PPPL--2708

DE90 014821

Issues in Tokamak/Stellarator Transport and Confinement Enhancement Mechanisms

F. W. Perkins
Plasma Physics Laboratory
Princeton University
Princeton, NJ 08543

Abstract

At present, the mechanism for anomalous energy transport in low- β toroidal plasmas - tokamaks and stellarators - remains unclear, although transport by turbulent $E \times B$ velocities associated with nonlinear, fine-scale microinstabilities is a leading candidate. This article discusses basic theoretical concepts of various transport and confinement enhancement mechanisms as well as experimental ramifications which would enable one to distinguish among them and hence identify a dominant transport mechanism. While many of the predictions of fine-scale turbulence are born out by experiment, notable contradictions exist. Projections of ignition margin rest both on the scaling properties of the confinement mechanism and on the criteria for entering enhanced confinement regimes. At present, the greatest uncertainties lie with the basis for scaling confinement enhancement criteria. A series of questions, to be answered by new experimental/theoretical work, is posed to resolve these outstanding contradictions (or refute the fine-scale turbulence model) and to establish confinement enhancement criteria.

DISCLAIMER

This report was prepared as an account of work sponsored by an agency of the United States Government. Neither the United States Government nor any agency thereof, nor any of their employees, makes any warranty, express or implied, or assumes any legal liability or responsibility for the accuracy, completeness, or usefulness of any information, apparatus, product, or process disclosed, or represents that its use would not infringe privately owned rights. Reference herein to any specific commercial product, process, or service by trade name, trademark, manufacturer, or otherwise does not necessarily constitute or imply its endorsement, recommendation, or favoring by the United States Government or any agency thereof. The views and opinions of authors expressed herein do not necessarily state or reflect those of the United States Government or any agency thereof.

COPIES OF THIS DOCUMENT ARE UNLIMITED

MASTER

1. Introduction.

Plans are now being developed for a compact magnetic confinement machine¹ which will be principally heated by thermonuclear reactions and may realize the long-sought goal of controlled thermonuclear ignition. Success would unequivocally demonstrate the scientific feasibility of magnetic fusion power. The cost of such a device, while modest in comparison to the superconducting supercollider² (~\$6B) or space station³ (~\$30B), is nonetheless high enough (~\$ 0.8-1.5B) that a careful tradeoff between device parameters and cost could result in appreciable savings. What additional cost is certainty of ignition worth, given the many important goals a power amplification experiment with an amplification $Q \approx 5$ could fulfill? Energy confinement dictates the size, and hence the cost, of the planned experiments. Uncertainties in extrapolating energy confinement physics translates into uncertainties as to whether a particular device will attain $Q \approx 5$ or ignition. The questions this article addresses are: First, what is our knowledge of the energy confinement mechanisms and their scaling in tokamaks and stellarators that would form the basis of a tradeoff? Second, will enhanced confinement regimes, found routinely in many tokamaks, be realized in an ignition device? There are significant unresolved issues and contradictions. The role of this article is to make the issues and the context in which they arise specific. Our "conclusions" take the form of a Table of 21 questions whose answers require new experimental and theoretical research.

Let us remark that, while energy confinement times in stellarators are appreciably smaller than tokamaks, they are in accord with theoretical tokamak confinement scaling formulas, suggesting a common mechanism for anomalous energy transport. Present stellarator experiments⁴⁻⁷ are comparable in size to the tokamaks of the early 1970's (minor radius ≈ 20 cm). Thus, one can hope that new, larger stellarator experiments such as Wendelstein VII-X⁸ and the Large Helical Device,⁹ which will have minor radii ≈ 60 cm almost as large as TFTR, will find correspondingly longer confinement times and enhanced confinement regimes. One also looks to stellarators as devices where nondimensional parameters (e.g., the q-profile, magnetic well geometry, and trapped particle fraction) differ from tokamaks and hence a source of insightful fluctuation experiments and important data on confinement mechanisms. In addition to anomalous transport, stellarators suffer direct orbit loss by dint of their 3-dimensional geometry¹⁰; these losses are outside the scope of this article.

The reliability of confinement projections for ignition experiments clearly depends on the degree of extrapolation from present results. The most meaningful extrapolation is the ignition margin - fusion power divided by power losses - which is proportional to the product $n_D T(0) T_i(0) \tau_E$ because the DT fusion reaction rate increases roughly as T_i^2 with temperature. Let density and temperature be given by parabolic profiles

$$T_i = T_i(0) (1 - r^2/a^2)^{\alpha_T} \quad n = n(0) (1 - r^2/a^2)^{\alpha_n} \quad (1)$$

Calculations by G. Bateman (private communication) then indicate that ignition requires

$$n_D T(0) T_i(0) \tau_E = (36 \pm 3) \frac{1 + 2\alpha_n + 2\alpha_T}{1 + \alpha_n + \alpha_T} \cdot 10^{20} \text{ m}^{-3} \text{ keV s} \approx 60 \cdot 10^{20} \text{ m}^{-3} \text{ keV s} \quad (2)$$

for reasonable values of α_n and α_T (e.g. $\alpha_n \approx 0.5$, $\alpha_T \approx 2.0$). Equal thermonuclear and auxiliary heating occurs when $Q=5$, which requires $n_D T(0) T_i(0) \tau_E = 30 \text{ m}^{-3} \text{ keV s}$. The best experimental results, summarized in Table 1, lie roughly an order of magnitude below these goals.

The proposed Compact Ignition Tokamak (CIT) facility¹ differs from the three experiments of Table 1 principally in its higher magnetic field and plasma current. Electron and ion temperatures are expected to be close in CIT because the rate of ion-electron temperature relaxation exceeds the inverse energy confinement time and α -particles preferentially heat electrons, rather than ions. Hence plasmas with $T_e \approx T_i$ are the appropriate starting point for extrapolation. CIT's goal is to study the physics of self-heated plasmas, defined by $Q \geq 5$.

Reactor engineering issues as well as fusion physics are the focus of the ITER device,¹¹ now under study by an international conceptual design team. These issues require an ignited plasma. Table 2 compares the various devices.

Recently, DIII-D has obtained discharges with $n_D T(0) T_i(0) \tau_E \approx 1.0 \cdot 10^{20} \text{ m}^{-3} \text{ keV s}$ with $T_i(0) \approx 4.5 \text{ keV}$, $T_e(0) \approx 3.5 \text{ keV}$ (R. Stambaugh, private communication, 1990). Extrapolating from DIII-D to CIT principally involves predicting the consequences of a factor-of-5 increase in toroidal field and plasma current; the shape and size are quite similar. This must yield a factor-of-60 increase in $n_D T(0) T_i(0) \tau_E$ for ignition. Fortunately,

theoretical microturbulence scaling [see Eqs.(23),(35)-(39)] predicts this factor will be attained; empirical scaling [Eq.(40)] falls short of ignition but is still sufficient to achieve $Q \approx 5$.

The plasma volume of the TFTR tokamak is almost identical to CIT, so that the required factor-of-30 increase in $n_D(0)T_i(0)\tau_E$ (under $T_i \approx T_e$ conditions) must be accomplished by doubling the toroidal field and elongating plasma shape in a way that increases plasma current. Because of its focus on the supershot regime with $T_i \gg T_e$, the highest values of $n_D(0)T_i(0)\tau_E$ for TFTR with $T_i \approx T_e$ occur under ohmic heating conditions¹² and have rather low temperatures $T_i(0) \approx 1.4$ keV. The circular TFTR device does not benefit from the improved β -limit and confinement time associated with elongation. Thus scaling TFTR data to CIT is subject to greater uncertainties than scaling DIII-D and JET results and is not attempted.

JET has successfully combined pellet injection with neutral beam injection (NBI) heating¹³ to reach $n_D(0)T_i(0)\tau_E \approx 5.8 \cdot 10^{20} \text{ m}^{-3} \text{ keVs}$ at $T_i(0) = 8.3$ keV in an elongated plasma, only a factor-of-10 below ignition requirements. One must note that high $nT\tau$ discharges in JET are characterized by a rapidly increasing energy content which is terminated by an abrupt release of carbon, called "carbon blooms", from hot spots on the limiters. Hence, the maximum value of $nT\tau$ for JET is expected to increase when the planned single-null divertor, which should eliminate hot spots and carbon blooms, is installed. Extrapolation of JET data trades off the benefits of higher plasma toroidal field and plasma elongation versus penalties of reduced size. Again, microturbulence scaling supports realization of ignition in CIT.

The point to be made here is that the extrapolation to ignition is appreciable. Factors of 10 to 60 in $nT\tau$ are involved. The most optimistic extrapolations predict $nT\tau \approx 100 \cdot 10^{20} \text{ m}^{-3} \text{ keVs}$ for CIT, a factor-of ≈ 2 above requirements. It is against this background that we must measure our knowledge of confinement.

Table 1. Values of $n_D(0)T_i(0)\tau_E$ Achieved in Experiment *

Experiment	$n_D(0) T_i(0) \tau_E \quad (10^{20} \text{ m}^{-3} \text{ keV s})$	
	$T_i \gg T_e$	$T_i \approx T_e$
JET ¹	8.9	5.8
DIII-D ²	0.5	1.0
TFTR ³	4.2	2.0

¹ A. Gibson and the JET team, "A Review of the JET Project", International Plasma Physics Conference, New Delhi, November 1989; G. Schmidt, private communication, January 1990.

² R. Stambaugh, private communication, January 1990.

³ M.G. Bell *et al* in Plasma Phys. Control. Nucl. Fusion Res. 1988 (Proc. 12th Int. Conf., Nice) p27; R. J. Hawryluk *et al* in Plasma Phys. Control. Nucl. Fusion Res. 1986 (Proc. 11th Int. Conf., Kyoto) p51.

* See also Table 7.

Table 2. CIT, ITER, and Present Tokamaks
($T_i = T_e$ conditions with maximum $n_D(0) T_i(0) \tau_E$)

Experiment	R_m	a_m	κ	B_T	I_{MA}	$q_{0.95}$
CIT	2.1	0.66	2.0	10	11	3.2
DIII-D	1.7	0.62	2.0	2.1	2.8	3.7
JET*	3.0	1.0	1.9	2.8	4.2	3.3
TFTR	2.4	0.7	1.0	5.0	1.6	3.2
ITER [†]	5.5	1.8	2.0	5.3	18	3.1

* Single-null H-mode. [†] Technology phase.

2. Identification of Tokamak/Stellarator Transport and Confinement Enhancement Mechanisms

How does energy flow across magnetic surfaces in a tokamak plasma? The easy answer is too fast; classical transport would be much slower than the observed "anomalous" heat fluxes. It is generally believed that fine-scale instabilities, driven by density and temperature gradients, arise in a tokamak plasma and generate a heat flux as a consequence of their non-linear self-consistent saturation. But there are a variety of instabilities and saturation mechanisms. First principles, quantitative predictions require theoretical calculations of turbulence in a 5-dimensional phase space -- even more challenging than ordinary 3-dimensional hydrodynamic turbulence. For example, to explain the difference between L-mode and H-mode confinement in terms of electrostatic microinstabilities (a leading candidate), one requires computations of sufficient accuracy to show that bulk plasma thermal diffusivity changes by a factor ≈ 5 in response to changes in plasma profiles. Furthermore, contradictions exist between fundamental theoretical scaling properties, derived from dimensional analysis, and empirical confinement scaling relations, derived from an extensive data base of tokamak experiments. Below, we argue that one could interpret the empirical scaling relations as contradicting the fundamental tenet of theory that plasma gradients are the ultimate cause of anomalous transport. The questions are then: Which mechanisms govern energy transport and confinement enhancement in current research tokamaks and what are their scaling properties? Can we identify these mechanisms from regression analyses of tokamak data bases? Which mechanisms will govern transport in CIT? And, can we develop useful understanding and control of these mechanisms in CIT?

The resources that can be brought to bear on answering these questions consist of an impressive data base of tokamak and stellarator global confinement time measurements and two decades of theoretical research on the linear instabilities which plague axisymmetric tokamaks as well as on their nonlinear consequences. Regression techniques applied to the data base have generated empirical scaling relations (which are often unfortunately called "scaling laws") expressing global confinement time in terms of parameters the experimentalist can control, such as plasma current I_p , toroidal field B_T ,

line average density n_e , plasma size and shape, etc. Perhaps the best-known empirical scaling relation is the 1984 Goldston-Aachen formula¹⁴

$$\tau_E (\text{sec}) = 0.037 I_{MA} P_{MW}^{0.5} R_m^{1.75} a_m^{0.37} \kappa^{0.5} (A_i/1.5). \quad (3)$$

Since then, candidate empirical formulas have multiplied profusely.¹⁵ Goldston's formula is representative of what has become known as L-mode confinement and characterizes the confinement time which tokamaks can realize with nominal operations. Additional features, such as defining the plasma boundary magnetically with a separatrix instead of by material limiters, can enhance observed confinement times by roughly a factor-of-two over L-mode. This change coincides with quite different fluctuation levels, density profiles, and temperature profiles in the vicinity of the separatrix and is called H-mode confinement.¹⁶ When intense central NBI heating and fueling is coupled with a first wall vacuum conditioning that absorbs impinging ions, a different enhanced confinement regime, dubbed "supershots",¹⁷ is obtained with $T_i \gg T_e$ and peaked density profiles. Peaked profiles have also improved confinement in ASDEX and JFT-2M.¹⁸

One might hope that theorists could overcome the technical difficulties engendered by a 5-dimensional phase space and offer at least qualitative guidance on how confinement will scale, even though experiment must be called on to identify and resolve quantitative issues, such as the L-mode/H-mode transition. Indeed, qualitative scalings¹⁹⁻²² can readily be found by dimensional analysis for electrostatic microinstabilities (where the bulk of theoretical research lies) and checked against the experimental data base. The predicament tokamak confinement scaling faces is that this experiment-theory comparison has lead to mixed results. A general qualitative and quantitative agreement is marred by distinctly different scaling in several important variables.

In this author's view, the program to identify a dominant transport mechanism has reached an impasse. Many experimentalists will not accept electrostatic microinstability scaling because of the contradictions which will be fully described below. Theoreticians, with good theoretical justification, continue to focus the bulk of their work on electrostatic microinstabilities which, we argue in Eq. (20), can never yield the Goldston Aachen formula. The highest priority in tokamak transport research must be with those

efforts which will identify the dominant transport mechanism and break this impasse. It is not clear at present whether the identification of thermal confinement scaling with regression analyses is subject to unappreciated systematic biases or whether theorists have been devoting the bulk of their effort to a nondominant transport mechanism.

2.1. Identification of Transport Mechanisms.

A first step in achieving a limited predictive capability of tokamak transport is to identify the dominant transport mechanism. Theorists offer two main possibilities: First, transport can result from turbulent, electrostatic $E \times B$ particle drifts across magnetic surfaces. Second, microscopic magnetic fluctuations can destroy the magnetic surfaces themselves, permitting high energy circulating electrons to stream along magnetic field lines and transfer heat from hot to colder regions. Experiments must be devised to distinguish between these phenomena. There is also the possibility that large-scale magnetic fluctuations, such as locked modes or error fields associated with low-order rational surfaces, influence tokamak transport.²³ Experimentally, discharges exist in which such phenomena degrade confinement even below L-mode. The degree to which large-scale magnetic transport persists in optimized discharges is not clear experimentally. Theoretically, this area is underdeveloped, but receiving renewed and deserved interest under the name of neoclassical MHD.^{24,25} Sophisticated computational methods²⁶ are required to evaluate the breakup of equilibrium magnetic surfaces into large scale islands in both tokamak and stellarator configurations.

Within the first class of electrostatic $E \times B$ transport, two important subclasses exist, local, saturated turbulence and marginal stability states, the latter particularly in the context of instabilities driven by the ion temperature gradient.²⁷ Pedagogically, it is useful to discuss the extreme marginal stability model while recognizing that, in real tokamaks, this model oversimplifies the physics, especially near the plasma periphery. The marginal stability model is based on the fact that the Ion Temperature Gradient (ITG) mode (often called the η_i -mode) has a marginal stability condition of the following generic form

$$\frac{1}{T_i} \frac{dT_i}{dr} = \left(\eta_{ic} \frac{1}{n_i} \frac{dn_i}{dr} - \frac{1}{R} f(q, s, T_e/T_i, \dots) \right) \quad (4)$$

where the second term has been related to the shear length or toroidicity effects depending

upon which author one consults.²⁸ Here η_{ic} determines the critical value of η_i

$$\eta_i = \left(\frac{1}{T_i} \frac{dT_i}{dr} \right) \left(\frac{1}{n} \frac{dn}{dr} \right)^{-1} \quad (5)$$

and itself can be a function of nondimensional parameters such as T_i/T_e and $L_s/L_n \equiv (Rq/sn)(dn/dr)$. The statement of marginal stability is that, according to calculations of renormalized hydrodynamic theories,²⁷ the heat flux associated with exceeding the marginally stable gradient (4) by a modest factor (e.g., factor-of-2) will greatly exceed the available flux corresponding to the auxiliary heating power. It follows that the ion temperature gradient must be everywhere close to its marginal value (4) and one can integrate (4) to find

$$T_i(0) = T_{i,edge} \exp \left\{ - \int_0^a dr \left[\eta_{ic} \frac{1}{n} \frac{dn}{dr} - \frac{f}{R} \right] \right\} \quad (6)$$

Succintly, the central ion temperature is directly related to the edge ion temperature. Edge confinement physics, which relates the edge temperature to the heat flux, governs $T_{i,edge}$ and hence $T_i(0)$. In the extreme marginal stability model, confinement in the bulk plasma is ineffective; it is edge confinement that determines tokamak performance. In reality, the reference thermal diffusivity is expected to fall with temperature [see Eq.(7)], and the transition from marginal stability to saturated turbulence will be gradual, but still localized to the outside regions of a tokamak. Because edge confinement determines central temperature, marginal stability is said to be a nonlocal transport mechanism. Experiments show edge confinement to be sensitive to edge density profiles which, in turn, respond to particle fueling (see Sec. 2.7).

It is important to devise experimental tests that can separate marginal stability states, which have a nonlocal relation between heat flux and gradients, from local saturated turbulence theories where a complicated (and still unknown) local nonlinear relation exists between heat flux and gradients. For example, marginal stability implies that $T_i(r)$ responds rapidly to changes in $n_i(r)$, which can be changed by pellet injection, ELM's,^{16,29} and gas puffs. Marginal stability also predicts that turbulence levels will respond rapidly if the plasma temporarily departs from marginal stability, suggesting that

microwave/laser scattering should observe bursts of turbulence coinciding with pellet injection and ELMs. A closely correlated question is: Do diffusivities and turbulence levels in the plasma interior respond instantaneously to the H-mode transition at the plasma edge, or do they change as a result of density profile changes propagating inward? In ASDEX,¹⁶ changes in profile could account for the improved global confinement. Very accurate, time-dependent observations of plasma energy balance profiles and fluctuations are required here.

One should also note that the argument leading to marginal stability relies on a theoretical derivation²⁷ which concludes that the heat flux would appreciably exceed the available heat flux unless the plasma is close to marginal stability. If this derivation should prove to overestimate heat flux, the case for marginal stability will be less compelling.

2.1.1 Fine-Scale Electrostatic Turbulence. Local, fine-scale, saturated electrostatic turbulence is a leading candidate for the mechanism governing confinement in tokamaks. *Fine-scale* is defined here to mean that the poloidal and radial scale size of the turbulence (and hence step-size in a diffusion process) is determined solely by the poloidal or toroidal ion gyroradius $\rho_{\theta,T} = (2T_i/M)^{1/2}/\Omega_{\theta,T}$ and other nondimensional parameters (β , q , s , L_n/L_s , η_i , T_e/T_i , R/a , etc.) some of which may be moderately large numerically. In particular, fine-scale means that the parameter ρ_i/a does not enter in determining scale size. Equivalently, the step size in microinstability diffusion is independent of device size. The widely-used "mixing length rule" saturation estimates, which state that fluctuating gradients become comparable to background gradients (or $E \times B$ velocities comparable to equilibrium diamagnetic drift velocities), fulfill the fine-scale ansatz provided the long-wavelength cutoff of the turbulent spectrum depends only on gyroradius. Fine-scale turbulence includes the well-known microinstabilities³⁰ — ion temperature gradient modes, dissipative and collisionless trapped electron drift modes. At the long wavelength end ($k\rho \approx 0.06$), trapped ion modes are often not small compared to machine size in high temperature tokamaks and the possibility of turbulent step size scaling with device size remains open. Fine-scale turbulence qualitatively differs from hydrodynamic turbulence where scale sizes proportional to system size (e.g., pressure gradient scale height in atmospheric and stellar convection) dominate transport. Experiments to measure the scale size of tokamak turbulence are of paramount importance. Confinement scaling changes according to whether the fine-scale ansatz is

fulfilled or not.

Trapped-particle instabilities are predicted to arise in stellarators as well.³¹ Stellarators are interesting configurations for investigations of microinstability physics because: (1) The helical connection length is short, making collisionless trapped-particle modes more accessible. (2) Trapped particle fraction can be regulated through vertical field control of direct orbit losses.^{9,10} (3) The rotational transform profile is reversed from that of tokamaks. And (4), the weak magnetic shear in the center of stellarators will affect microinstability mode-structure. From the point of view of microinstability transport, stellarator research is just beginning.

The importance of fine scale turbulence rests both on sophisticated linear stability codes³² which show that tokamaks are invariably unstable to a variety of microinstabilities while remaining stable to ideal MHD perturbations and on the observation of density fluctuations of the anticipated magnitude and frequency.⁷³ Clearly these instabilities could have transport consequences and this has lead to a concentration of theoretical research on fine scale turbulence. The question yet to be definitively answered is: Is fine-scale turbulence the dominant transport mechanism?

If turbulence is fine-scale, then dimensional analysis of the equations governing electrostatic turbulence yield, for the case of electrostatic microinstabilities³³

$$\chi = C_1 \frac{T^{3/2} c^2 M^{1/2}}{e^2 B^2 L_n} g(v^*, \beta, q, s, \eta_i, L_n/R, r/a, r/R, \dots) \quad (7)$$

where g is in principle a function of all the nondimensional parameters of a plasma (see Table 5) except ρ_i/a . Equation (7) can be understood as a diffusion process wherein the step size is $\Delta r \propto \rho_i = (T/M)^{1/2} (Mc/eB)$ and the rate at which steps occur is taken to be proportional to the natural frequency of electrostatic turbulence $\omega \propto (T/M)^{1/2}/L_n$. To estimate transport associated with a particular microinstability, one usually must modify both the step size and rate by the nondimensional arguments of g .

An important special case arises when the plasma is both low- β and collisionless. In this situation, straightforward scaling arguments²⁰⁻²² yield a set of nondimensional, nonlinear equations for fine-scale microinstability turbulence which contain only non-

dimensional numbers of order unity, such as q, s, η_i , etc., and are independent of v^* , β , and p_i/a . Thus, in the low- β , collisionless limit, the function g in Eq.(7) is independent of the latter parameters and the explicit diffusivity of Eq.(7) governs the scaling of an entire class of microinstabilities. The remaining nondimensional parameters determine which of the several fine-scale instabilities - electron drift waves, ion temperature gradient modes, or trapped ion modes - is dominant. The constant C_1 is expected to be of order unity and relates the turbulence level to that given by the mixing length rule which states that turbulence saturates when fluctuating gradients are to the background gradient. Computations²⁷ of turbulent wavenumber spectra for microinstabilities remain in a state of flux. Theorists³⁴ have frequently used a provisional estimate of $\langle kp \rangle \approx 0.3$. One expects that both C_1 and g will vary depending on which fine-scale instability is dominant.

When diffusivity (7) is used in a simple 1-dimensional heat diffusion equation,³³ one can determine the energy content for a given auxiliary heating power. The confinement time follows directly and, for fine scale turbulence, takes the form

$$\tau_E = 1.1 \times 10^{-9} \text{ sec} \left(\frac{1}{C_1} \right)^{0.4} \frac{a^{2.4} R^{0.6} B^{0.8}}{A^{0.2}} \left(\frac{n}{P} \right)^{0.6} F(\bar{v}^*, \bar{\beta}, \bar{q}, \dots) \quad (8)$$

where all quantities are in SI units. F is a function of ill-defined averages of nondimensional parameters which, at present, may be more reliably determined by experiment than by theory. An important consequence of Eq.(7) is that, if nondimensional parameters remain fixed from one device to another, then the explicit factors of Eq.(7) govern confinement scaling. Confidence in performance projections improves because uncertainties in scaling with nondimensional variables are eliminated. This is exploited in the Discussion section.

Forms (7) and (8) are becoming known as gyro-reduced Bohm or electrostatic microinstability scaling. Figure 1 gives a representative fit plotted against L-mode tokamak and stellarator data. It is clear that Eq.(8) accounts for much of the variation in tokamak confinement times. Since the older and new fits are virtually identical, one can claim that microinstability scaling predicted the performance of L-mode TFTR and JET. The prediction is not as crisp purely empirical L-mode (see Fig. 2), but nonetheless sufficient to give hope that a local turbulence model is on the right track. Further,

stellarator and tokamak data are fit by the same formula, suggesting a common transport mechanism.

A key unresolved question is: Does electrostatic turbulence fulfill the fine scale ansatz? If it does not, then the parameter ρ_i/a can enter the arguments of Eqs.(7)-(8), changing confinement scaling. On the theoretical side, one needs to understand the long wavelength limit of trapped ion modes⁴² which could be affected by device size. Experimentally, new diagnostics will attempt to visualize density fluctuations. One can also measure confinement scaling versus B at fixed v^* , q , β , etc. to ascertain whether the exponents predicted by Eq. (8) are observed. Analyzing fast phase fluctuations on interferometer chords can indicate -- or set upper limits on -- low- m , large-scale density fluctuations ($m \leq 4$).

2.1.2 Empirical Scaling. The standard against which theoretical confinement scaling is judged consists of empirical regression formulas which express confinement time in terms of parameters which experimentalists can control: toroidal field, plasma current, heating power, plasma density, size and shape. These formulas fit L-mode data remarkably well and Fig. 2 presents four examples. Because each fit is superior to those of electrostatic microinstability theory presented in Fig. 1 and because of contradictions examined in Sec. 2.1.4, many experimentalists have concluded that empirical scaling is more reliable than theory in projecting performance of future devices.

Fig. 1. Tokamak and stellarator L-mode confinement data plotted against theoretical microinstability scaling, Eq.(8), with $(1/C_1)^{0.4}F = 1.0$. a) Older tokamaks, 1985 fit from ref 33. b) 1989 fit provided by R. Goldston with recent large tokamak data. c) Data from ATF stellarator provided by M. Murakami from "Overview of Recent Results from the Advanced Toroidal Facility", ORNL/TM - 11453.

Fig. 2. Four representative fits to L-mode confinement data (see Table 3). a) Goldston Aachen. b) Kaye-all-complex, c) Goldston Thermal, d) Bickerton. This figure provided by S. Kaye.

Empirical scalings face three difficulties which suggest caution in their usage. First, since empirical scalings are not based on physical principles, one cannot assess their range of validity. Indeed, empirical scalings are often not expressed in terms of the nondimensional parameters which elementary plasma physics principles suggest should govern confinement.^{19,20} Goldston-Aachen, Kaye-all-complex, and Goldston-thermal share this property. Second, measured confinement times obviously reflect the physical processes bearing on plasma heating as well as confinement. But the use of empirical scaling in confinement projections is to extrapolate confinement of only the thermal energy component. The question is: Do systematic trends in plasma heating profiles and creation of high energy particles with appreciable energy content introduce important scalings into empirical regression formulas? Third, the assumption of a power law dependence is appropriate when confinement is governed by a single mechanism. Theory suggests that different transport regimes exist, in which case power-law formulas experience difficulties.

Two formulas of Table 3 -- Bickerton MHD and Goldston Thermal -- attempt to remedy these defects. In addition, these formulas have been selected to demonstrate potential variability of empirical confinement projections all of which fit the data base well. Bickerton found that a formula based entirely on ideal MHD quantities fit the data well. His regression formula is

$$\tau_E = \begin{cases} (\text{const.}) (a/V_A) \beta_0^{-2.2} (a/R)^{-3} q^{1.2} \\ (0.198) I^{1.08} n^{0.16} A^{0.16} a^{0.7} R^{0.6} p^{-0.69} \end{cases} \quad (9)$$

where equivalent theoretical and practical versions have been given. Bickerton's formula is an example of a regression analysis which has been constrained to be in accord with elementary plasma physics principles, in this case ideal MHD. It is remarkable in its simplicity and accuracy. The paradox is that ideal MHD linear stability studies of tokamaks invariably show stability with respect to fine-scale, ballooning modes for β values below the Troyon limit.³⁵ Recently, Goldston devised a procedure for subtracting energy stored in a high energy tail from L-mode experimental data. The regression formula fit to confinement of the remaining thermal energy content, in the units of Table 3, is

$$\tau_E(\text{sec}) = (0.078) I^{0.35} B^{0.22} n^{0.24} R^{1.28} a^{0.29} A^{0.12} \kappa^{0.75} P^{-0.72} \quad (10)$$

Note differences in density and power scaling from the Goldston Aachen formula (3). The next step is to correct for changes in the power deposition profile caused by beam attenuation at high densities.

Regression analyses can potentially be used to identify the class of theory (i.e. the transport mechanism) best-suited to fit observations. The technique is to assume various functional forms for τ_E motivated by theory. Equation (8) is an example for fine scale electrostatic turbulence. The form which exhibits the best fit is inferred to be the theory of choice. Difficulties arise when the theory so identified is judged to be inapplicable on theoretical grounds. For example, comparison of Figs. 1a,b with Fig. 2d suggest ideal MHD is to be preferred over fine scale electrostatic turbulence, even though tokamaks are stable to ideal MHD modes. A successful identification of transport mechanisms must yield scaling relations which follow from a justifiable theory.

2.1.3 Power Loss and Diffusivity Formulas. A more transparent view of the physics and extrapolation properties implied by empirical confinement scalings results from re-expressing them in terms of power loss and diffusivity by eliminating power in favor of temperature through a simple heat balance equation

$$\frac{3nTV}{\tau_E(P)} = P \quad V = 2\pi^2 R a^2 \kappa \quad (11)$$

and by replacing plasma current with an elliptical model cross section formula

$$I_p = \pi a^2 B_T (1 + \kappa^2) / R \mu_0 q \quad (12)$$

Table 3 presents the results along with corresponding formulas for collisionless electrostatic microinstability scaling. It is particularly interesting to compare the power loss formulas with the expression for fusion energy production

$$P_{\text{fusion}} = (\text{const.}) n_D^2 T^2 (2\pi^2 R a^2 \kappa) \quad (13)$$

based on an approximate T^2 dependence of the DT fusion rate. Evidently, different scaling relations suggest different experimental operations to optimize the ratio $P_{\text{fusion}}/P_{\text{loss}} = nT\tau$. In contrast to conventional fusion thinking, the Bickerton and

Goldston thermal scalings optimize at very low density because of their steep dependence of P_{loss} on density. We shall return to this point following Eq. (14).

2.1.4 Fine-Scale Electrostatic Microinstabilities and Empirical Scaling: Confluences and Contradictions

To what extent are the scaling predictions of fine-scale electrostatic microinstability theory borne out by experiment? Let us focus on the diffusivity column of Table 3 for it is there that the physics implications of empirical scaling are most transparent. Furthermore, the results of nonlinear turbulence calculations usually relate heat and particle fluxes to temperature and density gradients via a diffusivity, even though the diffusivity itself may depend on gradients. Three areas of significant agreement lend support to microinstability scaling. First, the absolute value of global confinement times predicted by microinstability theory correspond to a curve (Fig.1) which runs through the middle of tokamak and stellarator L-mode data *with no normalizing constants*. Second, theory interprets confinement degradation with power as resulting from diffusivities increasing with temperature. Table 3 shows that empirical diffusivities have this property and that the temperature-dependence exponents are in good accord. Similarly, increased magnetic fields reduce diffusivities. Again, theoretical and empirical exponents are close. These successes suggest electrostatic microinstability theory is on the right track.

But fine-scale electrostatic microinstability predictions qualitatively differ from empirical scalings in five important areas: confinement variation with hydrogen isotopes, density dependence, size scaling of diffusivities, diffusivity profiles, and sensitivity of turbulence levels to plasma profile and edge effects. A discussion of these contradictions follow.

1. First, the strong positive density dependence of Eq. (8) predicted by microinstability for global confinement is not found in the experimental base, although the ASDEX team¹⁶ finds an exponent ≈ 0.3 in H-mode. Goldston thermal scaling Eq.(10) also exhibits a modest increase of confinement time with density. When coupled with power degradation, this results in empirical diffusivities that increase sharply with density. Alternatively, the empirical diffusivity vanishes in the $n \rightarrow 0$ collisionless limit. One should ask if this is because some other physics in addition to microinstability transport is masking an inherent positive density dependence. Some suggestions include:

Table 3. Empirical and Microinstability Confinement, Power Loss, Diffusivity, and Ignition Margin Scaling

Name	Confinement Law* $\tau_E(\text{sec})$	Diffusivity	Power Loss	Ignition Margin
Electrostatic Microinstability	$(0.28)a^{2.4}R^{0.6}B^{0.8}n^{0.6}$ $p^{0.6}A^{0.2}$	$\frac{T^{1.5}A^{0.5}}{B^2a}$	$\frac{nT^{2.5}A^{0.5}}{B^2}(\frac{R}{a})$	$\frac{n\kappa a^3B^2}{T^{0.5}A^{0.5}}$
Goldston Aachen	$(0.030)I^{1.75}\kappa^{0.5}A^{0.5}$ $p^{0.5}a^{0.37}$	$\frac{nTa^{0.24}}{B^2A}(\frac{a}{R})^{0.5}\frac{\kappa q^2}{(1+\kappa^2)^2}$	$\frac{n^2T^2a^{1.24}}{B^2A}(\frac{R}{a})^{0.5}\frac{\kappa q^2}{(1+\kappa^2)^2}$	$\frac{B^2(1+\kappa^2)^2a^{1.76}A}{q^2}(\frac{R}{a})^{0.5}$
Kaye all Complex	$(0.066)\frac{I^{0.85}R^{0.85}a^{0.3}\kappa^{0.25}n^{0.1}B^{0.3}A^{0.5}}{p^{0.5}}$	$\frac{n^{0.8}T R \kappa^{1.5}q^{1.7}}{B^{2.3}A(1+\kappa^2)^{1.7}}$	$\frac{n^{1.8}T^2R^2\kappa^{1.5}q^{1.7}}{B^{2.3}A(1+\kappa^2)^{1.7}}$	$\frac{n^{0.2}aB^{2.3}A(1+\kappa^2)^{1.7}}{\kappa^{0.5}q^{1.7}}(\frac{a}{R})$
Goldston Thermal	$(0.078)\frac{I^{0.95}B^{0.22}n^{0.24}R^{1.28}a^{0.29}A^{0.12}\kappa^{0.75}}{p^{0.72}}$	$\frac{n^{1.7}R^{1.4}T^{2.6}\kappa^{0.91}q^{3.4}}{B^{4.2}A^{0.44}a^{0.65}(1+\kappa^2)^{3.4}}$	$\frac{n^{2.7}R^{2.4}\kappa^{0.91}q^{3.4}T^{3.6}}{B^{4.2}A^{0.44}a^{0.65}(1+\kappa^2)^{3.4}}$	$\frac{B^{4.2}a^{2.7}A^{.44}\kappa^{0.1}(1+\kappa^2)^{3.4}}{n^{0.73}T^{1.6}R^{1.4}q^{3.4}}$
Bickerton	$(0.198)\frac{I^{1.08}n^{0.15}A^{0.16}a^{0.7}R^{0.6}\kappa^{0.06}}{p^{0.69}}$	$\frac{n^{1.8}T^{3.22}R^{3.77}q^{3.48}\kappa^{3.03}}{A^{0.5}B^{3.48}a^{2.77}(1+\kappa^2)^{3.48}}$	$\frac{n^{2.7}T^{3.22}R^{4.77}\kappa^{3.03}q^{3.48}}{A^{0.5}B^{3.48}a^{2.77}(1+\kappa^2)^{3.48}}$	$\frac{B^{3.48}a^{4.77}A^{0.5}(1+\kappa^2)^{3.48}}{n^{0.7}T^{1.22}R^{3.77}q^{3.48}}$
ITER	$(0.048)A^{0.5}I^{0.85}B^{0.2}n^{0.1}a^{0.3}R^{1.2}\kappa^{0.5}$ $p^{0.5}$	$\frac{n^{0.8}TR^{0.3}\kappa q^{1.7}}{AB^{2.1}(1+\kappa^2)^{1.7}}$	$\frac{n^{1.8}T^2R^{1.3}\kappa q^{1.7}}{AB^{2.1}(1+\kappa^2)^{1.7}}$	$\frac{n^{0.2}a^2AB^{2.1}(1+\kappa^2)^{1.7}}{R^{1.3}q^{1.7}}$

* Units: P in MW, I in MA, B in Tesla, n in 10^{20}m^{-3} , R, a in m, A in atomic mass units.

(1) counting stacked beam energy in evaluating plasma energy, (2) plasmas with strong NBI ion heating (which raises T_i above T_e) in cases where electron heat loss dominates, (3) neutral beam penetration effects, (4) operational correlation of $\eta_i I_p$, etc. with density so that effects of density are mistakenly attributed to $\eta_i I_p$ and (5) a β -dependent confinement degradation mechanism. The change in density scaling between Goldston-Aachen and Goldston-thermal demonstrates that item (1) influences empirical scaling. Items (2-5) await definitive studies. Clearly, one expects a low- β limit to exist, in contrast to the empirical diffusivities of Table 3.

If one uses the freedom to introduce additional density scaling via the parameters v^* and/or β which are arguments of the undetermined function F and thereby alter the density dependence of (7) and (8), then the power degradation and magnetic field scalings, which do agree with experiments, are changed [see Eqs. (20,21)].

The empirical stellarator confinement scaling relation³⁶

$$\tau_E(\text{sec}) = 0.17 \left(\frac{P}{1\text{MW}} \right)^{-0.58} \left(\frac{n}{10^{20} \text{ m}^{-3}} \right)^{0.69} \left(\frac{B}{\text{Tesla}} \right)^{0.84} \left(\frac{a}{1\text{m}} \right)^{2.0} \left(\frac{R}{1\text{m}} \right)^{0.75} \quad (14)$$

finds the density exponent to be just that predicted by Eq. (8). Indeed, all exponents in Eq. (8) are quite similar to those observed in stellarator experiments. Results from JFT-2M tokamak ICRF heating experiments³⁷ agree with the theoretical density exponent, but differ with regard to power degradation.

Several empirical scalings of Table 3 have the property that power losses increase so rapidly with density (faster than n^2) that one can optimize ignition margin by operating at very low density. The algorithm for predicting CIT performance precludes this optimization by combining the auxiliary heating confinement time scaling law (as given in Tables 3 and 4) with the neoAlcator confinement time according to the formula

$$\tau_E = \min\{\tau_{E,\text{neoAlcator}}, \tau_{E,\text{aux}}\} \quad (15)$$

Here

$$\tau_{E,\text{neoAlcator}} = (0.07) n R^2 a q \quad (16)$$

in the units of Table 3. Empirically, Eqs.(15),(16) state that tokamaks can never achieve a

confinement better than ohmically-heated plasmas. Several questions can be raised concerning this procedure. First, if $\tau_{E,aux}$ represents a valid scaling law, why doesn't it give the appropriate ohmic confinement when ohmic heating power is used? Second, formula (15) implies there are two confinement regimes, one characterized by no degradation of confinement with power. What physics principles govern the boundary? Third, why should the *neoAlcator* scaling law not have a power degradation factor? The reason that ohmic confinement formulas lack a power degradation factor is because power cannot be independently controlled in ohmic experiments and not because physics principles preclude power degradation should power be increased over ohmic. Since the ohmic heating power P_{Ω} is given by

$$P_{\Omega} = (\text{const.}) B^2 Z_{eff} a^2 (q T^{3/2} R)^{-1} \quad (17)$$

one can construct a modified *neoAlcator* scaling

$$\tau'_{E,neoAlcator} = (0.07) n R^2 a q \left(\frac{P_{\Omega}}{P} \right)^{\nu} \quad (18)$$

where ν is an arbitrary exponent. The modified scaling law will, by construction, yield the correct confinement time for ohmically heated discharges. An alternative algorithm to Eqs. (15),(16) would be to match the exponent ν to the power degradation exponent in the auxiliary scaling law, thereby declaring that confinement degradation with power is a common feature in both regimes. Altogether, the algorithm for optimizing ignition margin is overly arbitrary when the empirical $nT\tau_E$ scaling has a strong density dependence. High power, low density confinement experiments are called for to clarify confinement scaling in the limit $n \rightarrow 0$. Furthermore, ohmic plasmas, together with their ohmic power, should be included in the tokamak data base for assessing confinement. A fundamental principle in our approach to confinement is that that heating method does not affect confinement physics.

2. The atomic mass dependence of Eq. (8) is opposite to the trend of observations which show that deuterium confinement is better than hydrogen confinement [F. Wagner, presentation at Ringburg Seminar(1989)]. While the extrapolation from pure deuterium plasmas to DT mixtures is not large, the isotope problem is very troubling to theorists because one expects proton plasmas, with their smaller gyroradius, to have better

confinement. Why experiments differ from such a simple principle is a real enigma, and the isotope effect constitutes a key objection to microinstability scaling. One should remark that, while an isotope effect is not always observed,³⁸ confinement in hydrogen discharges does not exceed that in deuterium, in contradiction to the smaller gyroradius principle.

Of the nondimensional quantities which can potentially control local turbulence levels, only two contain the ion mass. The first compares the effective electron collision frequency with typical drift wave frequencies

$$\frac{\nu_{e,eff}}{\omega_*} = \frac{1}{(k_\theta \rho_s)} \left(\frac{M}{m} \right)^{0.5} \left(\frac{r^{0.5} L_n}{R^{1.5} q} \right) v_e^* \quad (19)$$

This quantity enters into the transition from collisionless trapped electron modes to dissipative trapped electron modes and eliminates the theoretical isotope effect for the dissipative trapped electron mode, but not from collisionless modes. Further use of parameter (19) to give an isotope effect in the sense of observations would imply very severe confinement degradation with temperature which is not observed. Density scaling discrepancies would be aggravated as well. The second nondimensional parameter is $\rho_i \omega_{pe}/c = (M\beta/m)^{1/2}$. Use of this parameter to give the observed isotope effect would imply confinement improving with β , an unlikely prospect. Again, it appears that the resolution lies outside microinstability physics. Experiments with detailed n, T_e, T_i profiles at fixed B_T, I_p, n are needed separately for D^+, H^+ plasmas in OH, L- and H-mode conditions to attempt to understand whether electron-ion heat transfer, atomic physics, plasma fueling, NBI species, external neutral pressure, etc. can provide an understanding of this enigma, either directly or through their influence on profiles which, in turn, affect turbulence. Experimentally, the existence of enhanced confinement regimes suggests confinement is quite sensitive to profile details which may be inadequately measured. Empirical projections of confinement exploit confinement improvement with atomic mass in estimating ignition margin. Numerically, the effect is small, but it could play a role if ignition margin is close to unity, as is the case for CIT and ITER.

3. Microinstabilities arise from the presence of a pressure gradient in equilibrium. Consequently, the corresponding diffusivities depend on gradients. (While the density gradient appears explicitly in Eq. (7), it could be transformed into the temperature gradient

through the nondimensional parameter η_i .) Fine-scale microinstabilities obey the principle: the steeper the gradient, the higher the diffusivity. It is quite surprising then that, according to Table 3, empirical diffusivities exhibit the opposite trend. All empirical diffusivities contain machine size to a positive power which, in effect, states the gentler the gradient, the higher the diffusivity. Alternatively, this result could be understood in terms of a convection process or a diffusion process where the step size increases with machine size, i.e., a process which violates the fine-scale ansatz. Because of this scaling, empirical confinement projections for large machines like ITER are pessimistic with respect to theory. Again, one can question whether plasma heating physics, in particular NBI penetration effects, are masking a stronger positive size dependence. Also, since power is correlated with size (larger tokamaks have higher heating power), an underestimate of the power degradation exponent would lead a regression analysis to underestimate the positive size-scaling exponents. To test the latter hypothesis, Fig. 3 and Table 4 display two empirical fits, contributed by S. Kaye, of confinement time to the tokamak data base, the first unconstrained and the second constrained to have $(\text{length})^3$ scaling. It is clear that the fits are almost equivalent. But formulas for confinement time, power loss and $nT\tau_E$, given in Table 4, differ appreciably. The difference in predicted impact on tokamak optimization is dramatic. At fixed β , the unconstrained fit indicates $nT\tau_E$ improves rapidly with magnetic field and weakly with machine size, favoring CIT-like devices. With the constrained fit, $nT\tau_E$ increases rapidly with machine size and weakly with field, suggesting a large device such as ITER. Both fits share the bizarre trend of many empirical scalings to optimize $nT\tau_E$ at very low β , which can be traced to the fact that τ_E does not increase strongly with density. Indeed, none of the empirical diffusivities of Tables 3 and 4 show a clear collisionless limit as $n \rightarrow 0$. Experiments involving a series of discharges with auxiliary heating and *decreasing* density are required to establish an experimental low- β , low collisionality scaling for confinement and diffusivity in tokamaks.

90T0007

Fig. 3. Two fits to the L-mode tokamak data (see Table 4). a) Unconstrained fit "Kaye-Simple". b) Fit constrained to have $(\text{size})^3$ scaling. This Figure provided by S. Kaye.

Table 4. Two Empirical Fits to the Tokamak Data Base.*

Quantity	Unconstrained	Constrained†
$\tau_E(\text{sec})$	$(0.098) \kappa^{.52} I^{1.05} B^{.15} a^{.31} R^{.73} p^{.58}$	$(0.091) \kappa^{.46} a^{.45} R^{1.71} n^{.21} I^{.84} p^{.74}$
rms error	13%	16%
Power Loss	$\frac{n^{2.38} T^{2.38} R^{3.14} \kappa^{1.14} q^{2.5}}{B^{2.86} a^{.98} (1+\kappa^2)^{2.5}}$	$\frac{n^{3.03} T^{3.85} R^{.54} q^{3.25} \kappa^{.21}}{B^{3.25} a^{.54} (1+\kappa^2)^{3.25}}$
Diffusivity	$\frac{n^{1.38} T^{1.38} R^{2.14} \kappa^{1.14} q^{2.5}}{B^{2.86} a^{.98} (1+\kappa^2)^{2.5}}$	$\frac{n^{2.03} T^{2.85} q^{3.25}}{B^{3.25} a (1+\kappa^2)^{3.25}} \left(\frac{a}{R}\right)^{0.46}$
$nT\tau_E$	$\frac{B^{2.86} a^{.84} (1+\kappa^2)^{2.5}}{(nT)^{.38} \kappa^{.14} q^{2.5}} \left(\frac{a}{R}\right)^{2.14}$	$\frac{B^{3.25} R^{.46} a^{2.54} (1+\kappa^2)^{3.25}}{n^{1.03} T^{1.85} \kappa^{1.1} q^{3.25}}$

* Contributed by S. Kaye. Units are the same as Table 3.

† Size scaling is (length)³; this includes the linear size scaling contained in I.

Points 1 -3 above constitute fundamental scaling discrepancies between a entire class of microinstabilities and experimental data. If they cannot be resolved by a reinterpretation of the data, then a major redirection of theoretical confinement research, which has focussed on fine-scale electrostatic turbulence, must be undertaken. One could ask whether a judicious introduction of the nondimensional partameters v^* and β could bring microinstability scaling into accord with empirical scaling. Let us compare exponents in the formulas for power loss from Table 3

$$\text{Empirical Power Loss} \propto \left(\frac{n^2 T^2 a^{1.24}}{B^2 A} \right)_{\text{Goldston}} \quad (20a)$$

$$\text{Fine Scale Turbulence Power Loss} \propto \left(\frac{n T^{2.5} A^{0.5}}{B^2} \right) (v^*)^x \beta^y \quad (20b)$$

Can one choose exponents x, y to bring these formulas into agreement? The nondimensional parameters scale as

$$v^* \propto \left(\frac{A^{0.5} n a}{T^2} \right) \quad \beta \propto \frac{nT}{B^2}$$

where the (optional) A scaling in v^* arises from (19). There are five exponents to match and only two free parameters. Matching the B scaling requires $y=0$ and matching the n scaling gives $x=1$. One recognizes the resulting scaling as classical (or neoclassical, depending on multiplicative nondimensional parameters) with a diffusivity that falls with increasing temprature. Clearly, temperature and A scalings then have serious discrepancies. Fine scale electrostatic turbulence theories will not reproduce Goldston-Aachen scaling, no matter how sophisticated the analytic or numerical approach. And, Table 4 shows that when size scaling is constrained to reproduce fine-scale turbulence, power degradation and density exponents differ from Eq. (8).

4. Confinement time formula (8) exhibits a positive scaling with toroidal field, not with plasma current as is reported by virtually all tokamak experiments (see Table 3). A closely related discrepancy is the failure of diffusivity (7) to increase outside the $q \approx 2$ surface which is also a universally observed phenomena.³⁹ Some theorists have

responded⁴⁰ by introducing an empirical dependence of the function g on q in Eq. (7) which remedies this deficiency and happily depends only on a non-dimensional parameter $-q$ which, in principle, could affect turbulence levels. But attempts to derive such a dependence from nonlinear turbulence calculations have only recently been successful. Diffusivities with a positive q -dependence have been found in nonlinear computations of ITG mode turbulence⁴¹ and long wavelength ($k_{\parallel} < 0.1$) microinstabilities where neoclassical effects enter.⁴² Moreover, linear growth rates of collisionless electron modes increase with shear which itself increases appreciably between the $q=1$ and $q=2$ surfaces. Evidently, theoretical progress is being made and these works deserve further study. Experimental determination of whether shear exists within the sawtooth mixing region is an important input to microinstability diffusivity profiles. Thus the resolution of this discrepancy could lie within microinstability physics and calls for the ability to deduce analytically or computationally how turbulence levels and transport will respond to changes in non-dimensional quantities of order unity. Experiments should attempt to measure how turbulence levels depart from the mixing length value $\tilde{n}/n = 1/k_{\perp} L_n$ as various non dimensional parameters are changed, e.g., q , s , L_n/R , etc.

Heat pulse propagation studies⁴³ constitute a useful technique for studying diffusivity profiles and plasma transport properties in general. Most investigations have relied upon heat pulses following sawtooth crashes, and have found the diffusivity associated with thermal pulse propagation appreciably exceeds the steady-state power-balance diffusivity. But, the complications of sawteeth phenomena create uncertainties in interpretation. Fortunately, pulsed ECH heating should eliminate these complications and also provide for spatial control of the heated volume, greatly improving this technique.

5. Enhanced confinement regimes imply quite significant changes in the constant C_1 which is a measure of the turbulence level. Simple algebra [cf Eq. (8)] shows that a representative H-mode enhancement of confinement time by a factor-of-2 over L-mode corresponds to a factor of 5.6 reduction in C_1 (and to a factor of 5.6 increase in ignition margin). For empirical Goldston-Aachen scaling, a factor-of-2 in confinement time translates into a factor-of-4 in diffusivity. In ASDEX¹⁸ and TFTR,¹⁷ steepening the density profile improves confinement. Qualitatively, this implies that we can potentially - and in reality - exert a very significant control over turbulence levels. Within the context of fine-scale turbulence, the change in local turbulence levels must result from changes in

local nondimensional quantities, in particular the shape of the density profile which changes L_n/a , η_e , η_i , $\epsilon_n = L_n/R$, etc. Thus, identification of fine-scale electrostatic turbulence as the dominant transport mechanism calls for experiments that show turbulence levels responding to changes in profile shapes and other non-dimensional parameters. ITG mode theories suggest T_i/T_e is an important nondimensional parameter. As mentioned above, sensitivity to profile shapes could also be the cause of the isotope effect. A contrasting observation would be that core H-mode diffusivities respond instantaneously to the edge H-mode transition, even though local profiles have not been altered, and would serve to favor a non-local marginal stability state over saturated local microturbulence. Validation of local fine-scale turbulence calls for significantly improved nonlinear studies - both analytic and computational - on the saturation levels of collisionless trapped electron modes and ion temperature gradient modes. We should understand theoretically why transport responds to nondimensional parameters.

To summarize, validation of fine-scale turbulence scaling appears to require a reinterpretation of experimental data to address difficulties 1-3 and appreciably improved nonlinear theories to resolve points 4 and 5.

2.5 Identifying the Principal Electrostatic Micronistability.

Electrostatic instabilities come in a variety of forms -- fine-scale electron drift modes, both collisionless and dissipative, ion temperature gradient modes, as well as large-scale neoclassical trapped-ion modes. To the nonspecialist, there may appear to be an ad-hoc profusion of many modes with little in the way of unifying concepts. Unfortunately, jargon has obfuscated an underlying simplicity. Ion temperature gradient modes and electron drift waves are different eigenfunctions of the same equation. Collisionless, dissipative-trapped-electron, and resistive drift waves differ only in the equation chosen to represent the response of electrons to the parallel electric field and consequently depend on mean free path. Trapped-ion modes spring from the appropriate modification of these equations to describe ion response when the wave frequency is below the bounce frequency for magnetically mirroring ions. Resistive MHD and electromagnetic drift waves add magnetic fluctuations and induced electric fields arising from the parallel current density to these basically electrostatic, pressure-gradient-driven waves.

Nondimensional parameters govern which mode dominates and its transport

consequences. Unfortunately, the list of independent nondimensional parameters which can potentially govern turbulence levels is long. Table 5 lists 17 local, nondimensional parameters. The fine-scale ansatz eliminates only one of these - p_i/a . The others all enter the equations governing microinstability turbulence in ways whose scaling consequences are not transparent. Experimentalists can vary these parameters over meaningful ranges and, as mentioned above, have found variation in confinement time that imply very significant changes in microinstability turbulence. Furthermore, Table 5 lists only those nondimensional *plasma* parameters which can potentially affect turbulence levels and transport fluxes. Plasma parameters must be determined through a self-consistent solution of the transport equations which can depend on still more non-dimensional parameters, often involving atomic physics, such as λ_i/a and $F_0 I_0/q$ where λ_i denotes the ionization mean free path, F_0 the neutral particle influx, I_0 a typical ionization energy and q the outgoing heat flux. For example, the density profile shape, which influences turbulence levels through the parameters L_n/a and L_n/R , is surely affected by the plasma fueling profile — central NBI fueling in the case of TFTR supershots and edge fueling for DIII-D H-modes.

The long list of nondimensional parameters is what makes tokamak transport complex and, along with the fine-scale ansatz, stands in contrast to hydrodynamic turbulence where a wide "inertial" range of spatial scales occurs but the number of nondimensional parameters is small. The fine-scale hypothesis restricts the range of spatial scales in plasma turbulence to approximately a factor of 100 — $3 < k\rho < 0.03$ — a range amenable to simulation studies. Tokamak plasma turbulence simulations, like computational aerodynamics, must deal with inhomogenities and complex geometry, but are not plagued by the enormous range of spatial scales characteristic of hydrodynamic turbulence. Fundamentally, plasma turbulence occurs in a 5-dimensional phase space, but 3-dimensional fluid models are quite informative. It will require sophisticated analysis and computation to address the influence of nondimensional parameters theoretically.

Table 5. Nondimensional Plasma Parameters Governing Microinstability Turbulence

Parameter	Definition	Comment
ρ_i/a		gyroradius size; very small
L_n/a	$n/[a(dn/dr)]$	density profile shape
η_e	$n(dT_e/dr)/[T_e(dn/dr)]$	measures ∇T_e
η_i	$n(dT_i/dr)/[T_i(dn/dr)]$	measures ∇T_i
n_I/n_e		impurity concentration profile
a/R		aspect ratio
T_e/T_i		temperature ratio
A		atomic mass (AMU)
ν^*	$(R/r)^{3/2}(qR\nu_{ei})/(2T_e/m)^{1/2}$	collisional mean free path
β	$16\pi L_n T/B^2$	plasma pressure
q	rB_T/RB_θ	poloidal field
s	rq'/q	measure of current density
m/M		mass ratio
r/a		minor radius
$\omega R (2T/M)^{-1/2}$		toroidal rotation
$\omega' L_n (Mc/eB)$		toroidal rotation shear
κ , separatrix		elongation, shape

Even within the context of electrostatic turbulence, it is not known experimentally which modes are the dominant ones. Most theoretical research has concentrated on fine-scale trapped electron and ion temperature gradient modes which one presumes to have $k_{\perp}\rho \sim 0.3$. Laser and microwave coherent-scatter experiments are limited to the region $k_{\perp}\rho > 0.3$, and find the fluctuation spectrum still rising as wave number decreases as portrayed by Fig. 4. Since very long wavelengths, $m \leq 4$, would appear as rapid phase fluctuations on interferometer chords, which measure $\int n_e dl$, and these are not reported, one concludes a spectral peak exists.

At low ν_i^* , trapped-ion modes, which have $\omega_{bi} > \omega^* > \nu_{i,eff}$ and consequently $0.1 > k_{\perp}\rho > 0.03$, are also a possibility. Trapped ion modes in principle are large-scale modes where the diffusion step-size can scale with device size. Consequently, theoretical research is giving renewed attention to trapped ion modes,⁴² where neoclassical processes enter and the poloidal field plays an important role, as it must to understand empirical I_p -scaling.

It is crucial to experimentally determine where the spectrum peaks, so theorists can have some idea of what the dominant mode is and whether the spectral peak depends only on $k_{\perp}\rho$, as the fine-scale ansatz holds, or is affected by machine size. Heavy ion beam probe measurements,⁴⁴ although limited by the ambiguities of two point observations, suggest that the spectrum peaks at $k\rho \approx 0.1$. There are a variety of incoherent X-ray and optical imaging experiments - beam emission spectroscopy, parallel X-ray and visible bremsstrahlung imaging, 2-dimensional Thomson scatter imaging - as well as multi-detector heavy ion beam probe methods - which can potentially measure fluctuations to 1% accuracy. These should produce fascinating visualizations of tokamak density fluctuations in the regime $0.3 > k_{\perp}\rho > 0.03$. High priority is warranted for imaging experiments.

A rather detailed comparison is possible between measured heat and particle fluxes and the fluxes inferred from measured fluctuations, frequencies, etc. via quasilinear microinstability formulas with heavy ion beam diagnostics. Currently, agreement is marginal at best.⁴⁴ Development of multi-detector heavy-ion beam probe capability, which can eliminate the ambiguities of two-point measurements, coupled with improved

Fig. 4. Microwave scatter density fluctuation measurements. a) PDX data from T. Crowley and E. Mazzucato, Nucl. Fusion **25** 507 (1985). Open symbols denote an outside scattering volume and solid symbols an inside scattering volume of a diverted plasma. b) Unpublished TFTR Helium L-mode data provided by N. Bretz. These spectra assume a width $\Delta k_{\parallel}$ in parallel wavenumber space.

nonlinear theories is needed to resolve these discrepancies (or eliminate microinstabilities as the dominant transport mechanism).

Electron drift waves and ion temperature gradient modes differ in their source of instability. Electron cyclotron emissions (ECE) can potentially indicate trapped electron response, which drives drift waves, by contrasting inside and outside propagation paths. The principle here is that, because of large optical depth, fundamental ECE emission outward in major radius arises principally from electrons with a small relativistic shift which are highly collisional and hence unaffected by trapping. Emission inwards comes from high-energy trapped electrons. The fluctuating distribution function associated with drift waves cannot be characterized as simple temperature fluctuations. The role of trapped particles in driving fluctuations can also be investigated in stellarator experiments that eliminate trapped particles via direct orbit losses in certain configurations. (Direct losses will adversely affect confinement, however.) Ion temperature gradient modes depend principally on parallel ion motion for destabilization. Observation of short-scale density fluctuations in plasmas with $\nabla n_e \approx 0$, which eliminates cross-field advection as a source of fluctuations, would isolate the parallel compressibility destabilization of ion temperature gradient modes.

In summary, additional measurements of density fluctuations in the wave number range $0.3 > k\rho > 0.03$ ($40 > m > 4$ in TFTR), now probed only by heavy ion beams, are needed to identify which electrostatic instability is dominant as a function of nondimensional parameters and whether the fine-scale ansatz holds.

2.6 Transport by Magnetic Fluctuations. Let us turn next to the question of transport by magnetic braiding, wherein magnetic fluctuations destroy the picture of nested toroidal flux surfaces and permit electrons to stream along magnetic field lines from hotter to colder regions. While initial investigations⁴⁵ of magnetic braiding simply assumed fluctuation amplitudes, it is well-known theoretically that fine-scale magnetic fluctuations are inevitably associated with electrostatic $\mathbf{E} \times \mathbf{B}$ turbulence.^{42,46} These fluctuations arise from parallel currents associated with electrons streaming along magnetic field lines to neutralize ion fluctuations caused by $\mathbf{E} \times \mathbf{B}$ drifts. Development of a self-consistent theory of magnetic braiding transport is just beginning. The difficulty lies in combining the topology of temporally-evolving, fine-scale magnetic islands (not a quasilinear concept)

with collisionless electron dynamics. Static magnetic islands are another possibility. Rebut and Lallia⁴⁷ as well as White and Romanelli⁴⁸ argue that steady-state fine-scale magnetic islands, with a spatial spectrum centered on toroidal mode number N , can be self-consistently maintained via ergodicity effects on current density. A representative formula for the diffusivity

$$\chi_e \approx (0.5) \frac{\overline{v}_e R}{R^2 s^2 q^2 N^5} r^2 \quad (21)$$

exhibits great sensitivity to N which, according to Rebut and Lallia, is determined via nonlocal electron acceleration in a completely collisionless ergodic region. But, in present and planned tokamaks, electrons must collide many times in traversing the ergodic region, given the estimated island sizes and realistic mean free paths. The ergodic region is not collision-free and the argument that static islands can be maintained encounters difficulties. Dynamic islands, associated with microinstabilities, are quite possible, as discussed above. Bootstrap currents at rational surfaces are yet another mechanism to maintain islands.⁴⁹ For fixed N , formula (21) yields a diffusivity which increases with machine size, a feature of many empirical scalings (see Tables 3 and 4). Picturesquely, the tokamak is wired incorrectly and the step-size in the diffusion process is tied to machine size. It is also interesting to note that this diffusivity is independent of magnetic field strength because its physics rests only on parallel velocity and island size. It decreases rapidly towards the plasma boundary, whereas experimental diffusivities rise. Let us note that there is no direct theoretical connection between fine-scale islands and Rebut-Lallia confinement scaling,⁵⁰ which is a strictly empirical scaling that obeys Connor-Taylor constraints.²⁰ Yet another source of magnetic fluctuations, very short wavelength electron modes with $kp \gg 1$, has recently been proposed as a cause of magnetic fluctuation transport.⁵¹ Here too, the diffusivity is independent of magnetic field strength. Microwave and laser scatter experiments should be able to observe the associated high-wavenumber density fluctuations.

How can experiments distinguish between electrostatic transport and magnetic braiding? All approaches exploit the asymmetry in electron and ion transport arising from the disparity in parallel velocity, the magnetic fluctuations themselves being far too small to observe. One measure of magnetic fluctuation transport is the ratio $\alpha_e = q_e / T_e \Gamma$, where q_e is the electron heat flux, and Γ the ambipolar particle flux. If this ratio is large ($\alpha_e > 5$), then magnetic fluctuation transport is a possibility, but not proven. If this ratio is

small (the lowest value is $\alpha_e \approx 3/2$), then electrostatic $\mathbf{E} \times \mathbf{B}$ transport dominates. Magnetic fluctuations can also be probed by techniques which create spatially-localized bursts of superthermal, circulating electrons whose subsequent diffusion can be monitored by synchrotron emission and x-ray tomography techniques. If magnetic braiding is important, then the diffusivity of superthermal test particle electrons should exceed the thermal diffusivity χ_e . Pulsed oblique ECH and lower hybrid are potential excitation techniques. Measuring the diffusion of positrons introduced at the plasma periphery via their annihilation on a pellet is an alternative and innovative approach to this problem.

Helical convective cells, locked MHD modes,²³ and magnetic islands resulting from error fields or neoclassical currents at low-order rational surfaces constitute yet another potential cause of tokamak transport and can be viewed as an alternative to the fine-scale turbulence model, or as a process⁴⁰ which degrades confinement for $q \geq 2$ where low-order rational surfaces are close. It follows that error fields are a potential resolution of discrepancy 4 in Section 2.1.4. Experiments commonly find degraded confinement in connection with locked MHD modes. Furthermore, TFTR Thomson scatter data suggests magnetic islands at many low order rational surfaces.⁵² Theoretically such structures appear viable, and error fields are an inevitable consequence of construction tolerances. The issues here are: What is the extent of such islands and do they satisfy an island overlap criteria? Higher accuracy Thomson scatter data is needed. The theory of such structures falls into the domain of neoclassical MHD²⁴⁻²⁶; this theory needs to progress into the nonlinear simulation area. The challenge here is to reliably and self-consistently estimate magnetic island widths in an environment where neoclassical transport and stochastic electron orbits are influencing magnetic fluctuation evolution.

2.7 The Role of Edge Turbulence. Experiments have shown that the plasma edge is a region where effective control over transport can be exercised. Two enhanced confinement modes - H-mode and supershot,¹⁷ peaked profile¹⁸ modes - are triggered by edge physics. Furthermore, Langmuir probe measurements, which are possible only in the edge region, provide a much more detailed picture of the connection between fluctuations and transport, even though edge plasma physics instabilities are quite different than core plasma instabilities. Also, imaging⁵³ of the edge region in TFTR clearly shows plasma striations with a dominant wavenumber $k_{\theta\rho} \approx 0.02$ ($m=100$). In this sense, the edge constitutes an interesting testing ground for the theory of turbulent

plasma transport.

Let us first turn to the physics of the H-mode transition^{16,29,54} - a high leverage problem in tokamak transport. An empirical case can be made that the H-mode transition corresponds to the stabilization of microinstabilities in the vicinity of a magnetic separatrix. The evidence^{16,29,54} consists of (1) observation of an edge transport barrier in H-mode, (2) reduction of high frequency noise in reflectometry measurements at H-mode onset, and (3) development of very steep density and temperature profiles in the separatrix region. An experimental characterization of the instability which stabilizes at the transition is needed. Observational resources which can be brought to bear on this issue include (1) parallel imaging interferometry at the separatrix, (2) imaging reflectometry, (3) high spatial resolution H_α toroidal tangential imaging, (4) extensive poloidal magnetic probe arrays, and (5) heavy ion beam probes.

Although enhanced confinement via H-mode appears restricted to elongated (preferably separatrix) plasmas, H-mode edge transitions occur in TFTR⁵⁵ at high β_θ and large heating power (11-25MW). This transition does not improve confinement beyond its already-enhanced supershot value. The common element between separatrix and limiter configurations remains elusive, although the requirement of large β_θ in TFTR suggests that a significant Shafranov shift is needed. Perhaps a large variation in local shear is the key.

Experimental measurements of the equilibrium near a separatrix are also needed. Extremely high temperature and density gradients develop just inside the separatrix.⁵⁶ DIII-D observations^{57,67} show an inward radial electric field and sheared poloidal flow developing at the transition, but the cause/effect relationship is not yet clear. Also, the effect on main chamber neutral gas pressure on the transition (more is unfavorable) needs clarification.

Recent experimental and theoretical⁵⁸ work points to the role of the radial electric field and sheared poloidal rotation in the H-mode transition. Shaing and Krume⁵⁸ show that a bifurcation in poloidal rotation velocity results when the current required to neutralize orbit losses exceeds a critical value. Biglari *et al*⁵⁸ argue that the shear in poloidal rotation will suppress microinstabilities, in qualitative agreement with

observations. Both these papers work in simplified geometry and the role of the separatrix or local magnetic shear, as well as the trigger and threshold power which starts this process, remain unclear. A generalization of these works to evaluate separatrix and local shear effects is clearly important.⁵⁹ One should note that poloidal rotation is absent from Table 5 of nondimensional parameters. Neoclassical theory shows that poloidal rotation will be very small in the bulk plasma. But a sheath can arise within a poloidal gyroradius of the separatrix.⁶⁰ Let us note that sheaths and sheared poloidal rotation also occur⁶⁰ without H-mode transitions, suggesting that poloidal rotation is a necessary, but not sufficient, phenomenon to trigger H-mode transitions. While radial electric fields obviously suggest a sheath process driven by particle orbit losses, experiments have not established that sheaths are the key to the H-mode transition. Very high resolution diagnostics of the separatrix region are needed.

How will the threshold power for H-mode transitions scale? Hinton⁶⁰ argues that the power flux through a sheath varies as

$$P_s = (\text{const.}) 4\pi^2 R a n T_{\text{edge}} v_D \cdot \min(1, (\text{const.}) v^*) \quad (22)$$

where $v_D = T_{\text{edge}} c / e B R$ denotes particle magnetic drift velocity. In the collisionless limit, this power is independent of edge temperature. Observations^{16,57,67} indicate that the edge temperature must exceed a critical value ($T_{\text{edge}} > (\text{const.}) B_T$) to effect an H-mode transition. The physics governing the transition temperature remains to be identified, although one can form two expressions for temperature from the nondimensional parameters of confinement theory

$$T_1 = C_1(q, s, R/a, \dots) \left(\frac{\beta}{v^*} \right)^{1/3} (e^4 a B^2 Z_{\text{eff}} \ln \Lambda)^{1/3}$$

$$T_2 = C_2(q, s, R/a, \dots) \rho^{*2} B^2 e^2 a^2 / M c^2$$

where C_1 and C_2 are undoubtedly functions of magnetic geometry, aspect ratio, etc. Thus observations suggest there is a critical value of $\rho^{*2} \beta / v^*$ (or some similar combination) associated with H-mode onset. If $T_{\text{edge}} \propto B_T$, then Eq. (22) provides an interpretation for the DIII-D results^{57,67} that the power threshold scales as $n B_T$, since the edge of DIII-D lies in the high v^* regime. These scalings indicate a high edge

temperature and threshold power for H-mode in a high-density, high- B_T device such as CIT.

Theory in the vicinity of a separatrix is difficult and relies heavily on computation. Both MHD and neoclassical kinetic studies of a separatrix equilibrium with an arbitrary $J_{||}$ profile are needed. A sheath structure in the kinetic case,⁶⁰ as well as a self-consistent radial electric field and sheared poloidal rotation, is expected. Ionization is of demonstrated importance, as explained below. Given the equilibrium, then the physics of the H-mode transition may well be found in linear theory if the hypothesis that the transition corresponds to a stabilization of a linear mode is correct. If only linear theory is needed, prospects for a physics understanding of the transition are reasonable. But it could be that stabilization is an artifact, and not an essential element, of the H-mode transition. Ignition devices plan on a separatrix; if for no other reason, this dictates that development of separatrix theory should enjoy high priority.

Successes at JET with beryllium gettering,⁶¹ as well as the wall-conditioning requirements for supershots, point to the need for developing a scientific basis for the effect of plasma-wall interactions on confinement. Indeed, there is concern that the molybdenum walls of Alcator C Mod may produce a different confinement physics than the graphite tiles planned for CIT, especially since ICRF heating has a proclivity towards deleterious plasma-wall interactions. Therefore a material science basis for particle efflux from material walls and its response to various conditioning techniques should constitute an essential element of tokamak transport research. A parallel theoretical development of the role which ionization plays in the edge equilibrium is needed. Experiments show that reductions in gas fueling rate can induce a transition to a steep-density-gradient, improved confinement regime.¹⁸

Studies of edge plasma turbulence in machines such as TEXT⁴⁴ and University tokamaks constitute an interesting testing ground for turbulent transport because probe measurements provide much more detailed observations regarding the turbulence, especially the crucial phase-shift between density and potential fluctuations.⁶² The role of these studies is to serve as a paradigm that theory and observation do coincide where they can be measured, thus engendering confidence in turbulence theories of core transport. Edge studies should also attempt to clarify why low-recycling walls can trigger peaked density profiles with enhanced confinement.

Recent theoretical developments in edge microinstabilities have pointed out the destabilizing role of radiative losses.⁶³ Such a thermal-instability driving mechanism is needed in light of experiments⁵⁴ which have shown that edge turbulence is insensitive to the value of the toroidal electric field, ruling out the classic rippling mode. If fine-scale instabilities are thermally unstable, why isn't the equilibrium unstable to axisymmetric modes, calling into question the validity of the equilibrium? Thermal stability analyses of the equilibrium must contend with rather uncertain scalings of the anomalous heat flux. What is needed is a clear-cut signature that radiative cooling is the dominant instability drive and experimental observation of the signature. As mentioned above, the rippling-mode drive did not survive such a test. Furthermore, measurements indicate⁶⁵ ∇Z_{eff} is small, eliminating this drive for resistivity-gradient instabilities. Edge microinstability theories must also explain the dominant k_{gp} value observed in TFTR imaging experiments.⁵³

3. Identification of Confinement Enhancement Mechanisms: What Knobs Can We Turn?

A common theme in our discussion has been the role of the density profile in determining whether confinement is of the supershot, L-mode, or H-mode variety. It is just as important to identify the mechanisms governing particle transport as the mechanism of energy transport, because it is through control of particle transport that one hopes (plans?) to control energy transport. Pellets, NBI fueling, separatrixes, and control of wall recycling provide knobs which the experimentalist can employ to influence the source terms in the density equation, and hence, the profile. By contrast, in an ignited tokamak the heat source will always be α -particle electron heating at the plasma center.

It follows that theory should give particle transport high priority. Theory suggests that flat density profiles should reduce transport arising from electron drift modes

whereas steep density gradients potentially stabilize ion temperature gradient modes and ameliorate collisionless trapped electron modes. These qualitative arguments correspond to two observed enhanced confinement regimes, and underscore the importance of density profile control. An interesting observation is that DIII-D H-modes^{29,67} maintain very flat density profiles while TFTR supershots¹⁷ as well as ASDEX and JFT-2M improved confinement regimes¹⁸ have steep profiles. Hence, experiments suggest a bifurcation in particle transport mechanisms wherein microinstabilities associated with steep gradients cause an inward particle advection which is absent in H-mode. Edge physics triggers the bifurcation.

In many instances, confinement in ohmically-heated tokamaks is enhanced over Goldston-Aachen scaling, based on ohmic power. Within our working hypothesis that ohmic heating does not fundamentally differ from any other centrally-peaked electron heating method (e.g., electron cyclotron heating), the linear neo-Alcator regime constitutes a confinement enhancement technique which ASDEX¹⁸ has shown can be extended by density profile control. Theory interprets the improved confinement as resulting from higher collisionality via dissipative trapped electron mode physics.³³ High power, high density auxiliary heating experiments should search for similar enhancements when $v_e^* \geq 0.1$. Electron cyclotron heating, which has a peaked profile and avoids additional neutral gas pressure, is the technique of choice here.

Inherently, particle transport is more difficult to calculate than heat transport because of the ambipolarity constraint. Test particle methods give a good estimate of heat diffusivity, but lead to well-known errors in particle transport. Simulations of microinstability turbulence should include careful treatment of surface-average ($m=0$, $n=0$) equations to assure they obey conservation laws and include neoclassical effects, such as poloidal rotation damping. The mechanism causing inward advection of density⁶⁶ remains to be identified. It appears that we must rely on simulation to provide insight. It is for this reason that simulations must be free of spurious surface-average effects.

In addition to density profile, we can exercise control over the plasma shape, the existence and geometry of separatrixes, current density profile, wall recycling properties and locations, and, in present experiments, heating power and heat deposition profile. In H-mode plasmas, energy and particles are transferred through the plasma surface by series

of relaxation oscillations^{16,29} called ELMs (Edge Localized Modes) which are thought to be resistive or ideal MHD instabilities triggered by a steep pressure gradient at the separatrix. Heating power has been shown⁶⁷ to regulate the frequency and intensity of ELMs and, through ELMs, the impurity concentrations. Plasma shape is currently determined by ideal MHD stability considerations. Whether such shapes also optimize transport is an open issue. Plasma elongation, which increases the Troyon β -limit,³⁵ also has been found experimentally to improve confinement as evidenced by the $(1+\kappa^2)$ terms in Tables 3 and 4. The physics principles underlying this confinement improvement with elongation remain unclear. As for wall recycling, recent JET data open the possibility that beryllium tiles or gettering may be preferable to the conventional graphite tiles planned for CIT. Beryllium gettering enabled JET to attain H-modes with ICRF alone⁶⁸; the precise physical reason for this improvement remains to be determined.

TFTR's discovery of the supershot, high- T_i enhanced confinement regime¹⁷ with peaked density profiles stimulated investigation of $T_i \gg T_e$ discharges in JET and flat-profile, H-mode discharges with $T_i \gg T_e$ in DIII-D. Since ion confinement did not degrade in DIII-D, even though η_i was large, one can speculate that $T_i \gg T_e$ ameliorates transport arising from the ion temperature gradient mode in qualitative accord with linear theory.²⁸ Can one exploit $T_i \gg T_e$ in CIT? While α -particles preferentially heat electrons, a driven plasma with centrally concentrated ICRF⁶⁹ heating via second harmonic deuterium absorption could potentially raise T_i above T_e . The classical electron-ion energy interchange rate is

$$\frac{dP}{dV} = (500 \text{ kilowatts/m}^3) \left(\frac{n_e}{10^{20} \text{ m}^{-3}} \right)^2 \left(\frac{T_i - T_e}{10 \text{ keV}} \right) \left(\frac{10 \text{ keV}}{T_e} \right)^{3/2}.$$

It would not be unreasonable to concentrate ≈ 40 MW of ICRF power into the central 10% of the plasma volume (e.g., 3m^3 for CIT) leading to the possibility of a hot-ion core in a CIT-class device. Scientifically, a high- T_i core CIT could prototype a high-Q reactor as well as lead to studies of the effect of T_i/T_e on transport.

Sawtooth suppression via energetic particles created by ICRF or α -particles from thermonuclear reactions constitutes a confinement enhancement method of great potential

value to CIT. Usually, the dominant form of energy transport out of the central region of a tokamak, where thermonuclear burning occurs, is through sawteeth relaxation oscillations which greatly increase thermal conductivity (probably by magnetic ergodicity) throughout the region inside $q \sim 1.2$. In the absence of sawteeth, experiments find the thermal conductivity in this region to be quite low, a fact that can be understood theoretically as resulting from the combined effects of a flat density profile and low trapped-particle concentration. Recent experiments on JET⁷⁰ and theoretical work⁷¹ have shown that it is possible to suppress sawteeth with energetic particles, which, if realized with α -particles on CIT, should appreciably reduce the effective thermal conductivity inside $q \approx 1.2$ where most of the thermonuclear reactions will take place, and enhance ignition margin.

4. Discussion: Impact on Ignition Margin.

How does one optimize tokamak designs to obtain maximum ignition margin? As an example, let us start with the microinstability scaling expression for ignition margin $nT\tau_E$ from Table 3

$$nT\tau_E \propto \frac{n a^3 B^3}{T^{0.5}} \propto \frac{\beta B^4 a^3}{T^{1.5}} \quad (23)$$

High density operation leads to high ignition margin, but density is limited by the Troyon β -limit. Customarily, one envisions operating at a β -value roughly 50 - 70% of the ideal MHD Troyon limit.³⁵ Therefore, let us re-express the functional dependence of ignition margin $nT\tau_E$ by eliminating n in favor of β via the definition $n \equiv \beta B^2 / 16\pi T$. Ignition margin increases linearly with β , showing that a high β -limit is an important design criterion. Next, at fixed β , note that the temperature enters with a negative power $T^{-\sigma}$. This, of course, is an artifact of our oversimplified representation of the fusion rate $\langle \sigma v \rangle_e \propto T^2$. A $T^{-\sigma}$ dependence simply means that ignition margin maximizes at a temperature where

$$\frac{T}{\langle \sigma v \rangle_e} \frac{d \langle \sigma v \rangle_e}{dT} = 2 + \sigma \quad (24)$$

Here $\langle \sigma v \rangle_e$ denotes the effective fusion rate, including bremsstrahlung losses. If the exponent σ exceeds unity, then ignition margin maximizes at some temperature between 6

and 10 keV. The final factor $B^4 a^3$ serves as a figure-of-merit for comparing tokamak designs at fixed elongation and $q_{0.95}$ (or I/aB_T). The ITER value exceeds CIT by a factor of 1.6, giving ITER a somewhat larger ignition margin (assuming microinstability scaling). A key aspect is the high powers with which B and a enter the figure-of-merit -- a 10% increase in B increases ignition margin by 40%.

In an effort to reduce the uncertainties in projecting performance of CIT and ITER, the concept of similarity scaling^{19,20,22} has been introduced wherein one seeks to compare a present experiment with anticipated CIT discharges at fixed nondimensional parameters ρ^* , ν^* , β , q , κ , η_i and plasma shape. Here

$$\rho^* = \left(\frac{T}{M}\right)^{1/2} \left(\frac{Mc}{eBa}\right) \quad (25)$$

denotes the normalized gyroradius. Collisionality ν^* and β have their usual definitions

$$\nu^* = (\text{const.}) \left(\frac{n Z_{\text{eff}} a}{T^2}\right) \quad , \quad (26)$$

$$\beta = \frac{16\pi n T}{B^2} \quad (27)$$

The idea behind similarity scaling is that, if basic plasma physics principles govern confinement, then the confinement time must be expressed as^{19,20}

$$\tau_E = \left(\frac{Mc}{eB}\right) \psi(\rho^*, \nu^*, \beta, q, \kappa, \frac{a}{R}, \eta_i, \dots) \quad (28)$$

where the nondimensional arguments of ψ could comprise all entries in Table 5. If one could find a tokamak discharge with all nondimensional parameters identical with those proposed for CIT, then those general basic plasma physics principles predict that confinement should scale as $\tau_E \propto A/B$. One need not be concerned with how collisionality and finite- β effects are influencing microinstability transport, for example. Their relative contribution should remain fixed under similarity scaling. A sequence of discharges with identical nondimensional parameters is parameterized by magnetic field strength according to

$$a = (\text{const.}) \left(\frac{\beta A^3 Z_{\text{eff}}}{\rho^{*6} B^4 v^*} \right)^{1/5} \quad (29)$$

$$n = (\text{const.}) \left(\frac{\beta^3 \rho^{*2} v^{*2} B^8}{A Z_{\text{eff}}^2} \right)^{1/5} \quad (30)$$

$$T = (\text{const.}) \left(\frac{\beta^2 B^2 A Z_{\text{eff}}^2}{\rho^{*2} v^{*2}} \right)^{1/5} \quad (31)$$

Unfortunately, no discharges exist which can be scaled to CIT via Eqs.(29)-(31). One can, however, test the hypothesis that $\tau_E \propto B^{-1}$ by substituting expressions (29)-(31) into unconstrained empirical confinement scalings using

$$I \propto aB \propto B^{1/5} \quad (32)$$

$$P \propto \frac{n T a^3}{\tau_E} \propto \frac{1}{B^{2/5} \tau_E} \quad (33)$$

which transforms empirical scaling laws into a form

$$\tau_E = B^{-x} \Psi (v^*, \rho^*, \beta, q, \kappa, \frac{a}{R}, \eta_i, \dots) \quad (34)$$

If the similarity hypothesis is valid, then $x=1$. Table 6 gives results for several unconstrained empirical laws.

Table 6. Tests of τ_E Scaling with B.

Scaling Law	Exponent α
Goldston Aachen	1.4
Kaye-all-complex	0.3
ITER	0.67
NeoAlcator	0.8
Goldston thermal	0.62

Using a similar analysis, Goldston has found that the rms error of a fit of form (34) to the tokamak data base has a shallow minimum at $\alpha=0.94$. The case for similarity scaling is satisfactory; certainly, it is consistent with the tokamak data base. But, as the examples of Table 6 indicate, many good fits to data violate similarity scaling. One concludes either (1) that physics beyond that of thermal confinement is at work, or (2) that because of unappreciated correlated variables (e.g., η_i and n), important systematic biases are affecting empirical scaling. One can also ask whether the functional form of Ψ in Eq. (34), as determined by regression analysis, is close to a form corresponding to a justifiable theory. In other words, an arbitrary combination of nondimensional parameters resulting from regression analyses does not necessarily correspond to theory.

Since a discharge similar to CIT in all three nondimensional parameters cannot be found, one must adopt a scaling law to project CIT performance. One can argue that fine-scale microinstability scaling is our best present choice for scaling tokamaks in that: (1) It gives a reasonable account of tokamak confinement data (Fig. 1). (2) The theoretical basis is well-established. Linear instabilities are predicted and the corresponding fluctuations observed⁷³ (see also Fig. 4). Nonlinear fluctuation levels are comparable to rough theoretical estimates. Alternative theories such as resistive or ideal MHD either can not be justified for high-temperature tokamaks or fail to predict linear instabilities. (3) Microinstability diffusivities have a reasonable, finite limit as $n \rightarrow 0$ and increase as gradients steepen. Against these arguments stand the concern that the fine-scale ansatz

may not be fulfilled and the theoretical-experimental contradictions set forth in Sec. 2. For example, microinstability scaling fares poorly in relating the two discharges of Table 7. The difficulty regarding empirical scaling projections is that many scaling relations, with differing predictions regarding the optimization of ignition margin, fit the tokamak data base well. Furthermore, systematic corrections to the data change the relations and the diffusivities associated with empirical scalings vanish as $n \rightarrow 0$.

Fine-scale microinstability scaling takes the form

$$\tau_E = \left(\frac{Mc}{eB} \right) (\rho^*)^{-3} \psi(v^*, \beta, q, \kappa, \frac{a}{R}, \eta_i, \dots) \quad (35)$$

One can still arrange that discharges have identical values of v^* and β by scaling discharge parameters according to

$$n \propto B^{4/3} (\beta^2 v^* / Z_{eff})^{1/3} a^{-1/3} \quad (36)$$

$$T \propto (\beta/v^*)^{1/3} (B^2 a Z_{eff})^{1/3} \quad (37)$$

With v^* and β constant, the relative contribution of fine-scale magnetic braiding and collisionality to microinstability transport will remain constant and those uncertainties are eliminated. The validity of confinement scaling (35) thus rests only on whether the fine-scale ansatz applies. Discharge parameters, confinement times, and ignition margins in Table 7 are based on Eqs. (35)-(37). The scaling of Table 7 will be valid if fine-scale microinstabilities govern confinement *and* if one can argue that other nondimensional parameters affecting confinement enhancement will also remain constant. In view of the observed variation of diffusivity between L-mode and H-mode by a factor-of-5.6 in response to nondimensional parameter changes, these additional arguments are essential. Succinctly, Table 7 assumes CIT achieves H-mode confinement.

For fine-scale microinstability scaling [formulas (23) or (35)], ignition margin will maximize at $T=8$ keV and a correspondingly high density to achieve a β -value roughly 60% of the Troyon limit. The merits of the "low" temperature, high-density approach to ignition have been recognized by confinement theorists.^{33, 72} Let us use Eq. (23) and the properties of high $nT\tau_E$ discharges in DIII-D and JET, summarized in Table 7, to

predict performance for CIT for the β and v^* values actually attained in these experiments. For DIII-D, one must increase T by a factor-of-2.8 to $T_i(0) = 12.7$ keV (obtaining an average temperature near 8 KeV), B by a factor-of-4.8 to $B=10T$, and A by 1.25. The result is

$$(nT\tau_E)_{CIT} = 100 \cdot 10^{20} \text{ keV-m}^{-3}\text{-sec} \quad (38)$$

yielding an ignition margin of almost 2. Scaling JET, one finds

$$(nT\tau_E)_{CIT} = 79 \cdot 10^{20} \text{ keV-m}^{-3}\text{-sec} \quad (39)$$

Again, the ignition margin exceeds unity. Goldston Aachen scaling, by contrast, is decidedly more pessimistic. Using the $nT\tau_E$ scaling of Table 3, one finds

$$(nT\tau_E)_{CIT} = \begin{array}{ll} 28 \cdot 10^{20} \text{ keV-m}^{-3}\text{sec} & \text{(based on DIII-D)} \\ 43 \cdot 10^{20} \text{ keV-m}^{-3}\text{-sec} & \text{(based on JET)} \end{array} \quad (40)$$

which suffices for $Q>5$. The reason for pessimism is the adverse density scaling of diffusivity, which traces back to the lack of positive density dependence in the confinement time scaling law. In carrying out these scalings, we have assumed that all other nondimensional parameters, including confinement enhancement associated with H-mode, remain fixed. Goldston scaling has the unexpected property that ignition margin is independent of β . Thus, according to Goldston scaling, the sole reason ignition margin improves in an elongated tokamak is the increased confinement which results from the higher current elongation permits. Electrostatic microinstability scaling, on the other hand, predicts ignition margin increases with β , and a high ideal MHD Troyon limit is doubly beneficial.

Since one can find DIII-D and JET discharges which are similar to CIT in v^* and β , the crucial question is: Does confinement in tokamaks scale with ρ^* according to formula (35)? This question can be addressed experimentally by creating a series of discharges with varying magnetic field and with density and temperature controlled via pellets, gas puffs, and auxiliary heating to obey relations (36) and (37). Using Eqs. (25),(36),(37) one finds that ρ^* will vary as

$$\rho^* \propto \left(\frac{\beta Z_{eff}}{v^*} \right)^{1/6} \frac{A^{1/2}}{B^{2/3} a^{5/6}} \quad (41)$$

Let us generalize confinement scaling relation (35) to

$$\tau_E = \left(\frac{Mc}{eB} \right) (\rho^*)^{-x} \Psi(v^*, \beta, q, \kappa, \dots) \quad (42)$$

where $x=3$ for fine-scale microinstabilities and $x=2$ for a (possible) Bohm-like scaling that does not fulfill the fine-scale ansatz. At fixed v^* and β , confinement in a given machine will scale according to

$$\tau_E \propto B^{(2x/3)-1} \quad (43)$$

Hence, the exponent x can be determined experimentally. Such experiments can be carried out even in ohmically-heated plasmas because the required temperature scaling (37) fortuitously coincides with that predicted by a dissipative trapped electron model³³ which reproduces the linear neoAlcator regime. These experiments will complement fluctuation imaging experiments in determining whether the fine-scale ansatz holds.

One notes that the DIII-D and JET discharges of Table 7 differ appreciably in v^* and β . Can fine-scale turbulence account for their measured differences? From Table 3, one can recast the formula for fine-scale turbulence power losses into the form

$$P = (\text{const.}) \beta T^{1.5} A^{0.5} \quad (44)$$

which implies that the JET discharge should have higher, not lower, power loss. This serves as an example of the difficulty encountered concerning the density dependence of τ_E and, by implication, the v^* and β dependence of thermal diffusivity. Another candidate is that profiles differ and JET is further into enhanced confinement.

Table 7. High $nT\tau_E$ Discharges;
CIT Discharges with Identical v^* and β

Quantity	DIII-D		JET	
	Observed	CIT	Observed	CIT
Shot #	65747		20222	
B (Tesla)	2.1	10	2.8	10
R (m)	1.7	2.1	3.1	2.1
a (m)	0.62	0.65	1.0	0.65
I (MA)	2.0	7.7	4.2	10.0
$n_D(0)$ (10^{20} m^{-3})	1.1	8.8	0.56	3.5
τ_E (sec)	0.21	0.89*	1.25	1.35*
$T_i(0)$ (keV)	4.5	12.7	8.3	16.8
$T_e(0)$ (keV)	3.5	9.9	7.6	15.2
$n_D(0)T_i(0)\tau_E$ ($10^{20} \text{ keV m}^{-3} \text{ s}$)	1.0	100*	5.8	80*
β	0.041	0.041	0.024	0.024
κ	2.1	2.1	1.9	1.9
Power [†] (MW)	12	64*	7.2	24*
$\rho^*_{\text{CIT}}/\rho^*_{\text{Device}}$		0.40		0.69

* Scaled according to microinstability scaling $\tau_E \propto B a^{5/2} A^{-1/2}$ with fixed v^* and β .

[†] Required heating power (either thermonuclear or auxiliary); dW/dt subtracted for JET.

Although our understanding of the H-mode transition power threshold P_H is rudimentary, one can use Eqs. (36),(37) in expression (22) for characteristic sheath power and find $P_H \propto B^{5/3} a^{4/3}$, i.e., Bohm scaling. Scaled from JET, the transition power is $P_H \approx 35 \text{ MW}$ and from DIII-D, $P_H \approx 100 \text{ MW}$, based on the highest transition power of $\approx 8 \text{ MW}$ observed in DIII-D. CIT should have high auxiliary power (at least

50MW) to investigate H-mode transitions with nonreacting plasmas before DT burning experiments commence. At present, one must conclude there is uncertainty regarding H-modes in CIT.

Fine-scale turbulence projections indicate that CIT will ignite, provided H-mode is attained. Even if confinement scaling proves to be Bohm-like, CIT will still achieve $Q > 5$, as can be verified by multiplying the projected $nT\tau_E$ values by ρ^*_{CIT} / ρ^* and noting the resulting product still exceeds $30 \cdot 10^{20} \text{ keV m}^{-3} \text{ sec}$. All other nondimensional parameters remain fixed, so confinement scaling with v^* , β , q , etc. does not enter the projection. While prospects that CIT will enter the self-heating regime appear excellent, much remains to be understood regarding transport physics, especially v^* and β effects.

For ITER, which must ignite to fulfill its technology objectives, it is essential that scaling follow from the fine-scale ansatz. If fine-scale turbulence scaling prevails, then ITER's ignition margin exceeds CIT by a factor of 1.8 at constant v^* and β . But since $\rho^*_{ITER} = (0.25)\rho^*_{CIT}$, Bohm-like scaling would call ignition into question. We note that scaling n, T from DIII-D to ITER via Eqs. (36),(37) produces a reasonable, but somewhat low temperature, discharge which favors ignition prospects according to Eq. (23).

The preceding discussion has tried to convey how uncertainties in confinement scaling impact estimates of ignition margin. Both empirical and theoretical confinement scalings face difficulties because the physics mechanisms and scaling governing confinement and confinement enhancement are not yet definitively identified. Fortunately, extensive research has given both phenomenological and theoretical understandings of transport that have much in common, qualitatively and quantitatively. A strong case can be made that CIT will exceed $Q=5$. But the needs of tokamak design require even more precision and definitiveness in predictive capability. As we move forward to reactor-class devices such as ITER, size-scaling becomes crucial. To illustrate the extremes of projections, the empirical Kaye-all-complex scaling relation has a figure of merit of $B^{2.7}a$ versus B^4a^3 for fine-scale turbulence [see Eq. (23)]. Consequently, the ignition margin of ITER is a factor-of-0.50 that of CIT, according to Kaye-all-complex. For fine-scale turbulence, we found a factor of 1.6. There remains much to be learned about transport mechanisms. Transport physics will remain important even for reactors because a thermally-stable, steady-state operating point necessarily involves a balance between

thermonuclear power generation and transport losses. Our "conclusions" then take the form of 21 questions, listed in Table 8, whose answers are needed to improve our understanding of transport physics and sharpen predictive capability.

Acknowledgments: Numerous persons have contributed to my understanding of these issues. In particular, I would like to thank D. Baldwin, G. Bateman, R. Bickerton, H. Biglari, N. Bretz, J. Callen, L. Chen, G. Cordey, S. Cowley, P. Diamond, P. Efthimion, R. Goldston, T. S. Hahm, W. Horton, S. Kaye, K. Lackner, G. H. Neilson, R. Parker, S. Prager, M. Rosenbluth, P. Rutherford, G. Schmidt, J. Sheffield, R. Stambaugh, J. Strachan, W. Tang, R. Waltz, A. Wooton, M. Zarnstorff, and S. Zweben. This work was supported by U.S. Department of Energy Contract No. DE-AC02-76CHO3073.

Table 8. Tokamak /Stellarator Transport Issues

Issue	Experimental Program	Theoretical Program	Resources Required
1. At what value of $k_{\perp} \rho$ does the density fluctuation spectrum peak? Does this value depend on machine size? Is the fine-scale ansatz satisfied? What is the size scaling of confinement?	Imaging experiments of density, temperature fluctuations using optical and X-ray emissions, Thomson scatter, and phase contrast techniques. Also multidetector heavy ion beam probe. Reflectometry with multiple receivers. B field scans to vary ρ^* while holding machine size, v^* , and β fixed. Size-scaling comparisons between JET, DIII-D, and Alcator C-Mod. Also TFTR, PLT, and Alcator C.	3-dimensional fluid and 5-dimensional phase space simulations. Saturation study of trapped ion modes. Wave propagation through density fluctuations to support reflectometry. Computer visualization of density fluctuations to compare with imaging experiments.	Imaging optical and X-ray framing cameras. IR detector arrays. Detector arrays for heavy ion beam probe. Phased array reflectometry receiver. Large supercomputer budget, workstations, high-speed data link. Algorithm specialists to speed up computation.
2. Is the thermal energy confinement time independent of density? Do correlated variables bias empirical scalings? What scaling prevails in the low-density limit? Can auxiliary laws describe ohmic plasmas?	Accurate, high-resolution n, T profiles. Sophisticated energy budget codes which distinguish thermal energy from energetic ion tails. Experimental scans at constant NBI penetration, or with ICRF, ECH heating. Check for η_i correlations with n_e . Low density confinement scans.	Use transport codes to compute synergistic effect of microinstability transport and other physics, e.g., electron-ion equilibration, ionization, NBI penetration etc.	Tokamak operating time with carefully controlled scans. Improved Thomson scatter, $T_i(r)$, $n_i(r)$, etc. ICRF and ECH high power heating. Add ohmic plasmas to data base.

Issue	Experimental Program	Theoretical Program	Resources Required
3. What is physics basis of isotope effect?	Accurate, high-resolution n, T profiles for H^+, D^+ discharges under ohmic, L-mode, enhanced confinement conditions. Accurate measurements of plasma fueling, H_{α} , radiative losses, neutral gas pressures. Measure of fluctuation spectrum via microwave/laser scatter, imaging, and HIBP techniques. Low n_e scans to separate electron and ions.	Transport code investigations of competition between microinstability transport, electron-ion heat exchange, atomic physics sources and sinks and their synergistic effect on profiles and turbulence.	Tokamak operating time. Neutral pressure instrumentation. Pellet injector fueling. Microwave scatter and imaging density fluctuation measurements.
4. Why does confinement scale with I_p , not BT ?	High resolution n, T profiles to deduce effect of q, s on diffusivities. Fluctuation measurements to measure dependence of fluctuation levels on q, s minor radius. Do error fields and magnetic islands play a role? I_p ramps to vary shear.	Analytic and computational studies of the effect of magnetic quantities such as q, s on turbulence. Development of codes for trapped ion modes.	Toroidal, annular simulations. Super-computer time, workstations, and algorithm specialists.
5. Is the tokamak in a marginal-stability state?	Observe response of $T_i(r)$ to $n_i(r)$ changes. Search for rapid response of fluctuations to deviations from marginal stability. Search for relation between T_i edge and $T_i(0)$. Vary power deposition profile.	Improve precision of η_i threshold in toroidal geometry. Revisit heat flux associated with non-linear η_i -modes.	Pellet injector, gas-puff, and ELM density perturbation technique.

Issue	Experimental Program	Theoretical Program	Resources Required
6. Does magnetic braiding transport play a role?	Create impulsive, local bursts of high-energy circulating electrons and follow their diffusion. Measure ratio $q_e/T_e\Gamma$. X-ray pulse height analysis of non-Maxwellian tails. Positron injection techniques. Runaway studies; detection of high energy electrons in plasma periphery. classical MHD simulations.	Develop self-consistent electromagnetic theory and simulation capability Separate rippled magnetic surfaces from magnetic braiding. Develop neo-	ECH and lower-hybrid impulsive heating capability. Synchrotron and X-ray diagnostics.
7. What instability shuts off at the H-mode transition? Is the H-mode transition characterized by $T_{edge} \propto B_T$? Are sheaths an essential element of the H-mode transition?	Visualization by imaging interferometry, imaging reflectometry, and beam emission spectroscopy. Poloidal magnetic probe array observations of turbulence near separatrix. Sensitivity of transition to $j_{ }(r)$, neutral gas pressure. Plasma flow (E_r) measurements in separatrix plasmas. Scaling of transition power threshold. High resolution n , T , ϕ profiles near separatrix to locate sheath regions of poloidal gyroradius scale.	MHD and kinetic equilibrium and stability in the vicinity of separatrix. Neoclassical plasma sheaths. Self-consistent radial electric fields, poloidal flows, and instabilities. Ionization and radiative cooling effects. Linear stability of improved separatrix equilibria.	Fluctuation diagnostics focussed on separatrix vicinity. Extensive neutral pressure measurements. T_i, u_i spectroscopic measurements in separatrix plasma. Improved equilibrium codes with $J_{ } \neq 0$ at separatrix.

Issue	Experimental Program	Theoretical Program	Resources Required
8. What is level of density fluctuations when $\nabla n_e = 0$?	Microwave laser scatter and imaging fluctuation measurements in flat density plasmas. This isolates the negative compressibility ITG drive. Measure $\tilde{T}_i$.	Fluid and phase space simulations of fluctuations with $\nabla n_e = 0$. Investigate role of non-zero ∇n_i and ∇n_e .	Same as 1,4. Fast ion temperature measurement techniques.
9. How do fluctuations and transport respond to non-dimensional parameters, especially steep density, supershot profiles? Can a linear neoAlcator regime be found with ECH?	Generalization of 4 and 8. Of key interest here is the parameter $e_n = L_n/R$ which governs trapped particle - drift wave interaction, as well as q which observations show cannot be small in supershots. High power, 2 nd harmonic ECH heating of dense plasmas. Pellet injections.	Focus on marginal stability and saturation of instabilities at small e_n and n_i . Investigate density transport mechanism, especially density advection physics.	Pellet injectors. ICRF and ECH heating. Simulation codes with accurate, neoclassical surface-average equations.
10. How does the vacuum wall material affect the plasma edge?	Measurement of particle, photon, and rf wave fluxes incident on vacuum wall. Measurements of n_e , T_e , T_i at wall and limiters. Neutral gas measurements in divertor, main chamber. Detached plasmas as a perturbation technique.	Develop model for response of wall material to fluxes and rf waves. Compute effect of ionization and radiative loss on edge plasma equilibrium.	Neutral gas and low-emergy neutral particle diagnostics for particles incident on plasma. JET and ASDEX data crucial.

Issue	Experimental Program	Theoretical Program	Resources Required
11. Can high - T_i plasmas be exploited in CIT?	2 nd harmonic ICRF heating of deuterium, pellet-produced plasmas in TFTR and JET. Study relative ICRF heat deposition to electrons and ions.	Computation of electron versus ion power deposition profiles for fast waves near the 2 nd harmonic deuterium resonance.	High power ICRF at 80 MHz for TFTR and 55 MHz for JET.
12. How do shear, magnetic well, and poloidal field affect fluctuation amplitudes?	Exploit stellarator geometry, especially the ATF torsatron, where the rotational transform is largest at the plasma edge, in contrast to a tokamak. This should help clarify the role poloidal field on transport coefficients.	Develop theory of micro-instabilities in torsatron geometry.	Scattering, visualization and heavy ion beam probe measurements in ATF.
13. Are there regimes of different confinement physics and scaling? What marks the boundary between linear neo-Alcator and saturated ohmic confinement?	Experimental scans which emphasize large ranges of nondimensional parameters. Does collisionality characterize the transition between linear and saturated ohmic confinement?	Linear theory focussed on identifying regimes where various microinstabilities are expected to dominate. Nonlinear theory at identified regime boundaries.	Experiments with a full complement of active controls: NBI, ICRF, ECH, pellets, wall conditioning, etc.

Issue	Experimental Program	Theoretical Program	Resources Required
14. Are locked modes and neo-classical MHD modes contributing to transport?	More accurate Thomson scatter for T_e plateaus associated with low-order rational surfaces. Large-scale imaging diagnostics. Poloidal, toroidal arrays of peripheral magnetic probes. Fast interferometry of large-scale density fluctuations. Measure error fields. See also 6.	Develop non-axisymmetric equilibrium codes and neoclassical MHD time-dependant codes.	Supercomputer time. High laser pulse energies to improve Thomson scatter accuracy.
15. Does magnetic flutter at the c/ω_{pe} scale cause transport?	Microwave/laser scattering measurements in the $k\rho_i \gg 1$ range.	Develop nonlinear theory of η_e modes, especially in the presence of magnetic flutter.	Large-angle microwave/laser scattering geometries.
16. Is the H-mode non-local?	Measure time response of core diffusivities and fluctuations to H-mode transition. Accurate, time-dependent n , T profiles and energy budget analysis.	If yes, develop marginal stability model of core H-mode transport.	Same as 1,2.

Issue	Experimental Program	Theoretical Program	Resources Required
17. How does plasma shape, local magnetic shear, and separatrix geometry influence transport, edge E_r , and H-mode transitions?	Measure plasma transport at fixed B_T , a , q , n versus shape. Spectroscopic measurements of poloidal flow, E_r near separatrix.	Understand how plasma shape influences linear stability, nonlinear saturation of dominant instabilities. Development of neoclassical theory of sheath in separatrix region.	Tokamak with flexible shape, auxiliary heating, etc.
18. Are trapped electrons playing a crucial role?	Contrast inside and outside ECE fluctuation measurements. Large optical depth means that inside ECE is sensitive to high energy, trapped electrons. Imaging of superthermal X-ray fluctuations. Experiments in stellarator geometry, where trapped particle fraction is controllable.	Develop ECE emission theory for fluctuating non-Maxwellian velocity distribution.	Inside ECE observational capability. X-ray tangential imaging capability with response for superthermal X-rays.

Issue	Experimental Program	Theoretical Program	Resources Required
19. What triggers ELMs?	Dependence on $j_{ }$, plasma pressure profiles in separatrix region. Accurate pressure profiles needed to assess ideal MHD ballooning stability.	Ideal and resistive MHD simulations of plasma separatrix region.	High resolution Thomson scatter measurements of plasma pressure profiles in separatrix region.
20. Will energetic particle control of sawtooth enhance ignition margin in CFT?	ICRF creation of energetic particles to greatly increase sawtooth period. Exploit excellent confinement of $q \approx 1$ region.	Develop energetic particle algorithm for use in resistive MHD codes.	High-power fast wave ICRF heating in circular and shaped tokamaks. Super-computer time.
21. Can local micro-instability turbulence models account for heat pulse propagation results?	Attempt heat pulse propagation in situations free of complicating effects of the sawtooth crash, e.g., pulsed ECH studies. Accurate, high-time-resolution T_e and n_e profiles. Compare ECH, sawtooth-induced heat pulses.	Model heat pulse propagation using microinstability diffusivities, accounting for $\bar{T}_e$ effects on diffusivity.	Pulsed ECH heating.

References

- ¹R. Parker *et al.*, Plasma Phys. Controlled Fusion Res. 1988, Vol. 3 (Proc. 12th Intl. Conf., Nice, IAEA, Vienna 1989) p341.
2. M. Crawford, Science **246** 577 (1989).
3. E. Marshall, Science **246** 1110 (1989).
- ⁴ J. F. Lyon *et al.*, Fusion Technol. **10** 179 (1986); J. F. Lyon, Fusion Technol. **17** 19 (1990).
- ⁵ H. Zushi *et al.*, Nucl. Fusion **28** 1801 (1988); T. Obiki *et al.*, Plasma Phys. Controlled Fusion Res. 1988, Vol. 2 (Proc. 12th Intl. Conf., Nice, IAEA, Vienna 1989) p 337.
- ⁶ K. Matsuoka *et al.*, Plasma Phys. Controlled Fusion Res. 1988, Vol. 2 (Proc. 12th Intl. Conf., Nice, IAEA, Vienna 1989) p 411.
- ⁷J. Sapper and H. Renner, Fusion Technol. **17** 62 (1990); H. Renner *et al.*, "Status of the Advanced Stellarator Wendelstein VII-AS; First Results and Future Program," 7th Intl. Workshop Stellarators, Oak Ridge (April 1989)
- ⁸ C. Beidler *et al.*, "Physics and Engineering Design for Wendelstein VII-X," 7th Intl. Workshop Stellarators, Oak Ridge (April 1989)
- ⁹ A. Iiyoshi *et al.* , Fusion Technol. **17** 169 (1990); O. Motojima, "Recent Design Consideration for the Large Helical Device," 7th Intl. Workshop Stellarators, Oak Ridge (April 1989)
- ¹⁰L. M. Kovrizhnyk, Nucl. Fusion **24**, 851 (1984); H. E. Mynik, Phys. Fluids **25**, 325 (1982); K. Itoh, S.-I. Itoh *et al.*, Nucl. Fusion **29**, 1851 (1989); L. M. Kovrizhnyk and S. G. Shasharina, Nucl. Fusion **30** 851 (1990).

- ¹¹ K. Tombechi, Plasma Phys. Controlled Fusion Res. 1988, Vol. 3 (Proc. 12th Intl. Conf., Nice, IAEA, Vienna 1989) p215.
- ¹² R. J. Hawryluk *et al.*, Plasma Phys. Controlled Fusion Res. 1986, Vol. 1 (Proc. 12th Intl. Conf., Kyoto, IAEA, Vienna 1987) p 51.
- ¹³ A. Gibson, JET Team, "A Review of the JET Project," Intl. Plasma Physics Conf., New Delhi (1989).
- ¹⁴ R. J. Goldston, Plasma Phys. Controlled Fusion **26** 87 (1984).
- ¹⁵ ITER Concept Definition, Vol. 2, p8 (IAEA, Vienna 1989); S. M. Kaye *et al.*, "Status of Global Energy Confinement Studies," Princeton Plasma Physics Laboratory Report PPPL-2670 (1990).
- ¹⁶ ASDEX Team, Nucl. Fusion **29** 1959 (1989).
- ¹⁷ M. G. Bell *et al.*, Plasma Phys. Controlled Fusion Res. 1988, Vol. 1 (Proc. 12th Intl. Conf., Nice, IAEA, Vienna, 1989) p27.
- ¹⁸ G. Fussmann *et al.*, Plasma Phys. Controlled Fusion Res. 1988, Vol. 1 (Proc. 12th Intl. Conf., Nice, IAEA, Vienna, 1989) p145. F. X. Söldner *et al.*, Phys. Rev. Lett. **61** 1105 (1988); M. Mori *et al.*, Nucl. Fusion **28** 1892 (1988).
- ¹⁹ B. B. Kadomstev, Sov. J. Plasma Phys. **1** 295 (1975).
- ²⁰ J. W. Connor, Plasma Phys. Controlled Fusion **30** 619 (1988); K. Lackner and N. Gottardi, Nucl. Fusion **30** 767 (1990).
- ²¹ W. Horton *et al.*, Plasma Phys. **22** 663 (1980).
- ²² J. W. Connor, Nucl. Fusion **26** 193 (1986).

- 23 J. A. Snipes *et al.*, Nucl. Fusion **28** 1085 (1988); H. Ninomiya *et al.*, Nucl. Fusion **28** 1275 (1988); J. A. Snipes *et al.*, Nucl. Fusion **30** 205 (1990).
- 24 J. D. Callen and K. C. Shaing, Phys. Fluids **28** 1845 (1985); K.C. Shaing and J.D. Callen, Phys. Fluids **28** 1859 (1985); R. Fitzpatrick, Phys. Fluids **B1** 238 (1989).
- 25 M. Kotschenreuther, Phys. Fluids **28** 294 (1985).
- 26 J. L. Johnson and A. Reiman, Nucl. Fusion **28** 1116 (1988); A. Reiman and H. Greenside, Computer Phys. Commu. **43** 157 (1986).
- 27 P. W. Terry *et al.*, Phys. Fluids **31** 2920 (1988); G.S. Lee and P. H. Diamond, Phys. Fluids **29** 3291 (1986); N. Mattor and P. H. Diamond, Phys. Fluids **31** 1180 (1988); N. Mattor and P.H. Diamond, Phys. Fluids, **B1** 1980 (1989); H. Biglari *et al.*, Phys. Fluids **B1** 109 (1989); R. D. Sydora *et al.*, Phys. Rev. Lett. **64** 2015 (1990).
- 28 T. S. Hahn and W. M. Tang, Phys. Fluids **B1** 1185 (1989). R. R. Dominguez and R. E. Waltz, Phys. Fluids **31** 3147 (1988).
- 29 D. P. Schissel *et al.*, Nucl. Fusion **29** 185 (1989).
- 30 W. M. Tang, Nucl. Fusion **18** 1089 (1978); P. C. Liewer, Nucl. Fusion **25** 543 (1985); F. Romanelli, Plasma Phys. Controlled Fusion **31** 1535 (1989).
- 31 A. Bhanucharjee *et al.*, Phys. Fluids **26** 880 (1983).
- 32 W. M. Tang *et al.*, Phys. Fluids **29** 3715(1986); G. Rewoldt *et al.*, Phys. Fluids **30** 807 (1987); G. Rewoldt, Phys. Fluids **31** 3727 (1988).
- 33 F. W. Perkins and Y.C. Sun in *Basic Physical Processes of Toroidal Fusion Plasmas*, Vol. 2, p587 (Varenna, 1985).
- 34 N. Mattor and P.H. Diamond, Phys. Fluids **B1** 1993 (1989).

- 35 F. Troyon *et al.*, Plasma Phys. Controlled Fusion **26** 209 (1984).
- 36 S. Sudo *et al.*, Nucl. Fusion **30** 11 (1990); F. Sano *et al.*, Nucl. Fusion **30** 81 (1990).
- 37 K. Ida *et al.*, Plasma Phys. Controlled Fusion **31** 2127 (1989).
- 38 S. Fairfax *et al.*, Plasma Phys. Controlled Fusion Res. 1980, Vol. 1 (Proc. 8th Intl. Conf., Brussels, IAEA, Vienna, 1981) p.439; J. G. Cordey *et al.*, Plasma Phys. Controlled Fusion Res. 1984, Vol. 1 (Proc. 10th Intl. Conf., London, IAEA, Vienna, 1985) p.167.
- 39 S. D. Scott *et al.*, Phys. Rev. Lett. **64** 531 (1990); M. H. Redi *et al.*, Nucl. Fusion **27** 2001 (1987).
- 40 R. R. Dominguez and R.E. Waltz, Nucl. Fusion **27** 65 (1987).
- 41 S. C. Guo, L. Chen, S. T. Tsai, and P.N. Guzdar, Plasma Phys. Controlled Fusion **31** 423 (1989).
- 42 O. J. Kwon *et al.*, Phys. Fluids **B2** 291 (1990).
- 43 M. Soler and J. D. Callen, Nucl. Fusion **19** 703 (1979); G. L. Jahns *et al.*, Nucl. Fusion **26** 226 (1986); E. D. Fredrickson *et al.*, Nucl. Fusion **26** 849 (1986); B. J. D. Tubbing *et al.*, Nucl. Fusion **27** 1843 (1987); N. J. Lopes *et al.*, Nucl. Fusion **28** 1173 (1988); K. S. Riedel *et al.*, Nucl. Fusion **28** 1509 (1988).
- 44 A. J. Wootton *et al.*, Plasma Phys. Controlled Fusion Res. 1988, Vol. 1 (Proc. 12th Intl. Conf., Nice, IAEA, Vienna, 1989) p293. Ch. P. Ritz *et al.*, Phys. Rev. Lett. **62** 1844 (1989).
- 45 A. B. Rechester and M. N. Rosenbluth, Phys. Rev. Lett. **40** 38 (1978); J.A. Krommes, C. Oberman, and R. Kleva, J. Plasma Phys. **30**, part1, 11 (1983).
- 46 W. Horton *et al.*, Phys. Fluids **26** 1461 (1983); P. W. Terry *et al.*, Phys. Rev. Lett.

- 57 1899 (1986), Phys. Rev. Lett. **60** 966 (1988); A. A. Thoul *et al.*, Phys. Rev. Lett. **59** 1488 (1987), Phys. Rev. Lett. **60** 967 (1988);
- 47 P. H. Rebut *et al.*, Plasma Phys. Controlled Fusion Res. 1986 Vol. 2 (Proc. 11th Intl. Conf., Kyoto, IAEA, 1987) p187.
- 48 R. B. White and F. Romanelli, Phys. Fluids **B1** 977 (1989).
- 49 R. Carrera *et al.*, Phys. Fluids **29** 899 (1986).
- 50 P. H. Rebut *et al.*, Plasma Physics Controlled Fusion Res. 1988, Vol. 2 (Proc. 12th Intl. Conf., Nice, IAEA, Vienna 1989) p 191.
- 51 P. N. Guzdar *et al.*, Phys. Rev. Lett. **57** 2818 (1986); Y. C. Lee *et al.*, Phys. Fluids **30** 1331 (1987).
- 52 B. Grek *et al.*, Bull. Am. Phys. Soc. **34** 1965 (1989).
- 53 S. J. Zweben and S. S. Medley, Phys. Fluids **B1** 2058 (1989).
- 54 H. Matsumoto *et al.*, Nucl. Fusion **27** 1181 (1987); K. Hoshino *et al.*, Nucl. Fusion **28** 301 (1988); T. Ogawa *et al.*, Nucl. Fusion **30** 499 (1990).
- 55 C. E. Bush *et al.*, "Peaked Density Profile Circular Limiter H-Mode Plasmas on TFTR", Princeton Plasma Physics Laboratory Report PPPL-2693 (1990).
- 56 P. Gohil *et al.*, Phys. Rev. Lett. **61** 1603 (1988).
- 57 R. J. Groebner *et al.*, Proc. 16th European Conf. Control. Fusion Plasma Phys. (EPS, Petit-Lancy, Switzerland, 1989) Vol.1, p 245. T. Carlstrom *et al.*, Proc. 16th European Conf. Control. Fusion Plasma Phys. (EPS, Petit-Lancy, Switzerland, 1989) Vol 1. p 241.
- 58 K. C. Shaing and E. C. Krume, Phys. Rev. Lett. **63** 2369 (1989); H. Biglari *et al.*, Phys. Fluids **B2** 1 (1990).

- 59 S. Briguglio *et al.*, Phys. Fluids **B1** 1449 (1989).
- 60 F. Hinton, Nucl. Fusion **25** 1451 (1985); R. D. Hazeltine, Phys. Fluids **B1** 2031 (1989); F. L. Hinton and G. M. Staebler, Nucl. Fusion **29** 405 (1989); S. J. Kilpatrick *et al.*, "The Effect of Plasma Density on the Behavior of the Plasma Boundary in Helium Discharges in the Tokamak Fusion Test Reactor," Princeton Plasma Physics Laboratory Report PPPL-2683 (1990).
- 61 M. Keilhacker and the JET Team, Phys. Fluids **B2** 1291 (1990); J. Jaquinot *et al.*, "Effect of Beryllium Evaporation on the Performance of ICRH on JET," JET Report JET-P (89) 85.
- 62 C. P. Ritz *et al.*, Nucl. Fusion **27** 1125 (1987).
- 63 D. R. Thayer and P.H. Diamond, Phys. Fluids **30** 3724 (1987). J. F. Drake, L. Sparks, and G. Van Hoven, Phys. Fluids **31** 813, (1988);
- 64 K. Gentle and Y. J. Kim (Private Communication); Y. J. Kim, "General Description of Magnetic Fluctuations in TEXT," Ph. D. Thesis, University of Texas (1989).
- 65 D. Schissel *et al.*, Phys. Fluids **31** 3738 (1988).
- 66 J. D. Strachan *et al.*, Nucl. Fusion **22** 1145 (1982).
- 67 K. Burrell *et al.*, Plasma Phys. Controlled Fusion **31** 1649 (1989); G. L. Jahns *et al.*, Nucl. Fusion **29** 1271 (1989).
- 68 B. J. D. Tubbing *et al.*, Nucl. Fusion **29** 1953 (1989).
- 69 F. W. Perkins, 3rd Intl. Meeting on Theoretical and Experimental Aspects of Heating Toroidal Plasmas, Vol. 2 (Grenoble, 1976) p333.

⁷⁰ D. J. Campbell *et al.*, Phys. Rev. Lett. **60** 2148 (1988); J. Jacquinot *et al.*, Plasma Physics Controlled Fusion Res. 1986, Vol. 1 (Proc. 11th Intl. Conf., Kyoto, IAEA, Vienna 1987) p 449.

⁷¹ R. B. White, M.N. Bussac, and F. Romanelli, Phys. Rev. Lett. **62** 539 (1989); J. P. Christiansen, J. G. Cordey, D. G. Muir, Nucl. Fusion **29** 1505 (1989).

⁷² R. E. Waltz *et al.*, Nucl. Fusion **29** 351 (1989).

⁷³ P. C. Liewer, Nucl. Fusion **25** 543 (1985); T. Crowley and E. Mazzucato, Nucl. Fusion **25** 507 (1985); C. M. Surko and R. E. Slusher, Science **221** 817 (1983); D. L. Brower *et al.*, Nucl. Fusion **29** 1247 (1989).