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ENERGY CONFINEMENT IN TOKAMAKS

By

M. Sugihara and C. Singer

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PRINCETON UNIVERSITY
PRINCETON, NEW JERSEY

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M. Sugihara[†] and C. Singer

Plasma Physics Laboratory, Princeton University

Princeton, NJ 08544

ABSTRACT

In this letter, we make a straightforward generalization of the ohmic heating energy confinement scalings of Pfeiffer and Waltz and Blackwell et al. The resulting model is systematically calibrated to published data from limiter tokamaks with ohmic, electron cyclotron, and neutral beam heating. With considerably fewer explicitly adjustable free parameters, this model appears to give a better fit to the available data for limiter discharges than the combined ohmic/auxiliary heating model of Goldston.

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[†]Present address: JAERI, Naka, Ibaraki, Japan

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In this letter, we make a straightforward generalization of the ohmic heating energy confinement scalings of Pfeiffer and Waltz [1] and Blackwell et al. [2]. The resulting model is systematically calibrated to published data from limiter tokamaks with ohmic, electron cyclotron, and neutral beam heating. With considerably fewer explicitly adjustable free parameters, this model appears to give a better fit to the available data for limiter discharges than the combined ohmic/auxiliary heating model of Goldston [3].

We start with a dimensionally constrained [4,5] version of the Pfeiffer-Waltz type of scaling for low density, ohmically heated plasmas

$$\tau_{SN} = \alpha n a K^c R^{9/4} (T_{OH}/T)^c, \quad (1)$$

where

$$T_{OH} = \left\{ \frac{(0.5 \alpha n a K^c R^{9/4})}{3 \langle n \rangle k_B (2\pi^2 R a^2 K)} \frac{(2\pi R I^2 m_e^{1/2} Z_{eff}^2 \lambda)}{[6 \epsilon_0^2 (2\pi k_B)^{3/2}]} \right\}^{2/5}, \quad (2)$$

and τ_{SN} , n and $\langle n \rangle$, a , K , R , T , I , and Z_{eff} are an estimate of the global energy confinement time, the line and volume average electron density, midplane half-width, half-height/half-width, major radius, density-weighted volume average thermal plasma temperature, total toroidal plasma current, and volume average effective charge, respectively

(in SI units with temperatures in keV). T_{OH} is by definition the equilibrium temperature obtained with ohmic heating assuming Spitzer's resistivity [6] in a uniform plasma column with energy confinement time $\tau_E = \tau_{SN}/2$ [cf. comments following Eq. (11-13) below]. m_e , e , k_B , and ϵ_0 are the electron mass, electron charge, Boltzmann constant, and electrical permittivity of free space. In this paper, we approximate the Coulomb logarithm $\lambda = 37.8 - \ln(n^{1/2}/T) \approx 16$ and volume average density $\langle n \rangle = 3n/4$ (appropriate to parabolic density profiles). α , c , and t are adjustable constants to be determined from the confinement data. The response of conduction-dominated plasmas to different heating profiles is obtained by assuming a heat conduction coefficient $\kappa = \kappa_a \rho / \rho_a$ (cf. Singer et al., [7]), where $\rho < \rho_a$ is the average minor radius of a poloidal flux surface, and κ_a is proportional to $a^2 K / \tau_{SN}$. Integrating the conduction dominated steady-state total heat balance

$$-\frac{1}{\rho} \left(\frac{\partial}{\partial \rho} \rho \kappa \frac{\partial T}{\partial \rho} \right) = p_{\text{heat}}(\rho), \quad (3)$$

gives a generalization of the above model to

$$\tau_E = \tau_{SN} f_{\text{profile}} \quad (4)$$

For parabolic density profiles, fixed heating profiles, and $\kappa = \infty$ inside an assumed sawtooth inversion radius [8,9] at fractional minor radius

$$x_{\text{saw}} = q^{-1} = \{ [5 \times 10^6 a^2 B / (RI)] (1 + K^2) / 2 \}^{-1} \quad (5)$$

A good approximation to the results of this integration is

$$f_{\text{profile}} = (f_{OH}^{P_{OH}} + f_{NB}^{P_{\text{abs}}} + f_{EC}^{P_{EC}}) / P_{\text{tot}} ,$$

where P_{OH} , P_{abs} , and P_{EC} are the classical ohmic heating, absorbed neutral beam power, and electron cyclotron resonance heating power and $P_{tot} = P_{OH} + P_{abs} + P_{EC}$. The ohmic heating efficiency

$$f_{OH} = 0.3(1 + 0.7 q)^{1/2}(1 - x_{saw}^3) \quad (6)$$

is obtained by a fit to the results obtained from $P_{OH} \propto T_e^{3/2} \propto (1 - \rho^2/\rho_a^2)^{1.5\alpha_T}$ where $1.5\alpha_T = 0.7$ from Ejima et al., [10]. The neutral beam heating efficiency

$$f_{NB} = 0.6(1 - x_{saw}^3) \exp[-1.1 \times 10^{-19} n L f_{opac} (E_b, \text{keV}/A_b)^{-3/4}] \quad (7)$$

is obtained by an analytic fit to numerical integration of the pencil beam model for perpendicular injection [11,12]. The opacity correction factor

$$f_{opac} = [f_1 + 2^{-1/4} f_2 + 3^{-1/4} f_3] / (f_1 + 2^{-1} f_2 + 3^{-1} f_3) \quad (8)$$

is a linearized correction for the presence of partial energy fractions $f_1 + f_2 + f_3 = 1$ (by atomic current) in the beam. For nonperpendicular injection at an angle θ to the magnetic axis, the path length L is simply replaced by the distance

$$L = a\{[1 + 2(R/a) + (R/a)^2 \cos^2 \theta]^{1/2} - (R/a) \cos \theta\} \quad (9)$$

from the edge to the magnetic axis. The electron cyclotron resonance heating efficiency

$$f_{EC} = 1 - 1.2 \max(x_{saw}, x_{res}) + 0.2 [\max(x_{saw}, x_{res})]^3 \quad (10)$$

follows from assuming all the heating occurs at a fractional minor radius $x_{res} = \rho_{res}/(ak)^{1/2}$. We scale nonconductive losses by the square of the density normalized to the maximum obtainable density [7,13,14]

to obtain the final synoptic scaling formula

$$\tau_{SS} = \tau_{SN} f_{\text{profile}} f_{\text{Mu}}, \quad (11)$$

where

$$f_{\text{Mu}} = \frac{1}{1 + \alpha_{\text{Mu}} n^2 / n_{\text{MS}}^2}, \quad (12)$$

$$n_{\text{MS}} = \begin{cases} 1.5 \times 10^{20} \text{B}/(\text{Rq}) & \text{for gas puffing} \\ 2.1 \times 10^{20} \text{B}/(\text{Rq}) & \text{for pellet fuelling} \end{cases}, \quad (13)$$

and τ_{SN} is defined in Eq. (1). Typically, $\tau_{SS} \sim \tau_{SN}/2$. Here α_{Mu} is a fourth adjustable parameter to be determined from the confinement data. We used a convenient computer code for computing the thermal energy confinement time, τ_{SS} , for specified auxiliary heating (cf. Singer and Ku [15]). By the solution of the energy balance in Eq. (11), the a priori unknown temperature is eliminated from the energy balance, and the confinement time and temperature are both determined from a priori prescribed operational parameters. This code also computes the stored energy in fast ions for isotropic neutral injection into a uniform cylinder model, as needed for comparing with diamagnetic measurement of τ_E . We used the following method to determine the empirical constants - α (leading coefficient), c (K-scaling), t (controlling degradation of confinement with increasing temperature), and α_{Mu} (correction factor for edge radiation and recycling/convection; c.f. Singer et al. [7] for motivation for the functional form in Eqs. (12) and (13)). Using preliminary estimates $\alpha=0.3$, $t=1$, $\alpha_{\text{Mu}}=1$ (cf. Singer et al., [7]), we found $c=1$ minimized the variance between the model and experimental data on scaling of K (Nagami et al. [16]) at fixed $q_{\psi} \approx q^* \equiv [5 \times 10^6 a^2 \text{B}/(\text{RI})] [(1+K^2)/2] K^{1/2} (1+1.5 a^2/R^2)$, [17], from ohmic

and neutral beam heating in the D-III tokamak. This was done by stepping c in half-integral intervals (an accuracy approximately consistent with the scatter in the data).

We proceeded with a method designed to make maximum use of the scaling information in controlled parameter scans (which often exhibit apparently random deviations of order 15% or less in spite of the much larger deviation in τ_E ($\geq 30\%$) between different tokamak runs with nominally nearly identical global parameters). This method assumes the "temperature" and "density" scaling parameters t and α_{Mu} are universal constants of the complete data set. But the leading coefficient α is taken to be a random variable which has different values for each controlled parameter scan, the differences being due to different diagnostic methods, operator procedures, etc. Thus, we fix t and α_{Mu} and find the values of $\alpha_j^{\min}$ which individually minimize the variance

$$\sigma_j^2 = \sum_{i=1}^{M_j} \left[\frac{y(\vec{x}_{ij}, \alpha_j, t, \alpha_{Mu}; c) - Y_{ij}}{Y_{ij}} \right]^2 / M_j$$

separately between measured values, Y_{ij} and values computed from Eqs.

(1)-(3) at various points $\vec{x}_{ij} = (n, a, K, R, P_{tot}, I, Z_{eff})_{ij}$ for the M_j measurements in each parameter scan j . If the total of these minimized variances is small for a large data collection, then the assumption that there are unique values of t and α_{Mu} is reasonable. (For simplicity, we use the graphical estimates of Y_{ij} shown by the crosses in the figures described below, rather than the actual data shown by closed or open circles. To the extent that variance within a scan is smaller than variance between scans, as assumed here, this introduces negligible error.) Under this procedure the best pair (t, α_{Mu}) was $t = 1$ (searched in increments of 0.5) and $\alpha_{Mu} = 1.0$ (searched in increments of 0.25),

which gave an overall leading coefficient

$$\bar{\alpha} = \frac{1}{N} \sum_{j=1}^N \alpha_j^{\min}$$

and measure of variation

$$\sigma_{\alpha} = \left[\frac{1}{N-3} \sum_{j=1}^N (\alpha_j^{\min} - \bar{\alpha})^2 \right]^{1/2},$$

$$\alpha = 0.38 \pm 0.13,$$

for $N=15$ data sets [2,3,17-22]. Since the data in Figs. 2 and 3 are reported separately and not as part of an integrated scan, they are included separately in the treatment of this data. The fit was obtained by minimizing $\sigma^2 = \sum_{j=1}^N \sigma_j^2$ with respect to t and α_{Mu} . If truly representative of the overall data, the $\pm 40\%$ uncertainty in τ_E reflected in the variation in α would be comparable to or better than the deviations of the data from other confinement scalings (e.g., exhibited on log-log plots in [23,24]). The minimum ± 0.5 uncertainty in the exponent t determined here is probably of greater significance, but could not be explored in depth without a much more complicated statistical analysis.

The results are shown in Figs. 1-6. The solid curves in Figs. 1-5 show results computed with the best overall fit ($\alpha=0.38$; $t=\alpha_{Mu}=c=1$). The long-dashed curves in Figs. 1-4 show the results with the same fixed parameters but value of the random variable α_j determined individually by each parameter scan. Figure 6 shows that the K-scaling exponent $c=1$ gives the best fit to K-scaling data [16], with $\alpha=0.3$, $t=\alpha_{Mu}=1$.

For comparison, we show as short-dashed lines in Figs. 1-5, the

results calculated from Goldston's synoptic energy confinement model[3]*

$$\tau_{GS} = (\tau_{GN}^{-2} + \tau_{GW}^{-2})^{-1/2} ,$$

$$\tau_{GN} = 1.03 \times 10^{-21} \text{ s } a^{1.04} R^{2.04} q^{1/2} ,$$

$$\tau_{GW} = 1.37 \times 10^{-9} \text{ s } I^2 a^{-0.74} K R^{3.5} W_{tot}^{-1} ,$$

where $W_{tot} = P_{tot} \tau_{GS}$. This model added the eight adjustable parameters (given here in italics) to those previously used for circular ohmic plasmas at relatively high q (cf. Pfeiffer and Waltz [1]). All of the additional adjustable parameters in this model are to be constrained by a set of global energy confinement data. (In a systematic treatment of this type, the same data base would also have to determine the here-suppressed scaling of τ_{GN} with K and I and of τ_{GW} with density.) By contrast, the τ_{SS} model adds three parameters (c , t , and α_{Mu}) to be determined from global confinement data, and it evidently gives at least as good a fit to this data. Although additional assumptions may contribute to the fit obtained with this model (e.g., $x_{saw} = q^{-1}$, $\kappa \propto \rho$, Eq. 6 for f_{OH}), these are subject to verification with plasma profile shape data which supplement the data base on global energy confinement. Thus, while we have not performed a rigorous statistical analysis, it seems likely that the present model represents some progress in obtaining a statistically significant empirical fit to the data shown here. In our view, the need for so many arbitrary adjustable parameters in Goldston's and other similar models is primarily due to the arbitrary choice of $t=0$ for the temperature scaling exponent in the "neo-Alcator"

* This model approximately satisfies the dimensional constraint for a collisional, high- β plasma, but this constraint was not explicitly used in choosing the model.

scaling in Eq. (1), along with the neglect of above-mentioned physically motivated conclusions about the effects of changing heating profiles and approaching the density limit.

The application of the present model to additional data sets [23,24]* and to predictive analysis should be tempered by the following considerations. First, the scaling of nonconductive losses with the density limit is not meant to apply to plasma dominated by centrally peaked line radiation or bremsstrahlung. Major radiative losses from the inner half of the plasma volume should be subtracted from the input power before applying this model. Second, large sawteeth may enhance the transport well outside the inner q^{-1} fraction of the plasma, as indicated in the legend to Fig. 2. Third, the derivation of stored thermal energy for various heating profiles assumed zero pressure at the boundary. For ("H-mode") plasmas with a significant pressure pedestal at the boundary, at least a predictive model of the pressure stored "under the pedestal" is required to obtain the confinement enhancement. Fourth, the model is derived primarily to give a convenient description of existing data. Like any such model, it is strictly valid for predictive purposes only under restricted conditions determined by the limited utility of the dimensional constraint used and by the simplicity of the model used for nonconductive losses.

* We have not used the L-mode data base applied by Kaye and Goldston [23], because the data in this base are not sorted by parameter scan, so they might not give the most accurate possible estimate of the scaling exponents. Since these data are not accessible without an individual release by each of the contributors, it would also be much more difficult to confirm or extend our results than it is with the generally accessible literature cited below.

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Figure Captions

FIG. 1 Global thermal energy confinement times, τ_E , vs. line average electron density, $n = \bar{n}_e$, for the best overall fit to the data ($\alpha = 0.38$, $t = \alpha_{Mu} = c = 1$; solid lines) and for the same values $t = \alpha_{Mu} = c = 1$ for the value of α which best fits each parameter scan (computed to the nearest increment of 0.1; long-dashed lines), compared to the synoptic scaling of Goldston [3] computed using the code of Singer and Ku [15] (short-dashed lines) for (a) Alcator-C data from Fig. 1a of Blackwell *et al.* [2], upper abscissa in units of 10^{20} m^{-3} , and (b) and (c) D-III [16] and TFTR [18], respectively, lower abscissa in units of 10^{19} m^{-3} . In all cases in Figs. 1-6, we assumed $Z_{\text{eff}} = 2$ and took other parameters at the midrange of values quoted in the paper when specific numbers were not quoted with the figure from which the data were taken. Beam species mixes were assumed to be $f_1 = f_2 = f_3 = 1/3$ except where given in the data reference or (for D-III) quoted elsewhere [25]. Crosses (x) denote the approximations used instead of actual data (closed or open circles) when computing fits for long-dashed and solid lines.

FIG. 2 Change ΔW in stored energy between OH and ECRH resonance heating in the T-10 tokamak [19] vs. toroidal field, as in Fig. 1 for (a) $I = 0.17 \text{ MA}$ and (b) $I = 0.1 \text{ MA}$. Note that the estimation of the sawtooth inversion radius of Eq. (5) is not decisive. In fact, if we assume a larger region of profile flattening, e.g., $x_{\text{saw}} = 2^{1/2} q^{-1}$, the portion, which is insensitive to B_T , is widened and the fitting can be improved further (without changing the final results for α , t , α_{Mu} , and c).

FIG. 3 As a function of plasma current are plotted (a) for ISX-B [20],
 $\beta_T^{\text{dia}} = ((2/3) P_{\text{tot}} \tau_{\text{SS}} V^{-1} + p_{\text{fast-ion}}) (B^2/2\mu_0)^{-1}$, where $p_{\text{fast-ion}}$
 is the fast ion pressure computed for a uniform plasma [15]
 (solid and long-dashed lines, otherwise as in Fig. 1) and $\beta_T =$
 $(2/3) P_{\text{tot}} \tau_E V^{-1} (B^2/2\mu_0)^{-1}$ from Goldston's model [3,15] (short-
 dashed lines) and V is the plasma volume, and (b) ΔW as in Fig. 2
 for T-10 [19].

FIG. 4 (a) τ_E vs. total input power P_{tot} (OH plus absorbed beam) for
 TFTR [21], (b) $\tau_E^{\text{dia}} = \beta_T^{\text{dia}} (B^2/2\mu_0) V / P_{\text{tot}}$ for ISX-B [20] vs.
 total absorbed (beam only) power as in Fig. 3 and total thermal
 energy W vs. total input power P_{tot} (OH plus ECRH) in T-10 [19]
 for (c) $I = 0.27$ MA and (d) $I = 0.18$ MA.

FIG. 5 τ_E^{dia} as in Fig. 4(b) for ASDEX [22] at 14.5 keV/amu and 40 keV/amu
 full power beam injection energies. The two data points shown
 were not considered part of a controlled parameter scan and
 were treated separately when computing the best overall fit,
 $\alpha = 0.38$ (solid line).

FIG. 6 Diamagnetic beta β_T^{dia} (as in Fig. 4b) vs. elongation K for various
 values of K scaling exponent, c , for D-III [16] for the best
 fit α , $\alpha = 0.3$, and $t = \alpha_{\text{Mu}} = 1$.

Fig.1

Fig.2

Fig.3

Fig.4

Fig.5

Fig.6

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