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A Model Of Pedestal Structure

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Abstract

Predictions are developed for the structure of plasma parameter profiles of H-mode pedestals in transport quasi-equilibrium in tokamak plasmas. They are based on assuming paleoclassical radial plasma transport processes dominate throughout the pedestal. The key physical process in this model is that the electron temperature gradient in the pedestal increases to the magnitude required for paleoclassical electron heat transport to carry the large conductive radial electron heat flow from the hot core through the pedestal to the separatrix. The concomitant level of paleoclassical density transport is usually large in the pedestal compared to local fueling due to neutral recycling from outside the separatrix. Thus, in this model the pedestal density profile is usually determined not by edge fueling but rather by a combination of the separatrix density boundary condition and the pedestal density profile needed for the outward paleoclassical diffusive flux to be nearly balanced by the inward paleoclassical pinch flow. When neutral fueling effects are significant they add to the pedestal density and displace the density profile outward from the electron temperature profile. Model predictions are given for the electron density and temperature gradients, profiles and magnitudes in the pedestal. The transition into electron-temperature-gradient (ETG) driven anomalous radial electron heat transport in the core plasma determines the initial, transport-limited height of the electron pressure pedestal. Characteristics of the plasma toroidal rotation profile in the pedestal are also predicted. Model predictions are found to agree quantitatively (within about a factor of about two) with the properties of the recently studied 98889 DIII-D pedestal [J.D. Callen et al., Nucl. Fusion **50**, 064004 (2010)]. Applications to other outstanding H-mode pedestal structure and evolution issues in tokamaks are also discussed. Finally, a hierarchy of experimental validation tests are suggested.

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I. INTRODUCTION

A recent study [1] of plasma transport properties in a single DIII-D [2] H-mode pedestal indicated it is quite plausible that paleoclassical plasma transport [3] plays a significant role in the pedestal. Specifically, pedestal electron heat and density transport (diffusivities plus density pinch) were found to be in reasonable agreement with paleoclassical predictions. This report develops a comprehensive set of paleoclassical model predictions for plasma profiles in transport "quasi-equilibrium" H-mode pedestals, i.e., for the "steady-state" "structure" of pedestals that have properties similar to those in the 98889 DIII-D pedestal. These predictions are compared quantitatively with data from that DIII-D pedestal [1]. They are also used to interpret a number of properties in similar H-mode pedestals in DIII-D and other tokamaks, and the evolution toward ELMs in them. Three sets of experimental validation tests for this new pedestal structure model are also presented.

The fundamental physics of the model developed here can be understood as follows. First, one notes that often most of the heat flowing through pedestals is carried by electrons — about 75 % in the 98889 DIII-D pedestal [1]. Because most plasma heating occurs in the hotter and denser core, a large electron heat flux must be carried out through the lower density and temperature pedestal. Thus, from a Fourier heat flux law $\mathbf{q}_e = -n_e \chi_e \boldsymbol{\nabla} T_e$, the electron temperature gradient in the pedestal will increase until this large heat flux can be carried out through the pedestal to the separatrix. The needed electron temperature gradient is determined by the effective electron heat diffusivity χ_e in the pedestal.

The key assumption here is that paleoclassical processes dominate χ_e and all plasma transport channels (except for ion heat transport which has a large neoclassical component) throughout the pedestal and determine its structure. The inherent level of paleoclassical density transport is usually large in the pedestal compared to the typically small local fueling due to edge neutral recycling. Thus, in this model the density profile in the pedestal is usually determined not by local fueling in the pedestal but rather by what is needed for the paleoclassical density transport operator to be small, i.e., by the density profile needed for its outward diffusive flux to be nearly balanced by its intrinsic inward pinch flow. The pinch results naturally from the structure of the paleoclassical density transport operator.

Anomalous plasma transport due to microturbulence is typically dominant in the hotter, denser core; the transition to it at the top of the pedestal will determine the initial, transport-limited height of the electron pressure pedestal. Thus, while fluctuations are usually present in pedestals and may peak where the plasma pressure gradient is largest, they will be assumed to be too small to contribute significantly to plasma transport in the pedestal. Anomalous transport fluxes will be allowed for in the development of general plasma transport and flow equations. Then, the anomalous transport effects will be explicitly dropped in determining the structure of plasma profiles within the pedestal, but re-introduced to determine the initial, transport-limited pedestal height.

This report is organized as follows. The next two sections present formulas for key paleoclassical transport model parameters and the simplified plasma transport equations that will be used here. The following sections develop predictions for key properties of electron density, electron temperature and ion temperature profiles as well as the toroidal plasma flow profile in the pedestal region. Dimensionless variable scalings of the key results are developed in Section VIII. The following section discusses some key issues for and other applications of this model. The penultimate section suggests a hierarchy of experimental validation tests for this new pedestal structure model. The final section summarizes the main pedestal structure predictions developed in this report and their implications.

II. KEY PALEOCLASSICAL PARAMETERS

The initial paleoclassical papers [4–6] were based on a key hypothesis that charged particles diffuse radially along with thin annuli of poloidal magnetic flux in resistive, current-carrying toroidal plasmas. That is, they diffuse radially with the magnetic field diffusivity $D_\eta = \eta/\mu_0$, in which η is the plasma electrical resistivity. This hypothesis was later shown [7–9] to result from transforming the drift-kinetic equation from laboratory to poloidal magnetic flux coordinates, upon which Grad-Shafranov equilibria, neoclassical transport theory and gyrokinetic-based microturbulence-induced anomalous transport analyses are based.

The diffusive components of paleoclassical transport are all proportional to the plasma resistivity η whose Spitzer form scales as $T_e^{-3/2}$. Paleoclassical radial electron heat transport can be greater than fluctuation-induced gyro-Bohm-level transport ($\sim T_e^{3/2}/aB^2$) wherever $T_e < B(\mathrm{T})^{2/3}a(\mathrm{m})^{1/2}$ keV [10], which usually occurs in ohmic tokamak plasmas, or wherever fluctuations are significantly reduced. Thus, paleoclassical electron heat transport is likely to be dominant in H-mode pedestals [10] — for $T_e \lesssim 1$ keV in DIII-D and $T_e \lesssim 5$ keV in ITER [11]. Paleoclassical transport predictions compare favorably with experimental data from many ohmic-level toroidal plasmas [10] (tokamaks, STs, RFPs, spheromaks), in electron-cyclotron (EC) heated RTP plasmas [12] and in tokamak H-mode pedestals [1, 10].

Since D_η is fundamental to all paleoclassical plasma transport processes [13, 14], its evaluation and scaling will be specified first. In DIII-D pedestals the inverse aspect ratio is not small. Hence, trapped particles are a large fractional component and their effects are important. Thus, electron viscosity effects on parallel electron flows induced by trapped particles must be included in calculating the plasma resistivity η; they can increase the resistivity relative to Spitzer resistivity significantly — by factors of up to about 6 in pedestals.

The fundamental parameter of the paleoclassical transport model is [5, 6, 13, 14]

$$\boxed{D_\eta \equiv \frac{\eta_\parallel^{\rm nc}}{\mu_0} = \frac{\eta_0}{\mu_0}\,\frac{\eta_\parallel^{\rm nc}}{\eta_0},} \quad \text{magnetic field diffusivity, m}^2\text{/s.} \tag{1}$$

Since $\mu_0 \equiv 4\pi \times 10^{-7}$ N/A^2 (in SI, MKS units) is a constant of nature, the key parameter is really the parallel neoclassical resistivity $\eta_\parallel^{\rm nc}$. The reference, perpendicular resistivity $\eta_0 \equiv m_e \nu_e / n_e e^2$ can be written in the form of a magnetic field diffusivity as

$$\boxed{\frac{\eta_0}{\mu_0} \simeq \frac{1400\,Z}{[T_e(\text{eV})]^{3/2}}\left(\frac{\ln\Lambda}{17}\right),} \quad \text{reference magnetic field diffusivity, m}^2\text{/s.} \tag{2}$$

Here and below, $Z \to Z_{\rm eff} \equiv \sum_i n_i Z_i^2/n_e$ is the effective ion charge and $\ln\Lambda$ is the Coulomb logarithm (~ 14.5 at the separatrix but ~ 16 at the pedestal top for 98889 DIII-D parameters [1]). The parallel neoclassical resistivity $\eta_\parallel^{\rm nc}$ (Ohm-m) can be evaluated using formulas in the literature [15–19]; the NCLASS [20] or NEO [21] codes evaluate it most precisely.

A useful approximate formula for the ratio of neoclassical resistivity to η_0 is [5, 6]

$$\boxed{\frac{\eta_\parallel^{\rm nc}}{\eta_0} \simeq \frac{\eta_\parallel^{\rm Sp}}{\eta_0} + \frac{\mu_e}{\nu_e},} \quad \text{parallel neoclassical resistivity factor.} \tag{3}$$

The two components of the parallel resistivity in a tokamak are [5, 6]

$$\frac{\eta_\parallel^{\rm Sp}}{\eta_0} \simeq \frac{\sqrt{2}+Z}{\sqrt{2}+13Z/4}, \quad \text{Spitzer parallel resistivity factor, and} \tag{4}$$

$$\frac{\mu_e}{\nu_e} \simeq \frac{Z+\sqrt{2}-\ln(1+\sqrt{2})}{Z\,(1+\nu_{*e}^{1/2}+\nu_{*e})}\,\frac{f_t}{f_c} \overset{\nu_{*e}\to 0}{\underset{Z=1}{\Longrightarrow}} 1.53\,\frac{f_t}{f_c}, \quad \text{parallel electron viscosity effects.} \tag{5}$$

Here, f_c is the flow-weighted fraction of circulating particles [17] with Padé approximate [22]

$$f_c \simeq \frac{(1-\epsilon^2)^{-1/2}\,(1-\epsilon)^2}{1+1.46\,\epsilon^{1/2}+0.2\,\epsilon} \simeq 1 - 1.46\,\epsilon^{1/2} + \mathcal{O}(\epsilon), \quad \text{circulating particle fraction,} \tag{6}$$

in which $\epsilon \equiv (B_{\max} - B_{\min})/(B_{\max} + B_{\min}) \simeq r_M/R_0$ is the local inverse aspect ratio. The fraction of trapped particles is $f_t \equiv 1 - f_c$. Other approximate formulas for f_c have also been presented in the literature [23, 24]. For 98889 DIII-D pedestal parameters [1] $\epsilon \equiv r_M/R_0 \simeq 0.6/1.7 \simeq 0.35$ and hence $f_t/f_c \simeq 0.77/0.23 \simeq 3.35 \gg 1$. For the highly noncircular near-separatrix geometry in DIII-D, f_c should really be evaluated numerically [17, 18]. Finally,

$$\nu_{*e} \equiv \frac{\nu_e}{\epsilon^{3/2}(v_{Te}/R_0\, q)} = \frac{R_0\, q}{\epsilon^{3/2}\lambda_e}, \quad \text{neoclassical electron collisionality parameter,} \qquad (7)$$

in which $v_{Te} \equiv \sqrt{2T_e/m_e}$ and the electron Coulomb collision "mean free path" is

$$\lambda_e \equiv \frac{v_{Te}}{\nu_e} \simeq 1.2 \times 10^{16} \frac{[T_e(\mathrm{eV})]^2}{Z\, n_e(\mathrm{m}^{-3})} \left(\frac{17}{\ln\Lambda}\right), \quad \text{electron collision length, m.} \qquad (8)$$

The accuracy of the $\eta_\parallel^{\rm nc}$ from Eqs. (3)–(8) ranges from being about equal to (better than $1/\ln\Lambda \sim 7\%$ accuracy for $\mu_e/\nu_e << 1$), to as much as twice as large as (for $\mu_e/\nu_e \gg 1$) the most precise neoclassical resistivity evaluations [17–20]. It is important to note that ONETWO [25] usually uses the large aspect ratio Hirshman, Hawryluk and Birge (HHB) [16] formula for $\eta_\parallel^{\rm nc}$ which can be up to a factor of 2 too small in DIII-D pedestals when trapped particle (μ_e/ν_e) effects are significant. The magnetic field diffusivity D_η can be evaluated from ONETWO output files using the HHB value of $\eta_\parallel^{\rm nc} \equiv$ `eta` Ohm-cm to yield $D_\eta =$ `eta`$/(4\pi \times 10^{-5})$ m^2/s. The $Z_{\rm eff}$ and finite ϵ, ν_{*e} effects in Eqs. (1)–(8) are all very important in determining the magnitude and scaling of D_η in H-mode pedestals.

An approximate formula for D_η for the parameters of the 98889 DIII-D pedestal where $Z_{\rm eff} \simeq 2.83$ and $f_t/f_c \simeq 3.35$ can be written using $\ln\Lambda \simeq 15$ as

$$D_\eta \simeq \frac{1240\, Z_{\rm eff}}{[T_e(\mathrm{eV})]^{3/2}} \left(0.4 + \frac{4}{1 + \nu_{*e}^{1/2} + \nu_{*e}}\right) \ \mathrm{m^2/s}, \quad \text{for 98889 DIII-D pedestal [1].} \qquad (9)$$

Using (8), the neoclassical electron collisionality parameter is ($\sqrt{2}$ smaller than in [15, 18])

$$\nu_{*e} \simeq 6 \times 10^{-16} \frac{q\, Z_{\rm eff}\, n_e(\mathrm{m}^{-3})}{[T_e(\mathrm{eV})]^2}, \quad \text{for 98889 DIII-D pedestal [1].} \qquad (10)$$

This equation and Fig. 10a in Ref. [1] indicate $\nu_{*e} \propto q$ ranges from ∞ on the separatrix where $q \to \infty$, to 1.1 at the pedestal mid-point (density tanh fit symmetry point, $\rho_n \simeq 0.982a$), to about 0.6 at the pedestal top (here $\rho_{\rm t} \simeq 0.96a$). Thus, the $\eta_\parallel^{\rm nc}/\eta_0$ factor in parentheses in Eq. (9) is about 0.4, 1.67, 2.1 at these points. Hence, while $1/T_e^{3/2}$ decreases by a factor of $(360/90)^{3/2} \simeq 8$ from the separatrix to the 98889 pedestal mid-point ρ_n, the $\eta_\parallel^{\rm nc}/\eta_0$ factor increases by a factor of 4.2. Thus, from (9) the neoclassical resistivity and D_η only decrease by a factor of about 1.9 over the bottom half of the pedestal. The corresponding ONETWO ratio for 98889 is about 4.4. More precise quantification of these factors requires evaluation of $\eta_\parallel^{\rm nc}$ using the formulas in Refs. [18, 19], or by the NCLASS [20] or NEO [21] codes.

While most paleoclassical transport processes are governed by D_η, paleoclassical electron heat transport includes an additional factor M [5]. This factor arises from the addition of helically resonant radial electron heat transport contributions [5] in the vicinity of medium order rational surfaces — such as where $q = m/n$ with $n \sim 2$–4 and $m \sim 8$–20 in pedestals. The multiplier M can be written as a smoothed (Padé approximate) formula [5, 6, 10]

$$\boxed{M \simeq \frac{1/(\pi R_0 q)}{1/\ell_{\rm max} + 1/\lambda_e} \sim \frac{\min\{\ell_{\rm max},\, \lambda_e\}}{\pi \bar{R} q} \simeq \frac{\lambda_e}{\pi R_0 q},} \quad \text{paleoclassical helical multipler.} \tag{11}$$

It is determined by the minimum of the collision length λ_e defined in (8) and an effective parallel length over which field lines are diffusing radially [5]:

$$\ell_{\rm max} \equiv \pi \bar{R} q \, n_{\rm max}, \quad \text{parallel length of diffusing field lines, m,} \tag{12}$$

in which the maximum order n of the medium order rational surfaces is

$$n_{\rm max} \equiv \left(\frac{1}{\pi \bar{\delta}_e |q'|} \right)^{1/2}, \quad \text{maximum } n \text{ for diffusing field lines.} \tag{13}$$

Here, $\bar{\delta}_e$ is a dimensionless electromagnetic (em) skin depth factor defined by

$$\bar{\delta}_e \equiv \frac{\delta_e}{\bar{a}}, \quad \delta_e \equiv \frac{c}{\omega_p} \simeq 10^{-3} \left(\frac{3 \times 10^{19}}{n_e ({\rm m}^{-3})} \right)^{1/2} {\rm m} \sim 10^{-3} \ {\rm m}, \quad \text{em skin depth.} \tag{14}$$

Also, $q' \equiv dq/d\rho$ in which q is the "safety factor" (toroidal winding number or inverse of field line pitch). In the vicinity of flux surfaces at which q' vanishes (e.g., near the mid-point of H-mode pedestals because of the large bootstrap current there, as indicated in Fig. 9a of Ref. [1]), the magnitude of $n_{\rm max}$ is limited by q'' [5]. Throughout most of the 98889 DIII-D pedestal $\lambda_e < \ell_{\rm max}$ so $M \simeq \lambda_e / \pi R_0 q$; this approximation for M decreases from about 2.6 at the pedestal top, to 1.4 at the pedestal mid-point, to zero at the separatrix (where $q \to \infty$). However, $\ell_{\rm max}$ can be $\sim \lambda_e$ for $\rho < \rho_{\rm t} \simeq 0.96a$ and reduce M by a factor of ≤ 2 there.

The $\bar{a}$ parameter in (14) is the characteristic radius of the plasma for paleoclassical transport processes. It is defined in terms of the usual separatrix toroidal-flux-determined minor radius [1] $a \equiv \sqrt{\psi_{\rm t}({\rm sep})/\pi B_{\rm t0}}$ by (here $\langle \cdots \rangle$ indicates averaging over a flux surface)

$$\frac{a^2}{\bar{a}^2} = \left\langle \frac{|\boldsymbol{\nabla} \rho|^2}{R^2 \langle R^{-2} \rangle} \right\rangle, \quad \text{near unity paleoclassical geometry factor.} \tag{15}$$

This factor is slightly smaller than the $\langle|\nabla\rho|^2\rangle$ factor that arises from standard Fick's or Fourier law diffusive fluxes, which was important in the Ref. [1] transport analysis. It can be evaluated from ONETWO output files using the relation $a^2/\bar{a}^2 \equiv$ `gcap / <R0**2/R**2>`. In the 98889 DIII-D pedestal $a^2/\bar{a}^2 \simeq 1.6$ whereas from Fig. 7 in Ref. [1] $\langle|\nabla\rho|^2\rangle \simeq 2$.

III. PLASMA TRANSPORT EQUATIONS

In plasma transport quasi-equilibrium states shortly after an L-H transition or just before an edge-localized-mode (ELM) such as those discussed in Refs. [1, 26], the "steady-state" ($\partial/\partial t \to 0$) flux-surface-averaged (FSA) density and energy transport equations [14] can be written as

$$\langle \nabla\cdot\mathbf{\Gamma}\rangle \;\equiv\; \frac{1}{V'}\,\frac{d}{d\rho}(V'\Gamma) \;=\; \langle S_n\rangle, \quad \text{density}, \tag{16}$$

$$\langle \nabla\cdot\mathbf{q}\rangle \;\equiv\; \frac{1}{V'}\,\frac{d}{d\rho}[V'(\Upsilon + \frac{5}{2}\,T\,\Gamma)] \;=\; Q^{\rm net}, \quad \text{energy}. \tag{17}$$

Here: $\mathbf{\Gamma}$, $\mathbf{q}$ are the density, heat fluxes; $\Gamma \equiv \langle\mathbf{\Gamma}\cdot\nabla\rho\rangle$, $\Upsilon \equiv \langle\mathbf{q}\cdot\nabla\rho\rangle$ are the corresponding FSA fluxes; and $\langle S_n\rangle$, $Q^{\rm net}$ are the net FSA density, energy sources. Also, $\rho \equiv \sqrt{\psi_{\rm t}/\pi B_{\rm t0}}$ is the average radius of a flux surface, which has units of m and is the usual radial variable `rho` in ONETWO. Finally, $V' \equiv dV/d\rho$ (m^{-2}) is the radial derivative of the volume $V(\rho)$ (m^{-3}) of the ρ flux surface. In terms of ONETWO variables, it is $V' = 4\pi^2 R_0$`*rho*hcap`.

The density equation will be considered first. The paleoclassical density transport operator can be written [13, 14] in the standard form of the divergence of a density flux $\Gamma^{\rm pc}$:

$$\boxed{\Gamma^{\rm pc} = -\,\frac{1}{V'}\,\frac{d}{d\rho}(V'\bar{D}_\eta\,n_e) = -\,\bar{D}_\eta\,\frac{dn}{d\rho} \;+\; n\,V_{\rm pinch},} \quad \text{paleoclassical density flux.} \tag{18}$$

As indicated, $\Gamma^{\rm pc}$ naturally includes a pinch flow, which results from particle guiding centers being diffused radially (with no drag-type contribution [5]) along with diffusing thin annuli of poloidal magnetic flux. The pinch is usually inward (i.e., < 0) and is defined by

$$\boxed{V_{\rm pinch} \;\equiv\; -\,\frac{1}{V'}\,\frac{d}{d\rho}\,(V'\bar{D}_\eta),} \quad \text{paleoclassical pinch flow velocity, m/s.} \tag{19}$$

Here, the geometric factor $a^2/\bar{a}^2$ has been incorporated in a modified magnetic field diffusivity for simplicity of notation throughout the remainder of this report:

$$\boxed{\bar{D}_\eta \;\equiv\; \frac{a^2}{\bar{a}^2}\,D_\eta,} \quad \text{geometrically effective } D_\eta,\ \mathrm{m}^2/\mathrm{s}. \tag{20}$$

The (ambipolar) anomalous density flux $\Gamma^{\rm an}$ will be left general for now. Adding the paleoclassical and anomalous density transport fluxes, the density (continuity) equation becomes

$$\frac{1}{V'}\frac{d}{d\rho}\left(-\frac{d}{d\rho}(V'\bar{D}_\eta\, n) + V'\Gamma^{\rm an}\right) = \langle S_n\rangle. \tag{21}$$

Multiplying this equation by the differential volume $dV \equiv V'd\rho$ and integrating over ρ from the ρ flux surface to the separatrix ($\rho_{\rm sep} \equiv a$) yields the particle flow (#/s) equation that will be analyzed in the next section to determine properties of the pedestal density profile:

$$-\left[\frac{d}{d\rho}(V'\bar{D}_\eta\, n)\right]_\rho = \dot{N}(\rho) + V'\,\Gamma^{\rm an}. \tag{22}$$

Here, $\dot{N}(\rho)$ is rate of flow of charged particles (#/s) across the ρ surface:

$$\dot{N}(\rho) \equiv -\left[\frac{d}{d\rho}(V'\bar{D}_\eta\, n)\right]_a - \int_\rho^a V'(\hat{\rho})\, d\hat{\rho}\ \langle S_n(\hat{\rho})\rangle, \tag{23}$$

whose two terms represent the charged particle flow through the separatrix minus an integral correction due to the local density source $\langle S_n\rangle$ between the ρ surface and the separatrix.

Since the dominant heat flow through the pedestal is usually via electrons, the electron energy balance will be discussed next. The FSA paleoclassical radial electron heat transport operator is not in the usual form of the divergence of a radial electron heat flux. Rather, it is a multiplier $M+1$ times a divergence [5]. The paleoclassical electron heat transport operator can be written as (see Eq. (142) in [5] or the sum of Eqs. (47) and (72) in [14]):

$$\langle \boldsymbol{\nabla\cdot} \mathbf{q}_e^{\rm pc}\rangle = -\frac{M+1}{V'}\frac{d^2}{d\rho^2}\left(V'\bar{D}_\eta\,\frac{3}{2}\,n_eT_e\right), \quad \frac{\rm W}{\rm m^3}. \tag{24}$$

While the paleoclassical electron heat transport operator is non-standard, it usually yields dominantly diffusive radial electron heat transport with an effective $\chi_e^{\rm pc} \simeq (3/2)(M+1)\bar{D}_\eta$. The anomalous electron heat flux $\Upsilon_e^{\rm an}$ will be left general for now. Adding the paleoclassical and anomalous electron heat fluxes, the electron energy equation becomes

$$-\frac{M+1}{V'}\frac{d^2}{d\rho^2}\left(V'\bar{D}_\eta\,\frac{3}{2}\,n_eT_e\right) + \frac{1}{V'}\frac{d}{d\rho}\left[V'\left(\Upsilon_e^{\rm an} + \frac{5}{2}\,T_e\Gamma\right)\right] = Q_e^{\rm net}. \tag{25}$$

Multiplying this equation by $V'/(M+1)$ and then integrating over ρ from a given ρ surface outward to the separatrix at a yields the equation for electron heat flow (Watts) that will be analyzed in Section V to determine properties of the T_e profile:

$$-\left[\frac{d}{d\rho}\left(V'\bar{D}_\eta\,\frac{3}{2}\,n_eT_e\right)\right]_\rho = \hat{P}_e(\rho) + \int_\rho^a \frac{d\hat{\rho}}{M(\hat{\rho})+1}\,\frac{d}{d\hat{\rho}}[V'(\hat{\rho})\Upsilon_e^{\rm an}(\hat{\rho})]. \tag{26}$$

Here, $\hat{P}_e(\rho)$ is an effective conductive electron heat flow (Watts) through the ρ surface:

$$\hat{P}_e(\rho) \equiv -\left[\frac{d}{d\rho}\left(V'\bar{D}_\eta\frac{3}{2}n_eT_e\right)\right]_a - \int_\rho^a \frac{V'(\hat{\rho})\,d\hat{\rho}}{M(\hat{\rho})+1}\left[Q_e^{\rm net} - \frac{1}{V'}\frac{d}{d\rho}\left(V'\frac{5}{2}T_e\Gamma\right)\right]_{\hat{\rho}}. \quad (27)$$

This is the electron heat flow through the separatrix (since $M \to 0$ at the separatrix) minus usually small integral corrections due to local electron heating $Q_e^{\rm net}$ and convective electron heat flow between the ρ surface and separatrix. The local heating $Q_e^{\rm net}$ and convective heat flow are nearly negligible in the 98889 DIII-D pedestal (see Figs. 4a and 5a in Ref. [1]).

The ion energy equation is considered next. The paleoclassical ion heat transport operator is similar to the electron heat transport operator in (24), except that $M_i \to 0$ (because there is no helically resonant ion contribution since the ion toroidal precessional drift frequency is larger than the ion collision frequency — see discussion after Eq. (139) in Ref. [5]). Thus, unlike for paleoclassical electron heat transport, it can be written in the usual heat flux divergence form [14] with $\Upsilon_i^{\rm pc} = -(1/V')(d/d\rho)[V'\bar{D}_\eta(3/2)n_iT_i]$. Neoclassical ion heat transport is usually important in pedestals [1]. Its ion heat flux will be written as $\Upsilon_i^{\rm nc} = -n_i\chi_i^{\rm nc}\,dT_i/d\rho$. Then, labeling the anomalous ion heat flux as $\Upsilon_i^{\rm an}$ and performing the same type of analysis as was done for electrons to obtain (26) yields for the ion heat flow equation

$$-\left[\frac{d}{d\rho}\left(V'\bar{D}_\eta\frac{3}{2}n_iT_i\right)\right]_\rho + V'\left(-n_i\chi_i^{\rm nc}\frac{dT_i}{d\rho} + \Upsilon_i^{\rm an}\right) = P_i(\rho), \quad (28)$$

in which $P_i(\rho)$ is the conductive ion heat flow (Watts) through the ρ surface. That is, it has the same form as $P_e(\rho)$ in (27) except $M_i \to 0$ [same form as $\dot{N}$ in (23)] plus perhaps additional neoclassical and anomalous ion heat flows through the separatrix. This equation determines the T_i pedestal profile; properties of its solution will be discussed in Section VI.

IV. ELECTRON DENSITY PROFILE

Deuteron and impurity (carbon) densities in the pedestal are influenced by differing source, ionization and neoclassical pinch-type effects. These ion complications will be circumvented by determining the electron pedestal density profile. Neglecting the anomalous density flux $\Gamma_e^{\rm an}$ in Eq. (22) and integrating it over ρ from ρ to the separatrix at a yields

$$n_e(\rho)\,\bar{D}_\eta(\rho)\,V'(\rho) = n_e(a)\,\bar{D}_\eta(a)\,V'(a) + \int_\rho^a d\hat{\rho}\,\dot{N}_e(\hat{\rho}). \quad (29)$$

Thus, the combination of parameters $n_e\bar{D}_\eta V'$ is equal to its value on the separatrix plus the integral of the fueling source $\dot{N}_e$ between ρ and the separatrix. However, the effect of the fueling source is often negligible. For example, at the density mid-point ($\rho_n \simeq 0.982a$) of the 98889 DIII-D pedestal [1] the estimated ratio of the fueling to $n_e \bar{D}_\eta V'$ is negligibly small:

$$\frac{\int_{\rho_n}^{a} d\hat{\rho}\, \dot{N}_e(\hat{\rho})}{[n_e \bar{D}_\eta V']_{\rho_n}} \simeq \frac{(a-\rho_n)\, \dot{N}_e[(a+\rho_n)/2]}{n_e(\rho_n)\, \bar{D}_\eta(\rho_n)\, V'(\rho_n)} \simeq \frac{(0.018)(0.77)(2\times 10^{21})}{(1.65\times 10^{19})(1)(42.6)} \simeq 0.04 \ll 1. \qquad (30)$$

Neglecting $\dot{N}_e$ and the variation of $V' \simeq (\rho/a)V'(a) \simeq (\rho/a)(43.4)$, to lowest order (29) is

$$\boxed{n_e(\rho)\, \bar{D}_\eta(\rho) \simeq \text{constant} \quad \Longrightarrow \quad n_e(\rho) \simeq n_e(a)\, \frac{\bar{D}_\eta(a)}{\bar{D}_\eta(\rho)}, \quad \text{within the pedestal.}} \qquad (31)$$

For the 98889 DIII-D pedestal [1] the estimated ratio of $\bar{D}_\eta$ from the separatrix to pedestal mid-point ranges from $\simeq 4.4$ (ONETWO) to $\simeq 1.9$ from (9). These factors predict an n_e ratio within a factor of two of the experimental density ratio of $n_e(\rho_n)/n_e(a) \simeq 1.65/0.77 \simeq 2.14$. Note that if D_η were spatially constant (e.g., as in a Fick's diffusion law with a constant D), (31) would predict $n_e(\rho) =$ constant in the pedestal, i.e., no density pedestal at all.

The relation in (31) just reflects what is required to produce a small net electron paleoclassical density flux $\Gamma_e^{\rm pc}$ in (18) — by nearly balancing the outward diffusive density transport with a large inward pinch flow. That is, since in 98889 the fueling source flow $\int_\rho^a d\hat{\rho}\, \dot{N}_e$ (m/s) is small relative to the intrinsic paleoclassical transport flow $n_e \bar{D}_\eta V'$ (m/s), the density profile adjusts to produce a small net radial density flow through the pedestal. This scenario is precisely what was concluded from a pioneering interpretive analysis of the ion density pinch and transport [27] in the 98889 DIII-D pedestal; and paleoclassical predictions were consistent with the pinch and inferred diffusivity values obtained [1].

As indicated in (9), the magnetic field diffusivity D_η in pedestals varies as $Z_{\rm eff}/T_e^{3/2}$ times a function of the electron collisionality parameter $\nu_{*e}(T_e, Z_{\rm eff}, n_e)$. Thus, it depends strongly on the yet to be determined T_e profile. A general scaling relation for n_e cannot be developed from (31). However, since D_η is a mainly a function of T_e, the n_e and T_e profiles should be "aligned" (i.e., strongly correlated) whenever local fueling effects are negligible.

While fueling effects are negligible in the 98889 DIII-D pedestal, they can become significant when neutral beam (NB) core heating and fueling, and consequently edge recycling, are large. The fueling effects in (29) can be estimated by assuming $\dot{N}_e(\rho) \simeq \dot{N}_e(a)\, e^{-(a-\rho)/\lambda_n}$, in which λ_n is an effective neutral penetration distance [28, 29]. Then, Eq. (29) becomes

$$n_e(\rho)\, \bar{D}_\eta(\rho)\, V'(\rho) = n_e(a)\, \bar{D}_\eta(a)\, V'(a) + \dot{N}_e(a)\, [\lambda_n(1 - e^{-(a-\rho)/\lambda_n})]. \qquad (32)$$

The term in square brackets multiplying the separatrix fueling density $\dot{N}_e(a)$ is effectively just the minimum of the distance $a-\rho$ in from the separatrix and λ_n. When the edge fueling source $\dot{N}_e$ effect is large enough (particularly at the density mid-point ρ_n when $\lambda_n \gtrsim a-\rho_n$), it causes the electron density to increase aprroximately linearly with distance in from the separatrix: $n_e(\rho)\,\bar{D}_\eta(\rho)\,V'(\rho) \;\simeq\; n_e(a)\,\bar{D}_\eta(a)\,V'(a) \,+\, \dot{N}_e(a)\,(a-\rho)$ for $\lambda_n > a-\rho$. This causes the pedestal density profile to be shifted radially outward compared to the pedestal T_e profile which it would otherwise mimic if local fueling effects were negligible and n_e was predicted by (31). Such an outward shift of the n_e profile relative to the T_e profile has been observed in the high NB power (and high toroidal field, low ϱ_*) DIII-D data in the JET/DIII-D comparison experiments [30, 31]. In contrast, the JET pedestal profiles of n_e and T_e are aligned — apparently because the fueling per unit of spatial scale size $[\sim \dot{N}_e(a)\;\min\{a-\rho_n, \lambda_n\}/V'(a)]$ is smaller in JET than in DIII-D.

Processes that limit the height of the density pedestal will be considered next. There are three types of possibilities. First, note that, neglecting fueling effects, in moving radially inward from the separatrix toward the pedestal top D_η decreases as T_e increases, but becomes constant as T_e "saturates." When this effect dominates, n_e at the pedestal top would be predicted from (31). Second, in the absence of anomalous transport, the density gradient and gradient scale length at which pinch effects can be neglected (for $d\ln\bar{D}_\eta/d\rho \lesssim 1/a$) and fueling effects become significant can be estimated from an approximate diffusive form of (22): $-\,V'\bar{D}_\eta\,dn_e/d\rho \;\sim\; \dot{N}_e$, which yields $L_n \equiv n_e/(-dn_e/d\rho) \;\sim\; n_e\,\bar{D}_\eta V'/\dot{N}_e$. A prediction for the maximum density at the pedestal top ($\rho_{\rm t}$) can be inferred by requiring L_n to be less than the minor radius a:

$$\boxed{\max\{n_e^{\rm ped}\} \;\sim\; \frac{\dot{N}_e(\rho_{\rm t})\;a}{\bar{D}_\eta(\rho_{\rm t})\,V'(\rho_{\rm t})},} \quad \text{maximum pedestal density when} \;\; \left.\frac{d\ln\bar{D}_\eta}{d\rho}\right|_{\rho_{\rm t}} \lesssim \frac{1}{a}. \qquad (33)$$

At the top ($\rho_{\rm t} \simeq 0.96a$) of the 98889 DIII-D density pedestal [1], $\dot{N}_e \simeq 0.5\times10^{21}$/s (mainly due to core NB fueling), $a \simeq 0.77$ m, $V' \simeq 41.7$ m^2 and $\bar{D}_\eta \simeq 0.3$ m^2/s (ONETWO), which yields $n_e^{\rm ped} \simeq 3.1\times10^{19}$ m^{-3}. This prediction is close to the observed [1] $n_e^{\rm ped} \simeq 3\times10^{19}$ m^{-3}. The $\dot{N}_e$ scaling in (33) might be testable by comparing NB-heated H-mode plasmas with EC heated ones with no core fueling where edge fueling is dominant and $\dot{N}_e(\rho_{\rm t})$ is small.

The third possibility is that anomalous density transport becomes dominant as one moves from the pedestal into the core. Assuming a Fick's diffusion law $\Gamma_e^{\rm an} = -\,\langle|\boldsymbol{\nabla}\rho|^2\rangle D_e^{\rm an} dn_e/d\rho$ in (22), one needs $D_e^{\rm an} \gtrsim D_\eta$ ($\sim$ 0.2 m^2/s at the top of the 98889 DIII-D pedestal [1]) for anomalous transport to become dominant and thereby determine the density there. A scaling relation for testing this possibility could be developed from $D_e^{\rm an} \sim D_\eta$ if and when a formula for the microturbulence-induced $D_e^{\rm an}$ becomes available.

V. ELECTRON TEMPERATURE PROFILE

The analysis of the pedestal density profile in the preceding section was incomplete in that D_η, which depends on the electron temperature T_e, was not determined. This section uses various approximations in and integrals of (26) to predict a number of properties of the T_e profile in an H-mode pedestal.

Neglecting anomalous electron heat transport in the pedestal and using (22) with $\Gamma^{\rm an} = 0$ to obtain $-(d/d\rho)(V' \bar{D}_\eta \, n_e) = \dot{N}_e$, the electron heat flow Eq. (26) can be written as

$$-\,(V'\bar{D}_\eta \, n_e)\,\frac{3}{2}\,\frac{dT_e}{d\rho} = \hat{P}_e(\rho) \;-\; \frac{3}{2}\dot{N}_e T_e \,. \tag{34}$$

Electron heat lost via the fueling term is usually negligible compared to conductive electron heat flow through the pedestal: $(3/2)\dot{N}_e T_e/\hat{P}_e \sim (1.5)(2{\times}10^{21})(90)(1.6{\times}10^{-19})/(1.7{\times}10^6) \simeq 0.025$ at the separatrix in the 98889 DIII-D pedestal. Thus, neglecting the fueling term the electron temperature profile is determined to lowest order by

$$\boxed{-\,\frac{dT_e}{d\rho} \;=\; \frac{\hat{P}_e}{(3/2)(V'\bar{D}_\eta \, n_e)} \;\simeq\; \text{constant,}} \quad \text{electron temperature gradient in pedestal.} \tag{35}$$

Because local electron heating $Q_e^{\rm net}$ and convective heat flow are often negligible compared to electron heat flow through the pedestal, $\hat{P}_e(\rho) \simeq P_e(a)$ is often nearly constant (see Figs. 4a and 5a in Ref. [1]). Also, from (31) the combination of parameters $(V'\bar{D}_\eta \, n_e)$ is predicted to be approximately constant in H-mode pedestals. Hence, Eq. (35) predicts the electron temperature gradient is nearly constant in pedestals — until electron-temperature-gradient (ETG) microturbulence [32] causes significant anomalous electron heat transport in moving from the pedestal into the core plasma [1], as will be discussed at the end of this section.

The near constancy of $\hat{P}_e$ and $(V'\bar{D}_\eta \, n_e)$ also means that experimental validation tests of the prediction for $dT_e/d\rho$ in (35) can be made at any flux surface between about the mid-point of the pedestal and the separatrix. Evaluation at the separatrix would be simpler theoretically because the μ_e/ν_e trapped particle effects become negligible there and D_η is then determined solely by the Spitzer resistivity which just depends on $Z_{\rm eff}$ and $T_e^{-3/2}$ —

but obtaining experimental data at the separatrix for n_e, T_e and $Z_{\rm eff}$ can be problematic. Alternatively, (35) could be evaluated at the mid-point of the pedestal — but a very accurate evaluation of D_η is needed there. An estimate of the needed D_η can be bounded by the ONETWO-inferred value for a minimum and a maximum from Eqs. (1)–(8) [or Eq.(9) for 98889]. Thus, the most viable procedure is apparently to evaluate (35) at the mid-point ρ_n ($= 0.982a$ in 98889 [1]) of the n_e profile. Since ρ_n is often at a slightly larger radius than the T_e profile tanh fit symmetry point, evaluation there reduces possible complications from increasing ETG-induced anomalous transport in the top half of the T_e pedestal.

The electron temperature gradient prediction in (35) could in principle be integrated using a D_η formula such as that specified in (9). However, since $dT_e/d\rho$ is predicted to be nearly constant over most of the pedestal, the predicted T_e profile in the pedestal is simply $T_e(\rho) \simeq T_e(a) + (a - \rho)[-dT_e/d\rho]_{\rho_n}$ which should apply from the separatrix in to near the top of the n_e pedestal. Using this T_e profile in formulas for D_η given in (1)–(8) [or (9) for 98889] allows one, in principle, to determine the n_e profile and its gradient from (31) or (29).

In order to compare the electron temperature gradient prediction in (35) to experimental data it is useful to convert it to a prediction for the electron temperature gradient scale length. Dividing (35) by T_e and multiplying it by the minor radius a yields a prediction for the normalized electron temperature gradient scale length at the density tanh fit point ρ_n:

$$\boxed{\left.\frac{L_{Te}}{a}\right|_{\rho_n} \equiv \left[-\frac{a}{T_e}\frac{dT_e}{d\rho}\right]^{-1}_{\rho_n} = \frac{(3/2)\,[V'\bar{D}_\eta n_e]_{\rho_n} T_e(\rho_n)}{a\,\hat{P}_e(\rho_n)},} \quad T_e \text{ gradient scale length.} \qquad (36)$$

Very roughly scaling $\hat{P}_e \sim \overline{n_e T_e}\, V(a^2/\bar{D}_\eta)$, this yields $L_{Te}/a \sim [n_e T_e]_{\rho_n} / \overline{n_e T_e}$, which is the small ratio of the pedestal mid-point electron pressure to the volume-average electron pressure ($\lesssim 10^{-2}$). At the density mid-point ρ_n in the 98889 DIII-D pedestal $T_e \simeq 360$ eV, $V' \simeq 42.6$ m^2, $n_e \simeq 1.65\times 10^{19}$ m^{-3} and the conductive $\hat{P}_e \simeq 1.7 \times 10^6$ Watts, which yields $[L_{Te}/a]_{\rho_n} \simeq 0.046\,\bar{D}_\eta$. In the 98889 DIII-D pedestal $\bar{D}_\eta \simeq 1.6\,D_\eta$ so $D_\eta(\rho_n)$ estimates of 0.2 (Spitzer), 0.45 (ONETWO) and 0.89 m^2/s [from Eq. (9)] yield predictions of $[L_{Te}/a]_{\rho_n} \simeq$ 0.015, 0.033 and 0.066. These values range from 0.75 to 3.3 times the measured $[L_{Te}/a]_{\rho_n} \simeq$ 0.02 (see Fig. 9b in [1]). The variability of the predictions illustrates the need for accurate evaluations (or bounding) of D_η and hence of the neoclassical parallel resistivity $\eta_\parallel^{\rm nc}$.

The effect of ETG-induced anomalous electron heat transport will be considered next. Near the pedestal top (at $\rho_t \simeq 0.96a$ for 98889) the combination of parameters $V'\bar{D}_\eta n_e$ is nearly constant (on the L_{Te} scale length) not because of (31) but because each of its parameters is individually nearly constant there. Assuming M is nearly constant there and $\Upsilon_e^{\rm an}(a) \simeq 0$, the electron energy flow equation in (26) can be simplified near $\rho_{\rm t}$ to

$$-\,V'\bar{D}_\eta n_e \,\frac{3}{2}\,\frac{dT_e}{d\rho} \,+\, \frac{1}{M+1}\,V'\Upsilon_e^{\rm an} = \hat{P}_e(\rho). \tag{37}$$

A formula for the anomalous electron heat flux induced by ETG-driven microturbulence is needed to proceed further. While no such expression is available to the author's knowledge, one can be constructed from results of the gyrokinetic-based numerical simulations of ETG-induced microturbulence by Jenko et al. [33] at the top of an Asdex Upgrade (AUG) H-mode pedestal [34] that has parameters similar to the 98889 DIII-D pedestal [1]. From Table I in [33], at this AUG pedestal's top $\eta_e \equiv d\ln T_e/d\ln n_e = L_n/L_{Te} = \omega_{Te}/\omega_n = 24.8/11.2 \simeq 2.2$; similarly, Fig. 10b in Ref. [1] indicates that near $\rho_{\rm t}$ in the 98889 pedestal $\eta_e \gtrsim 2$. Thus, from Fig. 7 in Ref. [33] since $\eta_e \gg \eta_{e,{\rm crit}} \simeq 1.2$ [32], one is apparently well beyond the ETG threshold regime and in a nearly asymptotic ETG regime where $\chi_e^{\rm ETG}$ no longer increases with η_e. Hence, one can approximate the ETG-induced anomalous electron heat transport by a standard Fourier heat flux law in the form $\Upsilon_e^{\rm ETG} = -\,n_e\langle|\boldsymbol{\nabla}\rho|^2\rangle\chi_e^{\rm ETG} dT_e/d\rho$. For the AUG parameters the ETG simulations [33] found $\chi_e^{\rm ETG} \simeq 0.83$ m^2/s, which is a factor of about two larger than the $\chi_e \sim 0.4$ m^2/s obtained from interpretive transport analysis [34]. Using parameters for this AUG pedestal of $T_e \simeq 0.69$ keV, $L_{Te} = a/\omega_{Te} \simeq 0.65/24.8 \simeq 0.026$ m and $B_{\rm t0} \simeq 2.4$ T, in terms of an anticipated electron gyroBohm diffusivity $\chi_e^{\rm gB} \equiv (\sqrt{T_e/m_e}/\omega_{ce}L_{Te})(T_e/eB_{\rm t0})$ this ETG-induced electron heat diffusivity is $\chi_e^{\rm ETG} \simeq 3\,\chi_e^{\rm gB}$; the numerical factor would be $\simeq 1.4$ if the interpretive analysis χ_e were used. Thus, one infers

$$\chi_e^{\rm ETG} \;\simeq\; f_\#\chi_e^{\rm gB} \;\equiv\; f_\#\,\frac{\rho_e}{L_{Te}}\frac{T_e}{eB_{\rm t0}} \;\simeq\; 0.075\,f_\#\,\frac{[T_e({\rm keV})]^{3/2}}{L_{Te}({\rm m})\,B_{\rm t0}^2({\rm T})^2}\;\;{\rm m}^2/{\rm s}, \tag{38}$$

in which from the AUG pedestal transport analysis [34] and simulations [33] $f_\# \simeq 1.4$–3.

The effect of the resultant anomalous electron heat flux $\Upsilon_e^{\rm ETG}$ in (37) will now be considered. Since $\chi_e^{\rm ETG} \propto T_e^{3/2}$ whereas $\chi_e^{\rm pc} \simeq (3/2)MD_\eta \propto T_e^{1/2}$, at the pedestal top where $M \gg 1$ ($M \simeq 2.6$ at $\rho_{\rm t} \simeq 0.96a$ in 98889), the ETG-induced anomalous electron heat transport will begin to "saturate" the increase in T_e and cause a change in the electron temperature gradient when $\chi_e^{\rm ETG} \gtrsim \chi_e^{\rm pc} \simeq (3/2)\,(\sqrt{2T_e/m_e}/\pi R_0 q)\,\delta_e^2$. Equating these electron heat diffusivities and for simplicity neglecting the $(\bar{a}^2/a^2)\langle|\boldsymbol{\nabla}\rho|^2\rangle \simeq 2/1.6 \simeq 1.25$ factor yields an approximate prediction for the electron pressure $p_e \equiv n_eT_e$ at the pedestal top:

$$\boxed{\beta_e^{\rm ped} \;\equiv\; \frac{n_e^{\rm ped}T_e^{\rm ped}}{B_{\rm t0}^2/2\mu_0} \;\sim\; \frac{3\sqrt{2}}{\pi f_\#}\,\frac{\eta_\parallel^{\rm nc}}{\eta_0}\,\frac{L_{Te}}{R_0 q},} \quad \text{electron pressure at top of the pedestal.} \tag{39}$$

This is just a rough criterion for where T_e begins to saturate with decreasing ρ near the T_e pedestal top. However, it may reflect the typical experimental pedestal data [1] in which T_e often increases slowly as one moves from the pedestal into the core, which makes it difficult to clearly identify a specific position where the "top" of the T_e pedestal occurs. A more approximate pedestal height prediction than (39) that used a more generic gyroBohm scaling for $\chi_e^{\rm an}$ was given previously in the last paragraph of Section 6 in Ref. [10].

At the top ($\rho_{\rm t} \simeq 0.96a$) of the 98889 pedestal $L_{Te} \simeq 0.05a \simeq 0.0385$ m and $q \simeq 4.4$, so Eq. (39) predicts $\beta_e^{\rm ped} \simeq (0.007/f_\#)(\eta_\parallel^{\rm nc}/\eta_0)$. Since M and $\beta_e^{\rm ped}$ are about halved because $\ell_{\max} \simeq \lambda_e$ here, this is about twice the ONETWO value $\beta_e(0.96a) \simeq 0.002$ for the range of values for $\eta_\parallel^{\rm nc}/\eta_0$ (~ 1.5 [ONETWO] to 3 [Eq. (9)]) and $f_\#$ (~ 1.4–3). The electron heat diffusivities at $\rho_{\rm t} \simeq 0.96a$ in the 98889 DIII-D pedestal are $\chi_e^{\rm ETG} \simeq \chi_e^{\rm pc} \sim 0.3$–$0.65$ m^2/s, which are of order the spatially varying interpretive transport results there (see Fig. 8a of Ref. [1]).

VI. ION TEMPERATURE PROFILE

Ion heat transport in the pedestal is apparently [1] a complicated mix of neoclassical plus paleoclassical throughout, possible kinetic effects in the pedestal's bottom half and a transition to ion-temperature-gradient (ITG) driven anomalous transport in the core. Neglecting $\Upsilon_i^{\rm an}$ and kinetic effects in (28), the pedestal ion temperature gradient is determined by

$$-\frac{dT_i}{d\rho} \simeq \frac{P_i(\rho)/V'}{(3/2)n_i\bar{D}_\eta + n_i\chi_i^{\rm nc}}, \quad \text{ion temperature gradient in pedestal.} \qquad (40)$$

The ion relation deduced from (31) indicates $n_i\bar{D}_\eta = (n_i/n_e)\, n_e\bar{D}_\eta \simeq (n_i/n_e)\times$ constant. Assuming n_i/n_e is nearly constant in the pedestal so $n_i\bar{D}_\eta \sim$ constant and that $\chi_i^{\rm nc}$ does not vary too much in the pedestal ($\lesssim$ factor of 2 in 98889 — see Fig. 11 in [1]), the predicted ion temperature gradient may also be approximately constant in the pedestal. Kinetic effects due to ion banana drift orbits could modify $dT_i/d\rho$ very near the separatrix and even render T_i an inappropriate measure of the ion energy distribution there.

Dividing (40) by T_i/a and evaluating it at ρ_n yields the prediction

$$\boxed{\left.\frac{L_{Ti}}{a}\right|_{\rho_n} \equiv \left[-\frac{a}{T_i}\frac{dT_i}{d\rho}\right]^{-1}_{\rho_n} \simeq \frac{[(3/2)\bar{D}_\eta + \chi_i^{\rm nc}]_{\rho_n} n_i(\rho_n)\, T_i(\rho_n)}{a\, P_i(\rho_n)/V'},} \quad T_i \text{ gradient scale length.} \quad (41)$$

At the density mid-point point ρ_n in the 98889 DIII-D pedestal $n_i \simeq 0.634\times 1.65\times 10^{19}$ m^{-3}, $T_i \simeq 550$ eV, $V' \simeq 42.6$ m^2 and the conductive $P_i \simeq 0.65 \times 10^6$ Watts, so Eq. (41) yields

$[L_{Ti}/a]_{\rho_n} \simeq 0.078\,(1.5\bar{D}_\eta + \chi_i^{\rm nc})$. In the 98889 pedestal $\chi_i^{\rm nc}(\rho_n) \simeq 0.45$ m^2/s and $\bar{D}_\eta \simeq 1.6\,D_\eta$ so $D_\eta(\rho_n)$ estimates of 0.2 (Spitzer), 0.45 (ONETWO) and 0.89 m^2/s [from Eq. (9)] yield predictions of $[L_{Ti}/a]_{\rho_n} \simeq 0.073$, 0.12 and 0.20. These values are somewhat larger than the measured $[L_{Ti}/a]_{\rho_n} \simeq 0.06$ (see Fig. 9b in [1]); their variability illustrates the need for very accurate evaluations (or bounding) of both D_η and the neoclassical ion heat diffusivity $\chi_i^{\rm nc}$. It was already noted in [1] that the neoclassical ion heat diffusivity at the pedestal mid-point may be up to a factor of 3 too large there. It also appears that the paleoclassical prediction of $\chi_i^{\rm pc} \simeq (3/2)D_\eta$ is too large in the pedestal, by a factor of 2 (ONETWO) to ~ 3 [Eq. (9)].

Determination of the "top" of the T_i pedestal is even more problemmatic than for T_e — see discussion after Eq. (39) above. This is because firstly the pedestal ion temperature profile often does not have a very well defined point where its gradient changes significantly (cf., Fig. 3a in [1]), which would be indicative of a transition in the dominant ion heat transport processes. Secondly, unlike with ETG-driven anomalous transport, ITG-driven anomalous transport is likely to be at or near threshold conditions in which case a simple Fourier heat flux law is not appropriate. A rough estimate of the T_i pedestal height might be able to be developed from (28) if and when a formula for the ITG-microturbulence-induced ion heat flux $\Upsilon_i^{\rm an}$ near the $dT_i/d\rho|_{\rm crit}$ threshold becomes available.

VII. TOROIDAL FLOW PROFILE AND RADIAL ELECTRIC FIELD

Plasma flows are not so well characterized or understood in H-mode pedestals. The poloidal ion flows $V_{\rm p}$ should be determined [13, 14] by ion neoclassical processes [15, 17] since anomalous plasma transport is being neglected in this pedestal structure model. For deuterons it is predicted to be $V_{{\rm p}i} \simeq (c_{\rm p}/q_i B_{\rm t0})(dT_i/d\rho)$, in which $c_{\rm p} \equiv k_i \sim 1$ ($k_i \simeq 1.17$ for $\nu_{*i} \ll 1$, $\sqrt{\epsilon} \ll 1$ and $Z_{\rm eff} = 1$) is a factor determined by neoclassical theory [17] which can be evaluated using NCLASS [20]. Early measurements [35] indicated dominant ion (helium) poloidal flows larger than neoclassical theory predicted; however, charge-exchange frictional drag may have been important there [36]. Recent measurements of poloidal impurity ion flows in C-Mod [37, 38], MAST [39], NSTX [40] and DIII-D [41] pedestals are of order neoclassical predictions but often have large error bars that prevent definitive conclusions.

Recently a comprehensive equation has been developed [13, 14, 42] for the total plasma toroidal rotation frequency $\Omega_{\rm t} \equiv L_{\rm t}/(m_i n_i \langle R^2\rangle) \simeq V_{\rm t}/R$ where $L_{\rm t} \equiv m_i n_i \langle \mathbf{e}_\zeta \cdot \mathbf{V}\rangle \simeq$

$\sum_{\rm ions} m_s n_s \langle \mathbf{e}_\zeta \cdot \mathbf{V}_s \rangle$ is the plasma toroidal angular momentum density in which $\mathbf{e}_\zeta \equiv R\, \hat{\mathbf{e}}_\zeta$ is the toroidal angular moment vector. That is, the toroidal flows of the dominant hydrogenic ion and impurity species are added together here and $\Omega_{\rm t}$ will represent the total plasma toroidal angular rotation frequency. The separate toroidal flows of the dominant hydrogen ions and the impurities can be obtained from NCLASS calculations using the determined $\Omega_{\rm t}$ (radial electric field [13, 42]).

Neglecting resonant ($q = m/n$) three-dimensional (3D) magnetic field effects, in a transport quasi-equilibrium the toroidal plasma rotation equation can be simplified to [13, 14, 42]

$$0 \;=\; -\,\langle \mathbf{e}_\zeta \cdot \boldsymbol{\nabla} \cdot \boldsymbol{\pi}_{i\|}^{\rm 3D} \rangle - \,\langle \mathbf{e}_\zeta \cdot \boldsymbol{\nabla} \cdot \boldsymbol{\pi}_{i\perp} \rangle - \,\frac{1}{V'}\,\frac{d}{d\rho}\,(V'\Pi_{i\rho\zeta}) + \,\langle \mathbf{e}_\zeta \cdot \mathbf{S}_m \rangle. \qquad (42)$$

Here, $\boldsymbol{\pi}_{i\|}^{3D}$ is the parallel ion viscous stress induced by non-resonant 3D field effects, $\boldsymbol{\pi}_{i\perp}$ is the perpendicular ion stress due to collision-induced classical, neoclassical and paleolassical processes, $\Pi_{i\rho\zeta}$ is the ρ, ζ component of the microturbulence-induced Reynolds ion stress tensor and $\langle \mathbf{e}_\zeta \cdot \mathbf{S}_m \rangle$ is the torque density applied to the plasma at this flux surface by the net momentum sources and sinks $\mathbf{S}_m$.

Neglecting effects due to non-resonant (and resonant) 3D magnetic fields and fluctuation-induced Reynolds stress which are all usually small in H-mode pedestals, this equation becomes simply $\langle \mathbf{e}_\zeta \cdot \boldsymbol{\nabla} \cdot \boldsymbol{\pi}_{i\perp} \rangle \simeq \langle \mathbf{e}_\zeta \cdot \mathbf{S}_m \rangle$. There is usually no significant momentum source in H-mode pedestals. However, charge-exchange losses can be significant there; they will be represented by $\langle \mathbf{e}_\zeta \cdot \mathbf{S}_m \rangle \,\simeq\, -\,\nu_{\rm cx} L_{\rm t} \,=\, -\,\sum_{\rm ions} m_s n_s \nu_{{\rm cx}s} \langle \mathbf{e}_\zeta \cdot \mathbf{V}_s \rangle$. The collision-induced perpendicular ion stress is dominated by its paleoclassical component [13]. Thus, using the paleoclassical toroidal torque given by Eq. (103) in [13], the lowest order equation becomes

$$-\,\frac{1}{V'}\,\frac{d^2}{d\rho^2}\,[V'\,\bar{D}_\eta\,L_{\rm t}] \;\simeq\; -\,\nu_{\rm cx} L_{\rm t}, \quad \text{pedestal plasma toroidal rotation equation.} \qquad (43)$$

Multiplying by V' and integrating over ρ yields an equation analogous to the density flow equation in (22) and (23) or the electron energy flow equation in (26) and (27):

$$-\left[\frac{d}{d\rho}(V'\bar{D}_\eta L_{\rm t})\right]_\rho \simeq\, V'\Pi_{\rm t}, \quad V'\Pi_{\rm t} \,\equiv\, -\left[\frac{d}{d\rho}(V'\bar{D}_\eta L_{\rm t})\right]_a + \int_\rho^a V'(\hat{\rho})\,d\hat{\rho}\;\nu_{\rm cx}(\hat{\rho}) L_t(\hat{\rho}). \qquad (44)$$

Here, $V'\Pi_{\rm t} \,=\, V'(\rho)\,\Pi_{\rm t}(\rho)$ is the radial flow of toroidal angular momentum (N·m/s), which is comprised of the flow of L_t through the separatrix plus an integral correction due to charge-exchange losses between the ρ surface and the separatrix.

In low density pedestals with modest NB torque applied to the plasma and small charge exchange losses (e.g., from dominant X-point fueling rather than the main chamber walls), the $V'\Pi_i$ flow term can be neglected. Then, (44) becomes simply $[(d/d\rho)(V'\bar{D}_\eta L_{\rm t})]_\rho \simeq 0$. Analagous to the density relation presented in (31), this equation indicates that $V'\bar{D}_\eta L_{\rm t} = (n_i V'\bar{D}_\eta)\, m_i \langle R^2\rangle\, \Omega_{\rm t} \simeq$ constant. Neglecting the small variations of V' and $\langle R^2\rangle$ in the pedestal and using the relation in (31) assuming $n_i/n_e \simeq$ constant, this last relation yields

$$\boxed{\Omega_{\rm t}(\rho) \simeq \text{constant} \quad \Longrightarrow \quad \Omega_{\rm t}(\rho) \simeq \Omega_{\rm t}(a), \quad \text{within the pedestal.}} \tag{45}$$

Thus, to lowest order the total toroidal plasma flow $V_{\rm t} \simeq R\,\Omega_{\rm t}$ is predicted to be roughly constant within the pedestal — at its value on the separatrix. Hence, in this model the separatrix $V_{\rm t}$ boundary condition (usually > 0) is carried inward from the separatrix to the top of the pedestal. Toroidal flow of the carbon species has been found to be nearly constant in the pedestal region in both the 98889 and a similar 119436 DIII-D pedestal [27, 43], in qualitative agreement with the prediction in (45).

Possible effects of radial flow of $L_{\rm t}$ through the pedestal and charge-exchange momentum losses can be quantified by integrating (44) from ρ to the separatrix at a to yield

$$V'(\rho)\,\bar{D}_\eta(\rho)\,L_{\rm t}(\rho) \;\simeq\; V'(a)\,\bar{D}_\eta(a)\,L_{\rm t}(a) \;+\int_\rho^a V'(\hat\rho)\,d\hat\rho\;\Pi_{\rm t}(\hat\rho). \tag{46}$$

In high density pedestals with significant charge-exchange momentum losses (e.g., mainly from main chamber wall recycling) and modest applied plasma torques (e.g., in ECH H-mode plasmas or with weak NBI), one can apparently have a situation where radial flow of $L_{\rm t}$ through the pedestal is small and momentum loss via charge exchange dominates. Then, when Ω_t is positive (co-current directed) on the separatrix, charge-exchange momentum losses in the pedestal require a radial flow of $L_{\rm t}$ inward (i.e., < 0) across the separatrix. This situation can be modeled by representing the charge exchange momentum loss by $\nu_{\rm cx}(\rho)L_{\rm t}(\rho) \simeq \nu_{\rm cx}(a)L_{\rm t}(a)\,e^{-(a-\rho)/\lambda_n}$. Then, assuming the pedestal is wider than the neutral penetration depth λ_n, the inward momentum flow across the separatrix estimated by setting $V'\Pi_{\rm t}$ from (44) to zero for $a-\rho \gg \lambda_n$ yields $V'(a)\,\Pi_{\rm t}(a) \equiv -[(d/d\rho)(V'\bar{D}_\eta L_{\rm t})]_a = -\lambda_n\nu_{\rm cx}(a)L_{\rm t}(a)$. Using this result in (46) and again assuming $\lambda_n \gg a-\rho$ then yields

$$\Omega_{\rm t}(\rho) \;\simeq\; \Omega_{\rm t}(a)\,[1-(a-\rho)\lambda_n\nu_{\rm cx}(a)/\bar{D}_\eta(a)], \quad \text{for large cx losses in pedestal.} \tag{47}$$

Thus, charge-exchange losses cause $\Omega_{\rm t}$ to decrease linearly with distance in from its value on the separatrix, i.e., they cause $d\Omega_{\rm t}/d\rho = [\lambda_n\nu_{\rm cx}(a)/\bar{D}_\eta(a)]\;\Omega_{\rm t}(a) > 0$ in the pedestal. This is qualitatively consistent with recent results in DIII-D ECH H-mode plasmas [44] and experimental exploration of multiple ion species $dV_{{\rm t}s}/d\rho > 0$ cases in ASDEX Upgrade [45].

At the top of the n_e, T_e pedestal, radial transport of plasma toroidal angular momentum presumably transitions to being determined by ITG-induced microturbulence which produces a significant ion Reynolds stress component $\Pi_{i\rho\zeta}$ in the core plasma. In the absence of a specific prediction for $\Pi_{i\rho\zeta}$ it is difficult to develop a prediction for how $d\Omega_{\rm t}/d\rho$ might change and precisely where. Assuming $\Pi_{i\rho\zeta} \sim -\,\chi_{\rm t}\, dL_{\rm t}/d\rho$, at the radius where ITG-induced microturbulence causes $\chi_{\rm t} \gtrsim \bar{D}_\eta$, integration of (42) from $\rho = 0$ to $\rho \sim \rho_{\rm t}$ indicates the rotation profile gradient should change to some finite value: $d\Omega_{\rm t}/d\rho < 0$ for co-current core momentum input $\langle \mathbf{e}_\zeta \boldsymbol{\cdot} \mathbf{S}_m \rangle > 0$ or $d\Omega_{\rm t}/d\rho > 0$ for counter-current core momentum input.

On the outboard mid-plane of a tokamak the pedestal region is close to the toroidal magnetic field coils and possible magnetic field errors (and the I-coils in DIII-D). Thus, the 3D magnetic field components are largest there and can cause additional toroidal torques on the pedestal plasma. Ions trapped in the toroidal magnetic field ripple can drift radially out of the plasma and hence induce a radial "return current" [46] in the plasma. When this radially inward (negative) plasma return current is crossed with the poloidal magnetic field, it induces a toroidal torque on the plasma in the counter-current direction, which reduces the toroidal rotation in the pedestal. In addition, neoclassical toroidal viscosity (NTV) effects due to non-resonant 3D magnetic fields (from error fields, I-coils and toroidal field ripple) cause a FSA toroidal torque on the plasma ions of the generic form [13, 42]

$$-\,\langle \mathbf{e}_\zeta \boldsymbol{\cdot} \boldsymbol{\nabla} \boldsymbol{\cdot} \boldsymbol{\pi}_{i\|}^{\rm 3D} \rangle \,\simeq\, -\,m_i n_i\, \mu_\| \left(\frac{\delta B^{\rm 3D}}{B_{\rm t0}} \right)^2 \langle R^2 \rangle \; (\Omega_{\rm t} - \Omega_*), \quad \Omega_* \simeq \frac{c_{\rm p} + c_{\rm t}}{q_i R B_p} \, \frac{dT_i}{d\rho}. \qquad (48)$$

The order unity numerical coefficient $c_{\rm p} + c_{\rm t}$ is determined by the specific 3D radial banana center or ripple-induced drift process and ion collisionality regime involved [42]. The NTV effects due to non-resonant field components damp the plasma toroidal rotation toward a counter-current-directed $\Omega_* \propto dT_i/d\rho < 0$. Thus, both direct ripple-trapped ion losses and NTV effects decrease a positive plasma toroidal rotation frequency $\Omega_{\rm t}$ in the pedestal. They can even cause $\Omega_{\rm t}$ to become negative and approach $\Omega_* < 0$ for very large NTV effects [47]. Experiments in JET observed a decreasing toroidal flow Mach number in the edge (and thus throughout the) plasma as the magnetic field ripple was increased (see Fig. 6 in [30]), in qualitative agreement with these predictions.

Determining the $\Omega_{\rm t}$ profile in the pedestal when paleoclassical, charge-exchange momentum losses and ripple loss plus NTV effects are all present requires detailed numerical solutions of (42) for $L_{\rm t}$. Once the toroidal plasma rotation frequency $\Omega_t \equiv L_{\rm t}/m_i n_i \langle R^2 \rangle$ is known, the radial electric field is determined from radial ion force balance which yields [13]

$$E_\rho \equiv -\,\hat{\mathbf{e}}_\rho \boldsymbol{\cdot} \boldsymbol{\nabla}\Phi_0 = |\boldsymbol{\nabla}\rho| \left(\Omega_{\rm t}\psi_{\rm p}' + \frac{1}{n_i q_i}\frac{dp_i}{d\rho} - \frac{c_{\rm p}}{q_i}\frac{dT_i}{d\rho} \right), \tag{49}$$

in which $\hat{\mathbf{e}}_\rho \equiv \boldsymbol{\nabla}\rho/|\boldsymbol{\nabla}\rho|$ is the unit vector in the radial direction and $\psi_{\rm p}' = RB_{\rm p}$. In many H-mode pedestals the toroidal rotation effects are small [because from sheath effects in the scrape-off-layer (SOL) outside the separatrix $\Omega_{\rm t} \sim T_e(a)/L_{\rm SOL} \equiv T_e(a)|d\ln T_e/d\rho|_{\rm SOL}$ is usually less than $T_i(\rho_n)/L_{n_e} \equiv T_i(\rho_n)|d\ln n_e/d\rho|_{\rho_n}|$) and the poloidal ion flow effects are small (because $L_{Ti} \gg L_{n_e}$). Thus, in pedestals the radial electric field is typically nearly equal to to ion pressure gradient: $E_\rho \simeq (1/n_i q_i)(dp_i/d\rho)$. The degree to which these two dominant terms in (49) do not cancel produces the net plasma toroidal flow in the pedestal; this observation provides another interpretation of the key processes involved in the new interpretive density analysis procedure [27]. The resultant net toroidal rotation provides the small net thermodynamic drive for density transport in the pedestal [1].

VIII. DIMENSIONLESS VARIABLE SCALINGS

The natural dimensionless variables of the paleoclassical transport model are very different from the usually used ones — see Eqs. (63) and (64) and the surrounding discussion in Ref. [6]. The key parameter in the paleoclassical model is the magnetic field diffusivity D_η. It can be scaled to the reference diffusivity $\eta_0/\mu_0 = (m_e\nu_e)/(n_e e^2\mu_0) = \nu_e\delta_e^2$ in which $\delta_e \equiv c/\omega_p \propto 1/\sqrt{n_e}$ is the electromagnetic skin depth defined in (14). Noting from Eqs. (3)–(6) that $\eta_\parallel^{\rm nc}/\eta_0 = f^{\rm nc}(Z_{\rm eff}, \epsilon, \nu_{*e})$, in terms of δ_e and the natural paleoclassical variables the key electron temperature gradient scale length prediction in (35) can be written as

$$L_{Te} \simeq \delta_e\, f^{\rm nc}(Z_{\rm eff}, \epsilon, \nu_{*e})\, \frac{\nu_e\, \delta_e\, n_e T_e}{\hat{P}_e/V'}. \tag{50}$$

This form illustrates that in this paleoclassical-based pedestal structure model the electron temperature gradient scale length at the mid-point of the pedestal is a multiple of δ_e. The multiple is the ratio of the energy density $p_e \equiv n_e T_e$ (J/m^3) transported over the distance δ_e in an electron collision time $1/\nu_e$ divided by the energy flow per second (Watts) across the area $S \equiv \langle|\boldsymbol{\nabla}\rho|\rangle V' \simeq 1.2\, V'$ [1] of the ρ flux surface.

In this paleoclassical-based model the plasma gradients and profiles within the pedestal do not depend explicitly on the usual dimensionless parameters of normalized ion gyroradius

$\varrho_* \equiv \varrho_i/a$ or relative pressure $\beta \equiv P/(B^2/2\mu_0)$. (However, they could depend on β through their dependence on $\hat{P}_e/n_eT_e$.) The lack of a dependence on ϱ_* agrees with JET/DIII-D comparison experiments [30] which found essentially no dependence of the pedestal width or profile on ρ_* as it was varied by a factor of four. However, the height of the pedestal is determined in this pedestal structure model by the transition from paleoclassical to electron-gyroBohm-scaling ETG-driven anomalous electron heat transport. This introduces a β_e dependence into the initial, transport-limited pedestal height. To see this note that $\chi_e^{\rm gB} \sim (v_{Te}/L_{Te})\,\varrho_e^2$ (here the electron thermal speed is $v_{Te} \equiv \sqrt{T_e/m_e}$ and $\varrho_e \equiv v_{Te}/\omega_{ce}$ is the electron gyroradius) whereas the paleoclassical $\chi_e^{\rm pc} \simeq (3/2)M\nu_e\delta_e^2 \sim (v_{Te}/\pi R_0 q)\,\delta_e^2$. When these different transport mechanisms are equated to determine T_e at the top of the pedestal, they naturally yield $\varrho_e^2/\delta_e^2 = \beta_e^{\rm ped} \sim (v_{Te}/R_0 q)/(v_{Te}/L_{Te}) \sim (L_{Te}/R_0 q)$, as obtained in (39).

IX. DISCUSSION

The preceding sections have developed pedestal structure properties for the electron density profile (31), density fueling effects (32), maximum height of the density pedestal (33), electron temperature gradient (35), T_e gradient scale length (36), transport-limited height of the electron pressure pedestal (39), T_i gradient scale length (41), plasma toroidal rotation profile (45) and charge exchange momentum loss effects on $\Omega_{\rm t}$ (47). The predictions have all been shown to be within a factor of about two of the properties of the 98889 DIII-D pedestal [1]. Thus, all aspects of this new paleoclassical-based pedestal structure model seem to be in reasonable accord with experimental results — at least for the relatively low density, high T_e and hence low collisionality 98889 DIII-D pedestal. And maybe one should not expect better than factor of two agreement given the experimental error bars and the fact that paleoclassical theory has undetermined order unity coefficients [5, 6].

In all the evaluations for the 98889 pedestal it seems like D_η is systematically too large — by about a factor of two. There are two main possible reasons for this. First, in the 98889 data set the fractional density of fully ionized $C_{\rm VI}$ is taken to be a constant ($n_{C_{\rm VI}}/n_e \simeq 0.061$) and hence $Z_{\rm eff} = 2.83$ is also constant throughout the pedestal. However, $Z_{\rm eff}$ may decrease with increasing ρ in the pedestal as T_e decreases because carbon may be in lower ionization states there. While the $Z_{\rm eff}$ dependence in Eqs. (2)–(9) is complicated, D_η decreases approximately linearly with $Z_{\rm eff}$. Second, as discussed in connection with quasi-symmetric stellarators in Ref. [7], the paleoclassical model D_η should be multiplied by the fraction of local FSA poloidal magnetic field caused by current flowing in the plasma. Thus, the D_η in (1) should really be multiplied by a factor $(\partial\psi_J/\partial\rho)\,/\,(\partial\psi_{\rm p}/\partial\rho) \leq 1$ in which ψ_J is the poloidal magnetic flux induced by the current density on that surface and $\psi_{\rm p}$ is the total poloidal flux there. Since at the separatrix X-point the current-induced poloidal magnetic field just balances that produced by the divertor coils, just inside the separatrix the D_η might be reduced by up to a factor of two, but likely much less than that.

The pedestal height prediction in (39) can also be written in a form similar to a prediction for the normalized pedestal electron pressure gradient:

$$R_0 q\,\frac{\beta_e^{\rm ped}}{L_{Te}} \simeq \frac{3\sqrt{2}}{\pi f_\#}\,\frac{\eta_\parallel^{\rm nc}}{\eta_0} \sim 1. \tag{51}$$

This form resembles the ideal MHD instability parameters for high n ballooning modes [48] $\alpha \equiv -\,R_0 q^2 d\beta/d\rho$ and peeling-ballooning instabilities [49]. However, its derivation is based on entirely different physics (balance of paleoclassical and ETG-induced anomalous transport, rather than ideal MHD instability) and its implications are entirely different.

The relation in (51) indicates the electron pressure gradient a pedestal will attain in transport quasi-equilibrium. If, after an L-H transition, this is less than the peeling-ballooning instability criterion [49] for triggering an ELM, the pedestal should evolve to this state (in $\tau \sim (2L_{Te})^2/\bar{D}_\eta \sim [(0.04)(0.77)]^2/[(1.6)(0.45)] \sim$ ms for 98889 parameters). However, its top would slowly evolve further as the core plasma continues to evolve on the global plasma energy confinement time scale ($\tau_E \sim 150$ ms in 98889 [1]). In this pedestal structure model the n_e and T_e profiles within the pedestal should remain nearly constant on this longer time scale. However, the top of the pedestal could move further inward as the core electron temperature increases which reduces the T_e curvature and gradient (and hence ETG-drive) at the "top" of the pedestal. This continuing growth and inward spreading of the top of the pedestal T_e [and n_e via (31)], which is observed experimentally [26], would eventually cause the peeling-ballooning instability boundary to be exceeded and precipitate a Type I ELM. Presumably the same scenario would transpire between repetitive Type I ELMs. On the other hand, if just after an L-H transition the effective conductive electron power $\hat{P}_e$ flowing through the pedestal were larger, it would require a larger T_e gradient to carry the needed electron heat flow through the pedestal. The resultant larger pedestal electron pressure gradient could exceed the peeling-ballooning instability criterion before the transport quasi-equilibrium discussed here could be established and precipitate more frequent Type I ELMs with linearly increasing T_e between ELMs. Or perhaps if the high-n ballooning limit were exceeded, particularly in the bottom half of the pedestal, Type II ELMs might occur.

The addition of large enough anomalous electron density and temperature transport in the pedestal causes the plasma to revert to an L-mode state. In this pedestal structure model this will occur when the anomalous density and heat fluxes exceed the corresponding paleoclassical fluxes in (22) and (26). For the most likely types of microturbulence-induced anomalous transport caused by drift-waves or fluid-like modes (see for example [50] and references cited therein), the effective anomalous diffusivities for electron density and heat transport are comparable: $D^{\rm an} \sim \chi_e^{\rm an}$, physically because of the simultaneous advection of density and temperature by the $\tilde{\mathbf{E}}\times\mathbf{B}$ flow fluctuations. In contrast, the effective paleoclassical diffusivities are quite different in the pedestal: $D_{\rm eff}^{\rm pc} \sim f_D D_\eta \ll D_\eta$, in which $f_D \ll 1$ (in 98889 [1] $f_D \sim 0.1$) represents the degree to which the diffusive density flux is cancelled by the pinch flow in the paleoclassical transport model, whereas $\chi_e^{\rm pc} \simeq (3/2)(M+1)D_\eta \gg D_\eta$. Thus, a small amount of additional transport induced by microturbulence or some other process can in principle change density transport in the pedestal without significantly influencing the electron heat transport there. Specifically, if $D^{\rm an} \gtrsim f_D D_\eta \sim 0.1\, D_\eta$ but $\chi_e^{\rm an} \ll \chi_e^{\rm pc}$, the $n_e \bar{D}_\eta \sim$ constant relation in (31) would no longer hold, but the $dT_e/d\rho$ prediction in (35) would still be valid in the pedestal, albeit with $dT_e/d\rho$ varying spatially in the pedestal. Thus, if a small "controlled" additional pedestal density transport flux is added it could reduce $|dn_e/d\rho|$ but not change $dT_e/d\rho$ much and thereby prevent ELMs; this may explain the ELM-free EHO-induced Quiescent H- (QH-) modes in DIII-D [51] and EDA H-modes in C-Mod [52]. Also, this scenario could provide an explanation for the I-mode regime discovered recently in C-Mod [52] in which the density profile had no transport barrier (i.e., was L-mode like) but the T_e profile exhibited a normal H-mode-like transport barrier. For large anomalous transport that causes $\chi_e^{\rm an} > \chi_e^{\rm pc}$ and $D^{\rm an} > D_\eta$ the plasma should revert to an L-mode state with no transport barriers or pedestals in either n_e or T_e.

This pedestal structure model also provides a new interpretation of the ELM-free state achieved via resonant magnetic perturbations (RMPs) in DIII-D [53]. A key observation of those experiments is that RMPs cause density "pumpout" — reduction of the electron density on the separatrix and throughout the pedestal by about a factor of two. The pedestal T_e gradient simultaneously about doubles but T_e on the separatrix and the T_e pedestal height remain approximately constant. These RMP-induced changes can be understood in terms of this new pedestal structure model as follows. First, the relation in (31) implies that if the separatrix density $n_e(a)$ is reduced by a factor of two the density throughout the pedestal will also be reduced by that factor, and the value of $n_e \bar{D}_\eta$ will be similarly reduced. For a pedestal with half the density the T_e gradient must double to carry the same amount of electron heat flow, as indicated in (35). This causes the T_e gradient scale length L_{Te} in (36) to decrease by a factor of two and hence the pedestal electron pressure to also decrease by this factor, as indicated in (39). The reduced $\beta_e^{\rm ped}$ is more likely to be below the peeling-ballooning instability boundary and hence not lead to ELMs. Thus, the preceding discussion of RMP effects on this pedestal structure model is qualitatively consistent with some of the key experimental results [53]. The critical parameter for this interpretation is apparently the reduction of the separatrix density "boundary condition" $n_e(a)$ induced by the RMPs.

X. EXPERIMENTAL VALIDATION TESTS

The preceding sections have demonstrated that the H-mode pedestal structure model developed in this report is quantitatively consistent with many properties of the 98889 DIII-D pedestal [1] and qualitatively consistent with key features of QH-mode, EDA H-modes, I-modes and RMP pedestals. However, its key scaling properties need to be explored experimentally. Also, its predictions need to be tested quantitatively over much wider data sets. And finally, hopefully some of its predictions can lead to new regimes where ELMs are either prevented or better controlled. All of this is needed to truly validate this model.

The paleoclassical transport model provides lower limits on radial transport of plasma density, temperatures and toroidal flow (and hence E_ρ). It is based on Coulomb collision processes that cause the plasma resistivity which produces the magnetic field diffusivity $D_\eta \equiv \eta_\parallel^{\rm nc}/\mu_0$ that is so fundamental to paleoclassical transport and this pedestal structure model. However, like neoclassical transport theory [15, 17], there is no underlying phenomenology that can be tested to validate the fundamental phenomenologies of the model — beyond noting that the electrical resistivity in tokamak plasmas agrees with neoclassical predictions. Also, as noted in section II, paleoclassical transport predictions compare favorably with experimental data from many ohmic-level toroidal plasmas [10] (tokamaks, STs, RFPs, spheromaks), in EC-heated RTP plasmas [12] and in tokamak H-mode pedestals [1, 10].

The pedestal structure model developed in this report depends critically on the pedestal density profile being determined by the paleoclassical diffusive outward density flux nearly balancing its intrinsic inward pinch flux. Further interpretive transport modeling studies like those pioneered by Stacey and Groebner [27] are needed to identify the range of parameters over which density pinches play a critical role in determining density transport in the pedestal. Also, a direct experimental measurement of the density pinch flow would be of fundamental importance in validating this pedestal structure model.

A number of validation tests can be identified for the paleoclassical-based pedestal structure model developed in this report. A hierarchy of the most important tests are:

Fundamental Test #1: When edge fueling effects are negligible [e.g., as in (30) for the 98889 pedestal], is $n_e(\rho)\,\bar{D}_\eta(\rho)$ approximately constant within the pedestal, as predicted in (31)? Exploring this is complicated by the needs for an accurate $Z_{\rm eff}(\rho)$ profile in the pedestal and for a more accurate determination of the neoclassical parallel resistivity profile than ONETWO presently provides, as discussed after Eq. (10) above.

Fundamental Test #2: Is the T_e gradient in the pedestal approximately constant and of the magnitude predicted in (35)? The constancy of $dT_e/d\rho$ depends on the constancy of $n_e\bar{D}_\eta$ and the effective conductive electron heat flow $\hat{P}_e$ through the pedestal. When all other parameters are held constant, larger $\hat{P}_e/n_e$ ratios in the pedestal, particularly at the separatrix, should produce larger T_e gradients. Strongly shaped non-circular cross-section plasmas with larger V' should produce slightly smaller T_e gradients.

Fundamental Test #3: Is the scaling of the electron temperature gradient scale length at the pedestal density mid-point $[L_{Te}/a]_{\rho_n}$ as predicted in (36)? As implied by the preceding set of tests, when other parameters are held constant, the T_e gradient scale length should increase with non-cirularity ($\propto V'$), electron density n_e and temperature T_e at the mid-point of the pedestal density profile (ρ_n). In addition, it should decrease with increased conductive electron heat flow $\hat{P}_e$ at constant $n_e(\rho_n)$.

Fundamental Test #4: Can it be shown that long wavelength ($k_\perp \varrho_i \lesssim 1$) fluctuations within the pedestal do not contribute significantly to plasma transport there? Of course ITG-induced and perhaps trapped electron mode fluctuations should occur and cause significant transport at the top of the pedestal and moving into the core plasma.

Additional, secondary tests that are either less fundamental or result from adding information from other plasma transport models or additional sources and sinks are:

Secondary Test #1: Does the top of the density pedestal occur where $d\ln\bar{D}_\eta/d\rho \lesssim 1/a$ with a height predicted by the minimum of $n_e(a)\bar{D}_\eta(a)/\bar{D}_\eta(\rho_\mathrm{t})$ as indicated by (31) or $\max\{n^\mathrm{ped}\} \sim \dot{N}a/\bar{D}_\eta V'$ in (33)?

Secondary Test #2: Are the edge fueling effects on the pedestal n_e profile as predicted in (32)? And does this fueling effect cause the pedestal n_e profile to be shifted increasingly outward relative to the T_e profile as ϱ_* is decreased (via higher B_t and increasing NB power, fueling) in DIII-D but have no significant effect on the "aligned" profiles in JET, as observed in the JET/DIII-D comparison experiments [30, 31]?

Secondary Test #3: Is the "initial" (i.e., in the $\sim$ 10 ms after an L-H transition or ELM) quasi-stationary pedestal electron pressure height predicted by the β_e^ped in (39)? And at the top of the T_e pedestal do ETG-type fluctuations cause $\chi_e^\mathrm{ETG} \gtrsim \chi_e^\mathrm{pc}$ there?

Secondary Test #4: When charge-exchange effects on the toroidal momentum are negligible, is the total plasma toroidal rotation frequency $\Omega_\mathrm{t} \simeq V_\mathrm{t}/R$ nearly constant within the pedestal at its separatrix value $\Omega_\mathrm{t}(a)$ as predicted by (45)? When charge-exchange momentum losses are significant do they affect $\Omega_\mathrm{t}(\rho)$ as indicated in (46)?

Assuming the basics of this pedestal structure model are experimentally validated and the interpretations of QH-modes, EDA H-modes, I-modes and RMP effects in the preceding section borne out, some possible means for controlling the pedestal structure and thereby avoiding ELMs can be identified. Roughly speaking, the peeling-ballooning instability boundary for ELMs [49] is a combination of the pedestal pressure gradient $d\beta^\mathrm{ped}/d\rho$ being too large and/or the pedestal width $\propto L_P \equiv |d\ln P/d\rho|_{\rho_n}^{-1}$ being too large thereby causing too large a pedestal pressure for the given $d\beta^\mathrm{ped}/d\rho$. Thus, ELMs might be able to be avoided by reducing the pressure gradient in the pedestal and/or the pedestal height, while all other parameters are held constant. Some possible scenarios for doing this are:

Improvement Scenario #1: Reduce the pedestal height by reducing the electron separatrix density $n_e(a)$ for a given $\hat{P}_e$ (via increased neutral pumping or changing the magnetic structure outside the separatrix?) to decrease the $n_e(\rho)$ predicted from (31) and L_{Te} from (36), and thereby reduce the β_e^ped in (39), as apparently occurs with RMPs [53].

Improvement Scenario #2: Reduce the pedestal pressure gradient by reducing $\hat{P}_e/V'$ with larger V' (via more highly shaped plasmas) and/or by reducing the fraction of pedestal electron heat flow (via higher density to increase equilibration of T_e to T_i?) to reduce the T_e gradient as indicated in (35).

Improvement Scenario #3: Add a small controlled density flux in the pedestal (via controlled fluctuations or RF waves resonant there?) without much added electron heat flux to reduce the n_e gradient in the pedestal without affecting the pedestal T_e profile much, as apparently occurs in QH-modes [51], EDA H-modes [52] and I-modes [52].

Improvement Scenario #4: Prevent the pressure increase and inward growth of the pedestal "top" [26] on the plasma energy confinement time scale τ_E that ultimately precipitates an ELM, by decreasing n_e at the pedestal top by reducing $n_e(a)$ on the τ_E time scale?

Achieving control of the density buildup in H-mode pedestals, perhaps via Scenarios #1, #3 or #4, is a generally desirable goal. It may also be critical for ITER to be able to heat a low density H-mode startup plasma toward fusion burning conditions before adding density to increase the fusion power output [54].

XI. SUMMARY

The pedestal structure model developed here assumes paleoclassical plasma transport processes dominate in a quasi-equilibrium H-mode pedestal. The T_e gradient responds to the large electron heat flow through the pedestal by increasing until (35) is satisfied, which yields the T_e gradient scale length given by (36). The T_i gradient and its scale length are determined similarly and specified in (40) and (41). Since local fueling is usually small compared to the characteristic level of paleoclassical density transport in the pedestal, the n_e profile adjusts to satisfy (31). This $n_e(\rho)$ causes the net paleoclassical density transport to be small by nearly balancing its outward diffusive density flux component with its intrinsic inward pinch flux [cf., Eq. (18)]. It also causes "alignment" of the n_e and T_e profiles in the pedestal. When edge fueling effects are significant they can displace the n_e profile outward relative to the T_e profile, as indicated in (32). The initial, transport-limited height of the pedestal is determined by the transition from paleoclassical electron heat transport in the pedestal to ETG-induced electron-gyroBohm-scaling anomalous transport in the core plasma, as given by (39). When toroidal momentum sources and sinks in the pedestal are small compared to paleoclassical toroidal momentum transport processes, the total plasma toroidal rotation is predicted in (45) to be nearly constant in the pedestal, at its value on the separatrix. When charge-exchange momentum losses are significant, they cause $d\Omega_{\rm t}/d\rho > 0$ in the pedestal, as indicated in (47). The radial electric field in the pedestal is determined by substituting the predicted $\Omega_{\rm t}$ into the radial force balance, as indicated in (49).

These pedestal structure predictions have all been shown to agree within about a factor of two with data from the low collisionality 98889 DIII-D pedestal [1]. The predicted pedestal structure does not depend of ϱ_*; this agrees with results from recent JET/DIII-D comparison experiments [30, 31]. The model also provides a framework for interpreting pedestal structure properties in recent DIII-D QH-mode [51], C-Mod EDA H-modes and I-modes [52], and DIII-D RMP [53] experiments — as being due to slightly increased density transport with electron heat transport at the same level having no significant effect.

Finally, the model predicts long period Type I ELMs like those in Ref. [1, 26] occur when electron power flow through the pedestal is small enough so the transport quasi-equilibrium T_e gradient predicted by this model can be reached without triggering ideal MHD peeling-ballooning instabilities [49]. At higher power levels these ELMs are likely to become more frequent. If the transport quasi-equilibrium discussed here cannot be established before a Type I ELM occurs, T_e would increase nearly linearly between ELMs. Or perhaps Type II ELMs would be triggered by high-n ballooning modes in the bottom half of the pedestal before a Type I ELM occurs.

The preceding section discusses a hierarchy of experimental tests that are suggested to validate this pedestal structure model. It also discusses some possible avenues that emerge from this new model for reducing the pedestal electron pressure gradient or height and perhaps thereby avoiding ELMs.

The most important new paradigms that result from this pedestal structure model are: 1) the electron temperature gradient in the pedestal builds up to a large enough magnitude so the large electron heat flux emerging from the core can be carried via paleoclassical radial electron heat transport out to the separatrix; and 2) the pedestal density profile mostly adjusts itself to minimize the net paleoclassical density transport through the pedestal.

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