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PRELIMINARY SEMIEMPIRICAL TRANSPORT MODELS

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ABSTRACT

A class of semiempirical transport models is proposed for testing against confinement data from tokamaks and for use in operations planning and machine design. A reference model is proposed to be compatible with published confinement data. Theoretical considerations are used to express the anomalous transport coefficients in terms of appropriate dimensionless parameters.

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I. INTRODUCTION

A semiempirical transport model is determined by a set of transport coefficients chosen under two constraints. One constraint is that parameters in the model are adjusted to fit confinement results from an experimental data base. The other constraint is that the model include at least one testable physical hypothesis.*

In contrast, a purely empirical transport model is constrained only to fit existing experimental data, without any attempt to relate the functional form of the transport coefficients to the form of the underlying physics of the collisional Vlasov equations. Unless such a model accidently embodies a correct underlying physical hypothesis, it cannot be expected to extrapolate correctly to any new parameter regime. For example, the widely used INTOR/ALCATOR scaling (electron thermal diffusivity $\chi_e \propto 1/n_e$) failed to give the correct scaling of ohmic energy confinement with the size of the most recent generation of tokamaks [Pfeiffer and Waltz, 1979].

We define a theoretical transport model as an attempt to derive transport equations from the fundamental plasma physics equations without explicit constraints derived from the experimental data base. Unfortunately, none of the published theoretical models have been shown to fit even the limited data base considered in this paper.**

*A well-known example is an electron energy transport model [Coppi and Mazzucato, 1979] which fits confinement data from a variety of ohmically heated tokamak discharges. Unfortunately, the rationale for the physical hypothesis in this model was not clearly stated, and the model fails to fit observations of the scaling of energy confinement with machine size and auxiliary heating power.

**The only quantitative data used in this paper is energy confinement data from the first four papers presented at the 1982 IAEA Conference on Plasma Physics and Controlled Fusion Research [Johnson et al., 1982; Nagami et al., 1982; Wagner et al, 1982; Murakami et al., 1982], and the ohmic energy confinement data summarized by Pfeiffer and Waltz, 1979.

Here we take the view that purely empirical models are inadequate for extrapolation and that a complete and accurate theoretical transport model may not be forthcoming soon. Therefore, we propose a new class of semiempirical transport models and give reference parameter values which are most likely to fit experimental data. The transport coefficients are radially local quantities which are independent of the ratio of the Debye length to the system size. The transport coefficients are assumed to depend on the local radial scale height parameters which occur in many theories of plasma turbulence.

Since our reference transport model is based on a limited data base, it must be extended to account for transient phenomena such as rapid current rise, detailed evolution of sawteeth, and the transition from low to high recycling in a diverted scrape-off plasma. After describing the reference model and its motivation, we list some modifications which have been proposed to deal with these transient phenomena. We also make suggestions for how the parameters may be chosen more accurately, and how the list of empirically important parameter dependences may be made more complete to accommodate an expanding data base.

II. REFERENCE MODEL

The following semiempirical transport equations are proposed to describe the transport of thermal particles and energy in tokamaks [Pfeiffer et al., 1980].

$$\frac{1}{V'}\frac{\partial}{\partial t}\left(V'n_a\right) + \frac{1}{V'}\frac{\partial}{\partial \rho}\left(V'\Gamma_a\right) = S_a - Mn_a/\tau_{\parallel} ,$$

$$\frac{1}{V'^{5/3}}\frac{\partial}{\partial t}\left(\frac{3}{2}V'^{5/3}p_j\right) + \frac{1}{V'}\frac{\partial}{\partial \rho}\left[V'\left(q_j + \frac{5}{2}p_j\Gamma_j/n_j\right)\right] = W_j + \frac{e_j}{e}Q_\Lambda + \delta_{je}Q_{OH}$$

$$-\frac{3}{2}\, p_j/\tau_j + Q_{Ej} \;,$$

where V is the volume inside a surface of constant Φ, $\Phi = \int B_\phi \, dS_\phi$ is the toroidal flux passing through the poloidal plane surface S_ϕ, S_ϕ extends out to the effective minor radius ρ, ρ is given by $\pi\rho^2 B_{To} = \Phi$, B_{To} is a charactistic toroidal field that we take to be the vacuum toroidal field at the initial magnetic axis, ϕ is the toroidal angle, $V' = \partial V/\partial\rho$, and

$$n_a = \langle \tau^{-1} \int_o^{2\pi} d\phi \int_{t-\tau/2}^{t+\tau/2} n_a(\rho,\ \theta,\ \phi, \tau')\ d\tau' \rangle$$

is the density of species $\underline{a}$ (fluctuation-averaged over a time interval $\tau \ll |n_a/(\partial n_a/\partial t)|$ and flux-surface averaged using

$$\langle A \rangle = \int_{\theta_{min}}^{\theta_{max}} AJ d\theta \int_{\theta_{min}}^{\theta_{max}} J d\theta \;,$$

where the Jacobian $J = (\nabla\rho \times \nabla\theta \cdot \nabla\phi)^{-1}$, with θ an arbitrary poloidal angle coordinate normalized so that $\theta_{min} = 0$ and $\theta_{max} = 2\pi$ for closed equilibrium flux surfaces). The radial flux Γ_a (specified below) is the similarly averaged value of

$$n_a(\vec{u}_a - \vec{u}_\rho) \cdot \nabla\rho \;,$$

$$\vec{u}_\rho = J\left[(-\partial\rho/\partial t)\big|_{\vec{x}}\ \nabla\theta \times \nabla\phi - (\partial\theta/\partial t)\big|_{\vec{x}}\ \nabla\phi \times \nabla\rho\right]$$

is the velocity of surfaces of constant toroidal flux determined by solving the Grad-Shafranov equation (or by evolving the time derivative of the Grad-

Shafranov equation to determine evolution of torodial flux with respect to a specified set of magnetic coordinate surfaces [Hirschman and Jardin, 1979; Jardin, 1981]) to relate Φ and θ in terms of the Cartesian coordinates $\vec{x}$, $\vec{u}_a$ is fluid velocity of particle species a, S_a is the source of particle species a averaged as described above, and M and $\tau_\parallel$ are the effective parallel Mach number and loss time, specified below. In the heat balances, p_j is the averaged thermal ion or electron pressure q_j is the averaged ion or electron radial conduction flux, Γ_i is the ion flux and Γ_e is the electron flux, W_j is the averaged source due to interaction with radiation or particles other than thermal plasma, Q_Δ is the averaged coulomb energy interchange heating of ions by electrons [Braginskii, 1965]. Q_{OH} is the ohmic heating [Pfeiffer et al., 1980], and Q_{Ej} formally describe the remaining terms which arise when averaging the appropriate moments of the Fokker-Planck equations, as discussed below.

For plasmas with more than one thermal ion species, we sum the ion heat balances and assume a common ion temperature.

For the semiempirical transport models, we take

$$\Gamma_a = -\frac{n_a \chi_{OH}}{\rho}\left(\delta_1 + \delta_2 A_a^{\delta_3} A_i^{\delta_4} Z^{\delta_5} \frac{\partial \ln n_a}{\partial \ln \rho}\right) + \delta_o D_{Bohm} f_q ,$$

$$q_j = -n_j \chi_j \frac{\partial T_j}{\partial \rho} ,$$

where

$$\chi_e = (\chi_{OH} + \chi_I) + (\chi_q + \chi_{kink}) ,$$

$$\chi_{OH} = \chi_{OH*} \left(\frac{n_e}{n_{e*}}\right)^{\alpha_1} \left(\frac{\rho}{\rho_*}\right)^{\alpha_2} \left(\frac{\varepsilon}{\varepsilon_*}\right)^{\alpha_3} Z_{eff}^{\alpha_4} A^{\alpha_5} \left|\frac{\partial \ln n_e}{\partial \ln \rho}\right|^{\alpha_6} \left|\frac{\partial \ln T_e}{\partial \ln \rho}\right|^{\alpha_7} \left|\frac{\partial \ln q}{\partial \ln \rho}\right|^{\alpha_8},$$

$$\chi_I = \chi_{I*}\left(\frac{p}{p_*}\right)^{\beta_1}\left(\frac{n}{n_{e*}}\right)^{\beta_2}\left(\frac{\rho}{\rho_*}\right)^{\beta_3}\left(\frac{B}{B_*}\right)^{\beta_4}\left(\frac{q}{q_*}\right)^{\beta_5}\kappa^{\beta_6}A_i^{\beta_7}\left(\frac{\varepsilon}{\varepsilon_*}\right)^{\beta_8}\left|\frac{\partial \ln p}{\partial \ln \rho}\right|^{\beta_9}\left|\frac{\partial \ln q}{\partial \ln \rho}\right|^{\beta_{10}} ,$$

$$\chi_q = \gamma_o D_{Bohm} f_q ,$$

and

$$\chi_i = \gamma_1 D_{Bohm} f_q + \gamma_2 \chi_i^{CH} .$$

Here α_j, β_j, γ_j, and δ_j are empirically determined constants, χ_i^{CH} is the ion thermal diffusivity [Chang and Hinton, 1982], χ_{OH*} and χ_{I*} are reference values determined empirically near the reference parameters denoted by *, $\varepsilon = \rho/R$ where R is the time-dependent major radius of the magnetic axis, A_i is the mean atomic mass, and $\kappa = b(\rho)/a(\rho)$ where b is the maximum half-height and a is the maximum half-width of the equilibrium flux surface. The formula,

$$q = \frac{5\ \rho_m^2 B_T}{R_m I_{MA}}\left(1 + 1.5\ \frac{\rho^2}{R^2}\right)\left(\frac{1+\kappa^2}{2}\right)\kappa^{1/2} ,$$

(where $\rho_m = 10^2\rho$, $B_T = 10^4$ B, $R_m = 10^2 R$, $I_{MA} = 10^5 I/c$ with $I(\rho)$ the current inside the flux surface) gives a reasonable fit to the flux-surface average of the safety factor computed from the equilibrium for limiter discharges [Stanbaugh, 1983], but avoids the singularity at a divertor separatrix. $p = p_i + p_e$ is the total thermal pressure. The Bohm diffusion coefficient is $D_{Bohm} = cT_e/(16eB)$. f_q is a function with the property that $f_q \approx 1$ for $q < q_1$ and $f_q \approx 0$ for $q > q_1'$ where typically $q_1 \approx q_1' \approx 1$. (An agreed reference function to be used for f_q in various transport codes would be desirable.) χ_{kink} is zero when the ideal pressure-driven internal m=n=1 kink mode is stable, and it is a given function (yet to be determined) when this mode is unstable.

For a convenient reference normalized to $n_{e*} = 4 \times 10^{13} cm^{-3}$, $\rho_* = 27$ cm, $B_* = 12000$ G, $q_* = 3.6$, $\varepsilon_* = 0.3$, and $p_* = 5 \times 10^4$ erg/cm^3 (which corresponds to a reference temperature of 0.4 keV), we take

$$\chi_{OH*} = \chi_{I*} = 10^4 cm^2/s \ ,$$

$$\delta_o = 1,\ \delta_1 = \delta_2 = 0.5,\ \delta_3 = \delta_4 = 1,\ \delta_5 = 0,$$

$$\alpha_1 = -0.9,\ \alpha_2 = -1.0,\ \alpha_3 = 1.9,\ \alpha_4 = -0.1,\ \alpha_5 = 0,\ \alpha_6 = 1.0,\ \alpha_7 = \alpha_8 = 0,$$

$$\beta_1 = 1.0,\ \beta_2 = -0.2,\ \beta_3 = 0.4,$$

$$\beta_4 = -2.0,\ \beta_5 = 2.2,\ \beta_6 = -2.2,\ \beta_7 = -0.3,\ \beta_8 = 1.4,\ \beta_9 = 1.5,\ \beta_{10} = -1.5,$$

$\gamma_o = \gamma_1$, and $\gamma_2 = 2$, for reasons discussed below. The reference transport scalings are then

$$\chi_{OH} \propto n^{-0.9} \rho^{-1.0} \varepsilon^{1.9} Z_{eff}^{-0.1} |\lambda_{n_e}| ,$$

$$\chi_I \propto p^{1.0}\ n_e^{-0.2}\ \rho^{0.4}\ B^{-2.0}\ q^{2.2}\ \kappa^{-2.2}\ A_i^{-0.3} \varepsilon^{1.4} |\lambda_p|^{3/2}\ |\lambda_q|^{-3/2},$$

$$\chi_q = D_{Bohm} f_q \ ,$$

$$\Gamma_a = -(n_a \chi_{OH}/\rho)(0.5 + 0.5\ A_i A_a \lambda_{n_a}) + D_{Bohm} f_q ,$$

where

$$\lambda_x \equiv \partial \ln x / \partial \ln \rho.$$

For the other terms in the reference model, we take $M_\parallel = 0.5$ (appropriate for a limiter or a divertor with low recycling), $\tau_\parallel = L_\parallel / v_s$ with $L_\parallel = \pi q R$ (appropriate for a toroidally symmetric limiter or single-null divertor), $v_s = [(T_e + T_i)/m_i]^{1/2}$ with m_i the mean ion mass, $\tau_i^{-1} = 2/\tau_\parallel$, $\tau_e^{-1} = 2\gamma_e/\tau_\parallel$ with $\gamma_e = 2.9$ (appropriate in the absence of secondary electron emission) [Ogden et al., 1981]. For the coupling term Q_{Ej}, we take $Q_j = (e_j/e)(\Gamma_j/n_j)\partial p_j/\partial \rho$ (which ignores anomalous electron-ion coupling and neoclassical viscous damping of poloidal rotation [Pfeiffer et al., 1980; Hirshman and Jardin, 1979]). In the reference model for poloidal flux diffusion, the resistivity is taken to be classical [Braginskii, 1965], and the bootstrap current is ignored [Hogan, 1981]. Note that, for simplicity, these reference transport coefficients contain no neoclassical effects (except that the ion thermal conductivity has neoclassical scaling).*

III. METHODS

Here we describe the basis for our prescription of semiempirical models

*Sources in the reference model may be computed in an equilibrium using the first three equilibrium moments [Lao et al., 1981] heating and fueling by fast ions from the steady-state solutions to the Fokker-Planck equations in a periodic cylinder [Gaffey, 1976], neutral sources from integral transforms in an equivalent circle model [Tamor, 1981], radiation from coronal equilibrium of an average ion model [Post et al., 1977], and radio-frequency heating using traced rays interacting with a separately computed nonthermal particle distribution [Karney et al., 1983; Valeo, 1982; Kluge et al., 1982]. When using the reference transport with other methods of equilibrium and source computation, it would be desirable to make an estimate of the increase in accuracy or error which arises from using the nonstandard computation method.

of flux-surface-averaged radial transport of thermal ions and energy. (Only energy transport in tokamaks heated by neutral beam injection is discussed in detail. Energy confinement in ohmically heated tokamaks has been adequately discussed [Pfeiffer and Waltz, 1979], and there is insufficient data published on other heating methods to be useful here.) We begin with scalings of global energy confinement observed in the neighborhood of a specified reference set of plasma parameters. This data is converted to a scaling of global energy confinement with respect to dimensionless variables. By hypothesizing that energy transport is independent of the ratio of the Debye length to plasma minor radius [Kadomstev, 1975; Connor and Taylor, 1977], the scaling of energy confinement with machine size is obtained. Global confinement scalings are then converted to local transport relations, using experimental profile information and further physical hypotheses to prescribe the dependence of the radial fluxes on local radial scale heights. Transport equations for particles and the subdivision between ion and electron energy transport are then discussed.

A. χ_I

Data from the first four papers presented at the 1982 IAEA Baltimore conference are sufficient to deduce a global energy content scaling,

$$\beta \propto P^{0.5\pm0.2} I^{1.1\pm0.4} B^{-2.1\pm0.2} n^{0.1\pm0.1} \kappa^{a_5} a^{a_6} \text{ near } \vec{x}_1.$$

The notation "near $\vec{x}_1$" signifies that this scaling is only valid in the neighborhood of a convenient reference point

$$\vec{x}_1 = (2\text{MW}, 0.2\text{MA}, 1.2\ \text{T}, 4 \times 10^{13}\text{cm}^{-3}, 1, 0.4\ \text{m}),$$

where the data are particularly abundant. In other words, the exponents quoted above are merely the coefficients in a locally valid logarithmic power series expansion, with $\partial \ln\beta/\partial \ln P = 0.5\pm0.2$, $\partial \ln\beta/\partial \ln I = 1.1\pm0.4$, etc., when evaluated at $\vec{x}_1$. Here β is the total plasma-pressure/toroidal-magnetic-pressure ratio measured with a diamagnetic loop, P is the total absorbed heating power, I is the toroidal plasma current, B is the toroidal magnetic field, n is the midplane line-averaged electron density, $\kappa = b/a$ is the maximum half-height, b, divided by the midplane half-width, a.

The above scaling exponents were determined by drawing tangents by hand through the published parameter scans. Where the total absorbed power was not reported, an ohmic heating correction $P_{OH} \propto (n^{3/2}Ra/\beta_p^{3/2} I)$ was added to the quoted beam heating. This correction was only 0.07 MW at the reference point, $\vec{x}_1$. Where necessary, we took $\beta_p \propto (\epsilon/q)^2\beta$ with the proportionality constant normalized to results from ISX-B. Power scaling curves are assumed to intersect the point (β=0, P=0). The $\pm$ errors quoted above for $\partial \ln\beta/\partial \ln P$ and $\partial \ln\beta/\partial \ln I$ are standard deviations after averaging the exponents derived from PDX [Johnson et al., 1982]. Since these standard deviations are at least comparable to the scatter in the parameter scan data reported from each experiment, we believe that negligible additional error is introduced by our simple method of fitting the data. The quoted value of $\partial \ln\beta/\partial \ln B$ at $\vec{x}_1$ is an inverse-error weighted average of $\partial \ln\beta/\partial \ln B = -2.3 \pm 0.3$ (derived from text in the ISX-B paper) and $\partial \ln\beta/\partial \ln B = -2.0 \pm 0.1$ (derived from plotted PDX data). This quoted error in $\partial \ln\beta/\partial \ln B$ is only an approximate value, but it is also compatible with the data presented in the D-III paper. The density scaling exponent is taken from the D-III paper and is consistent with data from the other papers.

Scaling of β with elongation was reported as <u>linear</u> in D-III at fixed

P,B,n,a and q. Using this information (after eliminating the toroidal current in favor of q using the appropriate approximate scaling $q \propto B\kappa^{3/2}/I$) gives

$$\beta \propto P^{0.5} q^{-1.1} B^{-1.0} n^{0.1} \kappa\, a^{b_6} .$$

The experimentally determined limits on scaling of confinement with machine size are best examined by eliminating β in favor of the energy confinement time τ_E. Using the power balance, $\beta B^2 a^3 \kappa/\tau_E \propto P$ at fixed a/R, we obtain

$$\tau_E \propto P^{-0.5} q^{-1.1} B^{1.0} n^{0.1} \kappa^2 a^{c_6} ,$$

or, in a more familiar form (with q eliminated in favor of I),

$$\tau_E \propto P^{-0.5} I^{1.1} B^{-0.1} n^{0.1} \kappa^{0.35} a^{c_6 - 1.1}$$

(again at fixed a/R). A least squares fit to confinement times (interpolated to the reference point $\vec{x}_1$, using the power and current scalings derived from each data set) gives $c_6 - 1.1 \approx 1.5 \pm 1.2$. Here the quoted error is the difference between the best fit and the scaling of confinement with size between ISX-B and PDX (which gives the most pessimistic results). Thus, the scaling coefficient $c_6 \approx 2.6 \pm 1.2$ is not determined from the experimental data base with useful accuracy.

A theoretical hypothesis is evidently required to obtain size scaling near our reference point. Using the power balance to eliminate the heating power, P, in favor of the pressure, $p \propto nT \propto \beta B^2$, we obtain the form of confinement scaling analyzed by Connor and Taylor

$$\tau_E \propto (nT)^{-1.0}\, q^{-2.2}\, B^{2.0}\, n^{0.2}\, \kappa^3\, a^{(2c_6-3)}.$$

As first pointed out by [Kadomtsev, 1975] and later by [Connor and Taylor, 1977], casting a confinement scaling of this type into appropriate dimensionless variables only yields a result independent of $\lambda = \lambda_{Debye}/a$ if the exponents of the above dimensional variables satisfy a constraint. For a confinement scaling of the form $\tau_E \propto p^{d_1}\, q^{d_2}\, B^{d_3}\, n^{d_4}\, \kappa^{d_5}\, a^{d_6}$, this constraint is

$$d_6 = \left(5 + 10\, d_1 + 5d_3 + 8d_4\right)/4.$$

With the reference values derived above ($d_1 = -1.0$, $d_3 = 2.0$, $d_4 = 0.2$), this gives a size scaling exponent of $d_6 = 1.65$. Comparing to experimental measurements of size scaling coefficient $d_6 = 2c_6 - 3$, we get $d_6 = 2.6 \pm 2.4$ from the experiment and $d_6 = 1.65 \pm 1$ from the Kadomtsev-Connor-Taylor (KCT) constraint. Since the experimental limits on d_6 are very approximate, and since the value of d_6 derived from the KCT constraint contains an experimental uncertainty of order $\pm$ 1 due to uncertainties in the values of d_1, d_3, and d_4, there is no significant inconsistency between these two scalings. Experiments on TFTR and JET should soon provide an important test of the validity of the KCT constraint.

As a simple method of obtaining a self-consistent KCT-constrained global energy confinement scaling, we adopt the reference scaling

$$\tau_E \propto p^{-1.0}\, q^{-2.2}\, B^{2.0}\, n^{0.2}\, \kappa^3\, a^{1.65},$$

while noting that more extensive size scaling experiments may force a revision of these exponents. Making a final change to the equivalent circle radius, $\rho \propto a\kappa^{1/2}$, to be used in the transport equations given above (and dropping the extraneous significant digit), we obtain

$$\tau_E \propto P^{-1.0}\, q^{-2.2}\, B^{2.0}\, n^{0.2}\, \kappa^{2.2}\, \rho^{1.6}.$$

To define a model for the average plasma thermal flux Q across a magnetic flux surface, we note the definition $Q = \langle 3p/2\rangle_V\ (V/A)/\tau_E^{tr}(\rho)$ where $\langle\ \rangle_V$ is volume average inside the flux surface labelled by ρ, $V/A \approx \rho/2$ is the volume/surface ratio, and $\tau_E^{th}(\rho)$ is the confinement time for energy loss due to thermal plasma transport. Since $\tau_E^{th}(\rho)$ is comparable to the global confinement time τ_E over much of the plasma in the discharges considered here, the estimate $Q \sim p\rho/\tau_E$ should correctly give the observed global confinement scaling if we replace the global variables in the above τ_E scaling by their local values averaged over a flux surface. However, such a model might not reproduce the observed radial plasma profiles. This is because there are a number of physically significant dimensionless variables which are often roughly constant from one discharge to another but may vary in the radial direction in a given discharge. These include the inverse aspect ratio $\varepsilon = \rho/R$, the electron temperature profile factor $\lambda_{Te} = \partial \ln T_e/\partial \ln \rho$, and the additional dimensionless variables listed in the first column of Table 1. Motivated by theories which suggest that temperature gradients may drive an anomalous electron thermal transport which dominates energy losses, we have approximated $Q \approx q_e = -\chi_e n_e T_e \rho \lambda_{Te} = -n_e \chi_e \partial T_e/\partial \rho$ and written χ_e (near the reference point $\vec{x}_1$) in the form of χ_I given above. Since pressure gradients may enhance and shear may reduce the transport, and the inverse aspect ratio should be significant [Rewoldt et al., 1978; Carreras et al., 1983], we have included scalings with $\lambda_q = \partial \ln q/\partial \ln \rho$ and with ε. Scaling with mean ion mass A_i has been added to match experimental observations [Wagner et al., 1982].

Following a common convention [Pfeiffer et al., 1980; Hirshman and Jardin, 1979; Hawryluk, 1980], we have included a "convective" term $(5/2)T_e\Gamma_e$ in the transport equations, but this term will typically be small. (In the absence of a detailed and substantive derivation of the complicated turbulent averages which lead to the electron heat balance, the inclusion of this convective term should be viewed as a convenient artifice and not be given particular physical significance. As an extreme example of why this may be true, consider the form of the convective term for a hypothetical turbulent convection which moves one warm electron inward for every two cool electrons moved outward. Then even if $\Gamma_e = -D\partial n_e/\partial \rho$, it is possible to have n electron energy outflux when $\Gamma_e > 0$ and $\partial T_e/\partial \rho = 0$.)

The reference values of these additional transport scaling parameters have yet to be chosen with care. The pressure gradient and magnetic shear scaling exponents are chosen to be $\beta_8 = 1.5$ and $\beta_9 = -1.5$, respectively, in analogy with a theory of pressure-gradient-driven turbulence [Carreras et al., 1983]. If turbulence due to $\vec{E} \times \vec{B}$ drifts is important, the absolute magnitudes of these parameters may be lower [Rewoldt et al., 1978], but their signs should still be plus and minus, respectively. The scaling of χ_I with ε was determined by fitting radial logarithmic derivatives to "typical" profiles of the form $\chi_e \propto \rho^2$, $T_e \propto (1 - x^2)^2$, $p \propto (1 - x^2)^3$, and $q = 1 + 2.6\, x^2$, where $x = \rho/(a\kappa^{1/2})$ at a reference point $\rho = (2/3)a\kappa^{1/2}$, where transport analysis codes are typically most accurate. This analysis should be repeated with actual experimental profiles [Wieland, 1983; Kaye, 1983] from the existing unpublished data bases.

B. χ_{OH}

Pfeiffer and Waltz examined global energy confinement scalings for 118 ohmic discharges from eleven different tokamaks with mean plasma parameters $\vec{x}_2$ = (0.3 MW, 0.2 MA, 3 T, 4 × $10^{13}cm^{-3}$, κ = 1, 0.2 m) and with variation of inverse aspect ratio near ε = 0.2 [Pfeiffer and Waltz, 1979]. Applying the KCT constraint used above, Pfeiffer and Waltz obtained a good fit to the global confinement scaling of the form

$$\tau_E \sim n^{0.9}B^0\varepsilon^{-1.9}a^{3.0}Z_{eff}^{0.1}\left(T_e n^{0.1}B^{-0.1}\varepsilon^{0.8}a^{0.4}I^{-0.8}Z_{eff}^{-0.5}\right)^\gamma \text{ near } \vec{x}_2 \ .$$

Here the "temperature scaling exponent," γ, is unknown because the term in parentheses does not vary significantly from one ohmic heating experiment to another. There is no value of γ which makes this low-β scaling identical to the moderate-β scaling described above within the statistical uncertainty of the data base (± 0.1 to ± 0.2 in each exponent). (This leads to the unsurprising conclusion that a single power law of this type is inadequate to fit the composite data.) Converting this global ohmic energy confinement to a local thermal diffusivity scaling as described above, we obtain, with $\gamma = \alpha_7 = \alpha_8 = 0$ for simplicity,

$$\chi_{OH} \propto r^{-1.0}n^{-0.9}\varepsilon^{1.9}Z_{eff}^{-0.1}A_i^0\ \lambda^{\alpha_6} \ .$$

For a reference transport model, we assume that instabilities will be triggered, at least, by excessive density gradients [Rewoldt et al., 1978]. For simplicity, we set α_6 = 1 and ignore possible dependences on magnetic shear and other "hidden variables." A density profile factor α_6 = 1 gives a $\chi_{OH}(\rho)$ profile roughly consistent with measured values of χ_e even with no "temperature scaling exponent," so we set $\gamma=0$ in the reference transport model.

This reference model is, of course, somewhat arbitrary. As starting point for a more systematic search for a semiempirical model of this type, we would suggest varying γ to fit data from ohmic heating discharges cooled by significant radiative losses, fueled by pellet injection, and heated with low power levels of auxiliary heating. It should also be noted that recent studies of ohmic heating in D-III show improvement in confinement with plasma current and only weak improvement with $\bar{n}_e$ at high density [Nagami et al., 1982]. While this behavior is qualitatively consistent with summing the above thermal diffusivities, a quantitative fit is unlikely without a systematic reevaluation of the parameters in our reference model.

C. χ_q and χ_{kink}

Here we describe reference models for the time-average effect of sawteeth and pressure-driven internal kink modes. For an enhancement of transport to reproduce the observed profile flattening where $q < 1$, we add to each diffusivity a unit multiple of a large diffusion coefficient, arbitrarily taken to be D_{Bohm}, times a profile factor f_q. In the GOSPEL code [Larrabee et al., 1981], $f_q = c_q$ for $q < 1$ and $c_q \exp[(1 - q)/\epsilon_q]$ for $q > 1$, where typically $c_q = \epsilon_q = 0.1$; in the BALDUR code $f_q = 1.1(1 - q)^2$ [Silverman et al., 1983].

On theoretical grounds, it seems likely that pressure-driven internal kink modes will prevent β from indefinitely increasing with heating power, P, as fast as it does in the neighborhood of our 2 MW reference point. Although there is insufficient data for a detailed scaling analysis [Kaye, 1983], results from the highest power neutral injection experiments in PDX [Johnson et al., 1982] are compatible with such a saturation of $\beta(P)$. Because of incompletely characterized fast ion losses, however, thermal transport in such discharges is not well characterized [White et al., 1983].

For sensible extrapolation to high-β tokamak experiments and reactors with a semiempirical transport model, it is essential to include information from theoretical studies of β limits. Internal kink modes leading to experimentally observed fishbone oscillations appear to be most important on empirical and theoretical grounds. A theoretical estimate of transport due to this process is required which has the following properties: the dependence of the onset of this instability on plasma shape, currents, and pressure [Izzo et al., 1983] should be included. For studies of reactors and the PDX bean-shaped plasma, stabilization at very high β should be included in this description. A simple estimate of the effect of the internal kink mode of energy and particle transport should be included. This estimate should be compatible with the observation that global energy confinement in PDX with 6 MW heating is within a factor of two of the scaling derived above in the neighborhood of the reference point $\vec{x}_1$ where P = 2 MW. It should be noted that a complete computational model which includes χ_{kink} will also require a description of the effect of fishbone oscillations on fast ion distributions [White et al., 1983].

D. Particle transport

Since there have been no systematic studies of particle transport comparable to those described above for energy transport, the choice of models for the flux-surface averaged (and fluctuation averaged), Γ_j, must be based on qualitative considerations. Among these are the following [Coppi and Sharky, 1981; Strachan et al., 1982; Becker and Singer 1981]. (i) Hydrogen transport appears to have a component which is diffusive, in the sense that flux is down the density gradient (inward with hollow density profiles due to pellet injection and outward after termination of gas puffing). This is suggested by experiments with pellet injection. (No strong dependence on temperature gradient has been reported with discharges having hollow electron temperature profiles). (ii) There may be some difficulty in modeling both quasistationary and time-varying density profiles with a single transport relation of the form $\Gamma = -D\partial n/\partial\rho = -(nD/\rho)\lambda_n$, which has a simple linear dependence on the density profile parameter λ_n. In particular, density profiles in ohmic heating discharges cannot be matched by taking $D \propto \chi_e$ (as might be suggested by global confinement in discharges with low opacity, as well as by the most naive estimates of the effect of saturated microturbulence). This conclusion is not altered by changing the Ware pinch within the uncertainty of the underlying theory (about a factor of two). As the simplest model which may satisfy these qualitative considerations, we add fluxes of the form $\Gamma = -(nD/\rho)(\lambda_n + \delta)$ where $D \sim \chi_{OH}$ and $\delta \sim 1$. To obtain reasonable hydrogen confinement times and avoid the pinch coefficient being too large, reference values of $D = \chi_{OH}/2$ and $\delta = 1/2$ are taken. In the absence of detailed scaling studies, no mass or charge scaling is used in the reference models. These numbers should be recalibrated against experimental data. Note that only the low-β turbulence scaling is used in Eq. (27), since comparable enhancement of particle transport at moderate β is not necessarily predicted theoretically. Scalings for impurity diffusivity are determined from impurity loss rates [Marmar et al., 1982].

Since there may be convective plasma motion in the sawtooth cycle, a component $D = D_q$ should be added for sawtooth discharges. As present data analysis has not defined this contribution well (due to the weak particle source with the gas puffing generally used), for simplicity, we take $D_q = \chi_q$

in the reference model.

E. χ_i

Until now, we have assumed that energy losses are dominated by electron heat transport. For a more comprehensive model, especially at very high collisionality where ion losses are significant, we add a scaling for ion energy transport. We have chosen an admittedly oversimplified form for division of energy transport between ion and electron channels because we believe anomalous electron and ion energy losses are in any case likely to be correlated with anomalous division of the outflowing energy between ions and electrons. In particular, we include no anomalous electron/ion energy interchange, and we ignore a typically small coupling due to viscous damping of poloidal rotation [Hirshman and Jardin, 1981]. Instead, we choose an anomalous ion thermal conductivity equal to twice the value recently derived [Chang and Hinton, 1982], since a value of this order typically gives a reasonable fit to the observed ion temperature profiles [Kaye, 1983]. Since the combination of anomalous ion heat conduction and anomalous energy interchange cannot be empirically distinguished from a contribution to anomalous electron energy loss, there is little to be gained by using a more complicated ion heat balance in the absence of a very detailed anomalous transport theory. Finally, we note that, as for electron convection, the anomalous ion energy convection in our reference transport model should be regarded merely as a convenient form for comparison with results from conventional transport codes.

IV. EMBELLISHMENTS

A. High recycling divertors

Our standard transport model matches observations made in the absence of a high recyling divertor by assuming plasma is lost at ends of materially bounded flux surfaces at a substantial fraction, $M_\parallel$, of the sound of speed, v_s. Advection-dominated energy losses are computed from classical sheath theory (without secondary electron emission, for the reference model parameter values given above).

For plasmas with significant recycling occuring in a divertor or pumped limiter chamber, the above scrape-off model is inadequate. For such cases, we append a recycling channel to the ends of each of the above-described flux surfaces where they leave the main plasma chamber. The purpose of the recycling channel model is to obtain an order-of-magnitude estimate of how an intense recycling region alters the particle and energy outflux from the ends of the main-plasma scrape-off region. To this end, we neglect viscosity, radial and diamagnetic flows, and the variation of flux surface area along magnetic field lines [Singer and Langer, 1983; Morgan and Harbour, 1981]. In making these approximations, we neglect effects which are typically only comparable to those retained in the present model of flows parallel to the magnetic field in recycling channel surfaces with constant cross section. This should not compromise our goal of illustrating the main effects of a transition from low to high channel recycling. With these approximations, the relevant conservation equations for a plasma with a single hydrogen species are

$$\frac{\partial \Gamma}{\partial z} = S,$$

$$\frac{\partial}{\partial z} \left(mnu^2 + 2nT\right) = F ,$$

$$\frac{\partial}{\partial z} \left[\left(\frac{1}{2} mnu^2 + \frac{5}{2} nT\right) u + q\right] = W,$$

where $\Gamma = nu$ is the parallel ion flux, m is the ion mass, $T = (T_e + T_i)/2$ is the average temperature, $F = \int \Sigma_i mv$ is the non-Coulomb friction (often dominated by charge exchange with neutral atoms), $q = \kappa T^{5/2}(\delta T/\delta Z)$ is the parallel heat conduction [Braginskii, 1965], and $W = \sum_{j=i,e} \int \Sigma_j (1/2)mv^2$ is the energy source due to non-Coulomb collisions (Σ_j are the non-Coulomb collision operators, and integration is over velocity space). We have neglected small contributions from electron inertia and ion conduction and set $T_e = T \equiv (T_e + T_i)/2$ to obtain a one-fluid model for momentum and energy conservation.

We also neglect charge exchange in computing the source terms S,F, and W and set $F = 0$ and $W = ES$ where $E = 40$ eV to account for radiation losses. These approximations produce order-of-magnitude error only in the parallel particle flow at low temperature ($\lesssim 10$ eV). (At such low temperatures, a collision dominated 2-d model of the recycling channel is appropriate and necessary for more accurate modeling [Singer and Langer, 1983]). Neglecting non-Coulomb friction, the steady-state pressure balance integrates to [Morgan and Harbour, 1981]

$$mn_1 u_1^2 + 2nT_1 = mn_2 u_2^2 + 2n_2 T_2 ,$$

where subscript 1 denotes conditions at the entrance to the recycling channel and subscript 2 denotes conditions at the material boundary. With any significant recycling, the flow entering the recycling channel is very subsonic in this model. We therefore ignore the ion inertia except at the material boundary. This gives a parallel momentum balance of the form

$$2n_1 T_1 = mn_2 u_2^2 + 2n_2 T_2 = 4n_2 T_2 ,$$

for sonic flow at the material boundary (i.e., for $u_2^2 = 2T_2/m$). For the other two conservation equations, we are also interested in steady-state solutions, since core plasma quantities change slowly compared to scrape-off relaxation times in typical transport code simulations. It is convenient to integrate the particle and energy balances to obtain

$$\Gamma_1 = \Gamma_2 - \int S d\ell \ ,$$

$$Q_2 - \left(\frac{5}{2}\ \Gamma_1 T_1 + q_1\right) = \int W d\ell \sim E \int S d\ell .$$

Here, $Q_2 = 2(\gamma_i + \gamma_e)\Gamma_2 T_2$ is the energy flux assumed to flow through the plasma sheath.

To close this set of equations, we chose a simple neutral absorption model. We take

$$\int S d\ell = \Gamma_2 \left(1 - f_{pump}\right) \left[1 - \exp\left(-L_2/\lambda_2\right)\right].$$

It is assumed here that a fraction $(1 - f_{pump})$ of the ions striking the material boundary are returned to the recycling channel. A fraction $\exp(-L_2/\lambda_2)$ of these neutrals penetrates through the recycling channel to the main plasma chamber, and the remaining fraction $[1 - \exp(-L_2/\lambda_2)]$ is ionized in the recycling channel. We take $\lambda_2 = n_2 \langle \sigma v_e \rangle / v_o$ where $\langle \sigma v_e \rangle = \exp[\sum_{n=0}^{6} A_n \ell n^n (k^{-1} T_2)]$, A_n are coefficients given by Freeman and Jones, k is Boltzman's constant, and $v_o = (2E_o/m_o)^{1/2}$ is the energy of the recycled neutrals. While this model is necessarily simple, it should reproduce the qualitative features of a recycling channel. The effective channel length, L_2, and pumping fraction, f_{pump}, can be adjusted to attempt to match the properties of this model to the behavior of a complicated geometry.

Finally, we approximate the heat flux q_1 by the flux limited conduction formula

$$q_1 = \text{Min} \begin{cases} \alpha n_1 v_{the}(T_1 - T_2) \\ \kappa T_1^{5/2}(T_1 - T_2)/L_1 \end{cases} ,$$

where $0.03 < \alpha < 0.3$, and $v_{the} = (2T_1/m_e)^{1/2}$ for the usual flux-limited theories [Bell, 1981; Clause and Balescu, 1982]. To account for the possible effect of poloidal potential gradients on parallel electron heat conduction [Ohkawa et al., 1983] it suffices in the present model to include an option where $\alpha \lesssim 1$ and $v_{the} = (2T_2/m_e)^{1/2}$. (The use of T_2 instead of T_1 in this formula accounts for the inability of electrons in the main plasma chamber to surmount the potential barrier.) For the reference model, we take $\alpha = 0.1$ and $v_{the} = (2T_1/m_e)^{1/2}$. The value of κ is

$$\kappa = 3.16 \frac{3m_e^{1/2}}{4(2\pi)^{1/2} e^4 \ell n\Lambda} ,$$

where $\ell n\Lambda \approx 11$ for $n \sim 10^{14}\text{cm}^{-3}$, $k^{-1}T \sim 10$ eV (with $k = 1.6 \times 10^{-12}$ erg/eV) [Braginskii].

B. Sawtooth evolution

Here we express the "insulating region" model of time-dependent sawtooth evolution [Hwang 1983] in terms of dimensionless parameters. In this model, transport is inhibited in a small region near the outermost q=1 surface unless a sawtooth disruption is in progress. To accomplish this inhibition, we set

$$\left.\begin{array}{l} D_a[(\rho_1 - \Delta\rho_1) < \rho < \rho_1] \rightarrow C_{402}\, D_a \\ \chi_e[(\rho_1 - \Delta\rho_1) < \rho < \rho_1] \rightarrow C_{403}\, \chi_e \end{array}\right\} \text{for } -\lambda_p(\rho_1) < C_{404} ,$$

where the ρ_1 label is the position of the outermost $q = 1$ surface, Δ is a small number (generally taken to have a value which makes $\Delta\rho_1$ the width of a computational zone), χ_e and D_a are the electron thermal and hydrogen diagonal diffusivities calculated without application of the transport inhibition, and C_{402} and C_{403} are numerical coefficients with reference values of C_{402} and $C_{403} = 0.01$.

The sawtooth disruption is modeled by increasing diffusivities inside the $q = 1$ surface when the $\lambda_p = \partial \ln p/\partial \ln \rho$ exceeds a specified value:

$$\left.\begin{array}{l} D_a\ (\rho < \rho_1) \rightarrow C_{406}D_a \\ \chi_e\ (\rho < \rho_1) \rightarrow C_{408}\chi_e \end{array}\right\} \lambda_p(\rho_1) > C_{404}.$$

Reference values of the additional coefficients given here are $C_{406} = 10$, $C_{408} = 20$, and $C_{404} = 0.01$. Using the pressure gradient to trigger enhanced transport might seem more appropriate to modeling transport due to pressure-driven internal kink instabilities than to current-driven sawteeth. [In fact, the evolution of transport effects of pressure-driven fishbone oscillations can be simulated in the BALDUR transport code by setting $D_j\ (\rho < \rho_1) \rightarrow C_{406}D_a$ and $\chi_e(\rho < \rho_1) \rightarrow C_{408}\chi_e$ when $\lambda_p(\rho_1) > C_{405}$.] The use of this model for current-driven sawteeth is justified by the observation that electron temperature oscillations dominate the pressure variations in typical simulations. Since electrical conductivity and current density are proportional to $T_e^{3/2}$ in typical sawtooth simulations, the predominance of temperature oscillations in the pressure variation makes the model sufficiently accurate for its usual use, which is to generate a propagating heat pulse perturbation for use in the study of other transport phenomena.

C. Current startup

Transport of poloidal magnetic flux is known to be anomalous with rapid increase in toroidal current in tokamaks [Granetz et al., 1979]. To approximate correct current profiles soon after current startup, it has been suggested that the neoclassical parallel resistivity should be increased by a large factor [Larrabee, 1983], so we set

$$\eta = f_J \, \eta_{neoclassical} ,$$

where η_{neo} is the neoclassical parallel conductivity and

$$f_J = \begin{cases} 1, & \partial \ln J/\partial \ln \rho < 0 \\ C_J & \partial \ln J/\partial \ln \rho > 0 , \end{cases}$$

and

$$q - \varepsilon_J < m < q + \varepsilon_J ,$$

where m is any integer, and typically $\varepsilon_J = q/20$ and $C_J >> 1$.

D. Other processes

With completion of the internal kink model and correct values of the adjustable parameters, it is possible that the transport model described so far will reproduce almost all tokamak discharge plasma parameters. Here we describe a few likely exceptions to this generality.

One possible exception is a local flattening of electron temperature profiles near the q=2 surface. Although relevant transport models have been described, we do not think the experimental observations are yet clear enough to justify including this complication in the present reference transport model.

Another possible problem is poloidal asymmetry of the electron density. While this may provide valuable clues to the physics of anomalous transport, we again do not believe the observations are clear enough to justify an attempt to reproduce them in the present model.

Direct effects of collisions with circulating fast ions, for example, is another process which may be important in impurity transport. We have omitted such effects because it is not yet clear that direct effects of fast ions are significant compared to unavoidable indirect effects mediated by the thermal plasma.

Finally, there is no effect of toroidal plasma rotation in the anomalous transport model. While it is not yet clear that inertial effects due to toroidal rotation are of major importance in determining particle and energy transport, it would not be surprising if this were the case. Inclusion of such effects would first require solution of the toroidal momentum balance, which reportedly can be successfully modeled by setting the toroidal momentum diffusivity equal to the $\chi_i/\varepsilon^{3/2}$, where χ_i is the neoclassical thermal diffusivity [Howe, 1983]. Using this model appears at present to be the best course when constructing extended anomalous transport models where fluxes may depend on the toroidal Mach number, as noted below.

V. DISCUSSION

We have described a class of semiempirical transport models and suggested a set of reference parameters which, as far as we know, are compatible with results from all tokamak discharges which are free from major and minor disruptions and from unknown fluctuations in source terms. Here we discuss the most usual ways of applying such models and suggest methods of expanding them to accommodate an increasingly diverse data base.

Semiempirical transport models have two functions which cannot be adequately performed by purely empirical or theoretical models. One function is extrapolation of plasma parameters for use in machine design and operations planning. Similarity extrapolations (keeping the same relevant dimensionless parameters with different values of physical parameters) using the KCT constraint may be particularly useful. For other extrapolations, correct theoretical input is necessary to choose appropriate functional forms to supplement the Taylor series expansion around well-studied reference points. (An example is the inclusion of χ_{kink} in our reference model.) This should allow relatively confident planning of startup and early operations for large tokamaks as well as physically reasonable estimates of upgraded performance. No existing purely empirical or theoretical models are likely to be as adequate to accomplish these various tasks.

The other function of semiempirical models is to serve as an improved link between empirical observations and derivation of better theoretical models. By forcing reduction of data in terms of physically relevant and meaningful parameters, construction of semiempirical models leads to presentation of transport scalings in a form which should be more helpful to theorists who look to data analysis to guide their selection of appropriate physical processes to investigate.

This is not to imply that there is no role for transport simulation using purely empirical or theoretical models. Empirical transport models often contain a useful summary of the data base in direct terms of experimentally varied parameters, and the simplicity and easy accessibility of such models may suggest a limited utility in interpolative operations planning. On the other hand, testing of purely theoretical models using transport calculations is a most direct and efficient way of providing theorists with information about the accuracy of the approximations they use. But the large middle ground left for semiempirical models suggests that careful attention should be paid to their construction and testing. To this end, we discuss methods for increasing the accuracy and generality of the type of model we have proposed.

A. Data fitting

The broad scope of this report has left little room for accurate fitting of the reference model parameter values to a large data base. We therefore merely suggest a list of priorities for obtaining more accurate parameters for the reference model. The best approach is probably to fix all parameters in the reference model except for one parameter of greatest interest. This parameter can then be adjusted to give the best fit to data from transport analysis codes. From the above discussion of χ_I, it should be clear that the most fruitful initial choice to refine the reference model would be to better determine the inverse aspect ratio scaling exponent, β_8. Additional parameters could then be successively included and evaluated by linear regression analysis to the extent justified by the available data. A useful order for adjusting these parameters might be β_9 (pressure gradient scaling), β_{10} (shear), β_1, $\beta_2, \ldots, \beta_7$ (plasma parameters); α_6 (density gradient scaling for low-β plasmas), γ, α_1, $\alpha_2 \ldots \alpha_5$, α_7, α_8 (low-β plasma parameters); δ_2 (magnitude of particle diffusivity), δ_1 (particle pinch), δ_3, δ_4, δ_5 (mass and impurity charge); γ_0, δ_0, γ_1 (time-average sawtooth effects); and γ_2 (ion energy transport). Selected parameter scans or better theoretical estimates might be useful in determining some of these values without a full regression analysis. For particularly complete parameter scans, inclusion of some second order terms in the local logarithmic Taylor series expansion might also be justified.

B. Dimensionless variables

The list of variables in our reference model is likely to be insufficient to describe local transport scalings for a truly comprehensive data base. Table 1 gives a complete list of fundamental scaling variables found in recent theories of tokamak transport for a pure hydrogen plasma with a steady axisymmetric boundary. As far as we know, all variables in transport theories for such plasmas can be expressed in terms of the variables listed in Table 1 and ratios of fixed fundamental constants. For example, the electron thermal diffusivity recently derived by Carreras et al.[23] can be written as

$$\chi_e \propto \rho v_{the} \beta \nu \delta^3 q^4 \varepsilon^{-5/2} \lambda_p^{-3/2} \lambda_q^{3/2} ,$$

where $v_{the} = (2T_e/m_e)^{1/2}$, $\nu = \nu_e/\Omega_e$, $\nu_e \propto n_e/T_e^{3/2}$ is the electron-electron collision frequency, $\Omega_e = eB/(m_e c)$, $\delta = (v_{the}/\Omega_e)/\rho$, and $\lambda_x \equiv \partial \ln x/\partial \ln \rho$. The analogous χ_I term in our reference model is

$$\chi_I \propto \rho v_{the} \beta^{0.6} \nu^{0.2} \delta^{0.6} q^{2.2} \kappa^{-2.2} A_i^{0.3} \varepsilon^{1.4} |\lambda_p|^{-3/2} |\lambda_q|^{3/2}.$$

A comparison of these two χ_I scalings suggests that the theory of Carreras et al. may not fit the global confinement data used here.

There are two ways to improve semiempirical transport models by adding scalings dependent on the parameters listed in Table 1. First, an analysis of the generic implications of various theoretical hypotheses may be used to suggest that certain parameters are particularly important. Numerical values in the functional forms suggested by theory may then be refined and described in Section V.A. Second, new tokamak parameter scans of the types listed in Table I may allow testing of new scaling hypotheses. While there is rarely a unique correspondence between adjustable and natural dimensionless parameters in tokamaks, noting some of the correspondences listed in Table 1 might aid in extracting appropriate scalings by regression analysis. For multispecies plasmas, the list of potentially relevant variables is longer than that given in Table 1. For each species, at least, mass, charge, and density, scale height may be important, as indicated in our reference model particle diffusivities. Moreover, fast ion populations may stimulate or inhibit anomalous thermal plasma transport. (Most studies have concentrated on anomalous loss or deceleration of the fast ions themselves, a topic which lies beyond the scope of our reference transport model for thermal plasma.) However, it may be necessary eventually to add parameters which describe the direct effects of fast ions on thermal plasma transport.

For nonaxisymmetric boundary conditions, direct effects of static ripple and radio-frequency fields on thermal plasma transport may also be important. Aside from effects of toroidal field ripple on ion thermal conductivity, which has been described adequately elsewhere, there are no useful treatments of such effects, to our knowledge.

VI. CONCLUSION

We have outlined a comprehensive model of thermal particle and energy transport for tokamak plasmas. This model is an improvement over previous empirical transport models in several ways. First, the model is based on a broader data base for global energy confinement than previously published models. Second, the model is consistent with the physically reasonable Kadomstev-Connor-Taylor constraint, in that it does not depend on the ratio of the Debye length to system size when expressed in suitable dimensionless variables. Third, the transport coefficients depend on the radial gradients of plasma parameters, as is appropriate for anomalous transport processes due to the fluctuations related to such gradients. Fourth, the model is sufficiently comprehensive to reproduce a wide variety of transport phenomena, including impurity transport, sawtooth oscillations, and transition from low to high divertor recycling regimes. When completed with a suitable model of the effects of pressure-driven internal kinks, the model should provide a useful tool for operations planning of existing experiments and for design of experiment upgrades and fusion reactors.

Methods for refining and extending the reference transport model are also proposed. While the optimum method of accomplishing this has yet to be worked out in detail, we hope that the attempt will stimulate an effort to bring methods of empirical data analysis closer to a form compatible with the underlying physics of plasma transport. This should lead to the evolution of more accurate and useful semiempirical transport models.

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APPENDIX: Calibration of Semiempirical Transport Parameters

Here we present a set of parameters which have been adjusted to match results from an ASDEX discharge. The sole purpose of giving these parameters is to serve as a useful starting point for testing the semiempirical transport model against a wider data set. (Neither the author nor members of the ASDEX-team represent that this particular set of parameters provides a general description of tokamak transport.)

The simulation was constrained to match the following experimental parameters: $\bar{n}_e = 4 \times 10^{13} cm^{-3}$, $I_p = 375$ kA, $B_T = 2.2$ T, R = 164 cm, $r_{separatrix} = 40$ cm, tangential injection of D^o into an initially 95% H^+ plasma. Maximum applied beam heating power was 1.6 MW at 40 keV, 0.8 MW at 20 keV, and 0.5 MW at 17 keV. Total beam power was ramped (linearly in the simulation) from 2.1 MW to 2.9 MW in 40 ms. Parameters in the transport model were adjusted to fit measured temperature and density profiles just before and after the beam heating ramp. (For details of the simulation and comparison with the data, see the author.)

The resulting thermal diffusivities and particle fluxes were

$$\chi_i = 2\, \chi_i^{neo} ,$$

$$\chi_e = \chi_q + \chi_{OH} + \chi_I ,$$

$$\Gamma_a = -n_a \left[v_{OH} + 0.67\, v_{ware} + \frac{(D_q + D_{OH})}{n_a} \frac{\partial n_a}{\partial r}\right],$$

where χ_i^{neo} is the usual neoclassical ion thermal diffusivity (provided by P. Rutherford) used in the BALDUR transport code. The time-averaged effect of sawteeth of the transport coefficients is given by $D_q = 0.75[1.1(1-q^2)\ D_{Bohm}]$

for $q < 1$, $\chi_q = 1.5\ D_q$ (with $D_q = \chi_q = 0$ for $q > 1$). The ware pinch is the usual BALDUR formula (provided by R. Hawryluk). The new semiempirical transport contributions are

$$\chi_I = 1.5 \times 10^3\ (p/p_*)^{1.0}(n_e/n_*)^{-0.2}(r/r_*)^{0.4}(B/B_*)^{-2.0}(q/q_*)^{2.2}A_i^{0.0}$$

$$(\epsilon/\epsilon_*)^{1.4}|\lambda_p|^{1.0}|\lambda_q|^{-2.0}\ ,$$

$$\chi_{OH} = 6 \times 10^3(n_e/n_*)^{-0.9}(r/r_*)^{-1.0}(\epsilon/\epsilon_*)^{1.9}Z_{eff}^{0.0}\ |\lambda_n|^{0.5}|\lambda_{TE}^{+}|^{-0.5}$$

$$\left[\frac{T_e}{T_*}\ (\frac{n}{n_*})^{0.0}(\frac{B}{B_*})^{-0.9}(\frac{\epsilon}{\epsilon_*})^{0.0}(\frac{r}{r_*})^{0.40}(\frac{q}{q_*})^{0.8}Z_{eff}^{-0.5}\right]^{-0.5},$$

$$D_{OH} = 0.6\ \chi_{OH}\ ,$$

$$v_{OH} = 0.1\ \chi_{OH}/r_+\ ,$$

where $p = n_iT_i + n_eT_e$, $p_* = 2.4 \times 10^{13}$ keV/cm^3, $n_* = 2 \times 10^{13}$cm^{-3}, $B_* = 2.2$ T, $r_* = 27$ cm, $\epsilon_* = 27/164$, $q_* = 2.2$, $|\lambda_p| = |\lambda_n| + 0.5\ (|\lambda_{Te}| + |\lambda_{Ti}|)$, $\lambda_n = (r/n_e)\partial n_e/\partial r$, $\lambda_{Te} = (r/T_e)\partial T_e/\partial r$, and $\lambda_{Ti} = (r/T_i)/\partial T_i/\partial r$. The weak singularity in $\partial q/\partial r$ at the separatrix is approximated by using

$$|\lambda_q| = |(r/q_{cyl})\ \partial q_{cyl}/\partial r| + 0.1/(|q_{cyl} - q_{sep}| + 0.01),$$

where q_{cyl} is the safety factor computed by BALDUR and q_{sep} is the value of q_{cyl} at the separatrix. This formula gives a good fit to the $\partial q/\partial r$ profile (provided by K. Lackner) computed from an MHD equilibrium for ASDEX. (The number 0.01 in the denominator of the MHD correction to $|\lambda_q|$ is the fraction of the plasma volume occupied by a single computational zone in the simulation, and it is added to avoid a numerical singularity.) Temperature gradient stabilization occurs above a threshold given by the definition

$$|\lambda_{Te}^{+}| = 0.1 + |\lambda_{Te}|,$$

in the above formula for χ_{OH}. Finally, in computing the anomalous pinch velocity, v_{OH}, a small addition of the form $r_{+} = r + 5$ cm is added to the denominator to avoid $\chi_{OH} = 0/0$, as $\chi_{OH} \rightarrow 0$ when $r \rightarrow 0$.

While the above transport coefficients were calibrated only to Ohmic heating and so-called L-mode data, they have an appropriate form to describe H-mode confinement as well. Despite the unfavorable $\chi_I \propto p^{1.0}$ scaling in the χ_I term which dominates confinement with beam heating, deterioration of plasma confinement is arrested in this simulation near the onset of the H-mode. This results from propagation of the unfavorable pressure gradients out to the region stabilized by magnetic shear. A parameter search to optimize the fit to ASDEX and PDX H-mode data will be undertaken in Princeton.

As pointed out in the PDX discussion paper "Understanding the H-Mode" circulated earlier this year, the L → H-mode transition should be accompanied by changes in sources (e.g., of a factor of two) which are likely to be comparable to changes in transport on the closed flux surfaces near the plasma periphery. Estimating these changes requires an improved model of the response of the density (and temperature) at the separatrix to changes in the divertor parameters. To this end, a simple "two chamber" model of the parallel scrape-off losses was incorporated into the Garching version of BALDUR. (This model differs from a similar model developed by Langer et al. at Princeton only in that the effective divertor plasma temperature is to be input from experimental data rather than estimated from the poloidal energy transport equation.)

The particle loss rate from the main chamber is given in this model by

$$\tau_{\parallel}^{-1} = Mv_s/L_{\parallel},$$

where $v_s = [(T_e + T_i)/m_i]^{1/2}$, $L_{\parallel} = \pi q_{circ}R$, and M is the "parallel Mach number." As described in the paper "Preliminary Semiempirical Transport Models," M is related to the divertor pressure and the effective opacity encountered by recycling particles until they enter the main chamber from the opening of the divertor throat. In the simulation described here,

$$M = (\alpha/2)(T_e/T_{div})^{1/2} \quad \text{where } T_{div} = 10 \text{ eV},$$

$$\alpha = 1-.9(1-\exp(-\lambda/L_{eff})) \quad \text{where } L_{eff} = 70 \text{ cm},$$

$$\lambda = (n_{div} \langle\sigma v\rangle_{div}/(2kE_o/m_o)^{1/2})^{-1} \quad \text{where } E_o = \text{ave. neutral E},$$

$$n_{div} = n_e T_e/(2T_{div}),$$

where $\langle\sigma v\rangle_{div}$ is the electron impact ionization rate coefficients at the temperature T_{div}, computed from the usual Freeman-Jones formula used in BALDUR. Electron energy loss by heat conduction parallel to the magnetic field gradient was included using the classical Braginskii formula with $\nabla_{\parallel} T_e = (T_e - T_{div})/L_{\parallel}$. This heat flux was limited to a maximum value (in cgs units) of $q_{\parallel} = 0.3\ (3n_eT_ev_{the}/2)$ where $v_{the} = (2T_e/m_e)^{1/2}$.

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TABLE 1

Plasma Parameters	Tokamak Parameters	Dimensionless Parameters
"Connor-Taylor" parameters:		
$T=(T_e + T_i)/2$	Heating power, P_{heat}	$\beta=16\pi nT/B^2$
n(electron density)	Feedback control gas puffing	$\nu=\nu_e/\Omega_e$
B(toroidal field)	Toroidal field coil currents	$\delta=\rho_{Le}/r$
ρ(effective minor radius)	Limiter position	$\lambda=\lambda_{Debye}/r$
Multispecies parameters:		
A_i(mean atomic mass)	Working gas	A_i
Z_{eff}	Impurity injection	Z_{eff}
Geometric parameters:		
ϵ(inverse aspect ratio)	Inside vs outside limiter	ϵ
κ(elongation)	Elongation shaping coils	κ
τ(triangularity)	D/bean shaping coils	τ
q(safety factor)	Plasma current, I	q
Profile factors:		
$d\ell np/d\ell n\rho$	Heating profile(beam energy, etc.)	λ_p
$d\ell n/d\ell n\rho$	Fueling profile(pellet injection)	λ_n
$d\ell nq/d\ell n\rho$	Heating profile(beams, RF, α's)	λ_q
$d\ell nJ/d\ell n\rho$	Current ramp rate	λ_J
$d\ell nT_e/d\ell n\rho$	Heat pulse propagation	λ_{Te}
$d\ell nT_i/d\ell n\rho$	RF Heating method(ICRD,LH,ECRH)	λ_{Ti}
T_i/T_e	RF Heating method(ICRH,LH,ECRH)	T_i/T_e
Electric fields:		
E_{loop}(toroidal electric field)	Current drive	$E=E_{loop}/E_{Dreicer}$
$v_{toroidal}$(toroidal rotation velocity	Beam injection angle	$M_{tor}=v_{toroidal}/v_{thi}$
E_θ(poloidal potential gradient)	Varies slightly with neutral beam isotope and with many of the above parameters.	$\phi=e\tilde{\phi}/T_e$

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