucl. Mater. **415** S353
- 16 Gray T K, Maingi R, Soukhanovskii V A, Surany J E, Ahn J-W and McLean A G 2011 J Nucl. Mater. **415** S360
- 17 LaBombard B, *et al.* 2011 Phys. Plasmas **18** 056104
- 18 Fundamenski W, *et al.* 2011 Nucl. Fusion **51** 083028
- 19 Kaye S M, *et al.* 2001 Phys. Plasmas **8** 1977
- 20 Myra J R, Russell D A, D'Ippolito D A, Ahn J-W, Maingi R, Maqueda R J, Lundberg D P, Stotler D P, Zweben S J, Umansky M 2011 J. Nucl. Mat. **415** S605
- 21 Myra J R, *et al.* 2011 Plasmas **18** 012305
- 22 Garcia O E, Bian N H, Paulsen J-V, Benkadda S and Rypdal K 2003 Plasma Phys. Control. Fusion **45** 919
- 23 Ricci P, Rogers B N, and Brunner S 2008 Phys. Rev. Lett. **100** 225002
- 24 Sarazin Y, Garbet X, Ghendrih Ph and Benkadda S 2000 Phys. Plasmas **7** 1085
- 25 Diamond P H, Itoh S-I, Itoh K, and Hahn T S 2005 Plasma Phys. Control. Fusion **47** R35
- 26 Benkadda S, Klochkov D N, Popel S I, and Izvekova Yu N 2011 Phys. Plasmas **18** 052306
- 27 Russell D A *et al.* 2011 Bull. Am. Phys. Soc. **56**, 140 (paper JO4-7)
- 28 Russell D A, Myra J R and D'Ippolito D A 2009 Phys. Plasmas **16** 122304
- 29 Wakatani M and Hasegawa A 1984 Phys. Fluids **27** 611
- 30 Loarte A, *et al.* 1999 J. Nucl. Mater. **266-269** 587
- 31 Garcia O E, Horacek J, Pitts R A, Nielsen A H, Fundamenski W, Naulin V and Rasmussen J J 2007 Nucl. Fusion **47** 667
- 32 Russell D A, Myra J R and D'Ippolito D A 2007 Phys. Plasmas **14** 102307
- 33 Myra J R, Russell D A and D'Ippolito D A 2006 Phys. Plasmas **13** 112502
- 34 Agostini M, Zweben S J, Cavazzana R, Scarin P, Serianni G, Maqueda R J, and Stotler D P 2007 Phys. Plasmas **14** 102305